



VNIVERSITAT
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**NON MINIMAL
EXTENSIONS FOR**

Dark Matter, Neutrino Masses, and Leptogenesis

**A Phenomenologist's
Cookbook**

Programa de Doctorat en Física - 3126
Departament de Física Teòrica
Institut de Física Corpuscular (CSIC-UV)
Director: Juan Herrero García



DRONA VATSYAYAN
València, 2025



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**Non-Minimal Extensions for Dark Matter,
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Non-Minimal Extensions for Dark Matter, Neutrino Masses, and Leptogenesis: A Phenomenologist's Cookbook

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This dissertation is original work, written solely for this purpose, and all the authors whose studies and publications contributed to it have been duly cited.

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Certifica:

Que la presente memoria, **“Non-Minimal Extensions for Dark Matter, Neutrino Masses, and Leptogenesis: A Phenomenologist’s Cookbook”**, ha sido realizada bajo su dirección en el Instituto de Física Corpuscular de València por **Drona Vatsyayan**, y constituye su Tesis para optar al grado de Doctor en Ciencias Físicas.

Y para que así conste, en cumplimiento de la legislación vigente, presenta en la Universitat de València la referida Tesis Doctoral, y firma el presente certificado, en

Valencia, Junio de 2025

Juan Andrés Herrero García

To Chitrangada

Entities must not be multiplied beyond necessity.

-William of Ockham

[Occam's razor]

*If three things are not enough to verify an affirmative proposition about things,
a fourth must be added and so on.*

-Walter Chatton

[Chatton's Anti-razor]

*Everything should be made as simple as possible,
but not simpler.*

-Albert Einstein

LIST OF PUBLICATIONS

This PhD thesis is based on the compendium of following publications:

- *Large neutrino mass in cosmology and keV sterile neutrino dark matter from a dark sector*, C. Benso, T. Schwetz and D. Vatsyayan, *J. Cosmol. Astropart. Phys.* **04** (2025) 054, [arXiv: 2410.23926](#) [1]
- *Asymmetries in extended dark sectors: A cogenesis scenario*, J. Herrero-García, G. Landini and D. Vatsyayan, *J. High Energ. Phys.* **05** (2023) 049, [arXiv: 2301.13238](#) [2]
- *Lowering the scale of fermion triplet leptogenesis with two Higgs doublets*, D. Vatsyayan and S. Goswami, *Phys. Rev. D* **107** (2023) 3, 035014, [arXiv: 2208.12011](#) [3]
- *Multi-component dark sectors: symmetries, asymmetries and conversions*, A. Bas i Beneito, J. Herrero-García and D. Vatsyayan, *J. High Energ. Phys.* **10** (2022) 075, [arXiv: 2207.02874](#) [4]

Other research works carried out during the course of PhD:

- *From Capture to Collapse: Revisiting Black Hole Formation by Fermionic Asymmetric Dark Matter in Neutron Stars*, S. Robles, D. Vatsyayan and G. Busoni, [arXiv: 2507.22881](#) [5]
- *Neutrino masses from new Weinberg-like operators: Phenomenology of TeV scalar multiplets*, A. Giarnetti, J. Herrero-García, S. Marciano, D. Meloni and D. Vatsyayan, *J. High Energ. Phys.* **05** (2024) 055, [arXiv: 2312.13356](#) [6]
- *Neutrino masses from new seesaw models: Low-scale variants and phenomenological implications*, A. Giarnetti, J. Herrero-García, S. Marciano, D. Meloni and D. Vatsyayan, *Eur. Phys. J. C* **84** (2024) 803, [arXiv: 2312.14119](#) [7]
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- *Symmetry origin of Baryon Asymmetry, Dark Matter and Neutrino Mass*, S. Bhattacharya, R. Roshan, A. Sil and D. Vatsyayan, *Phys. Rev. D* **106** (2022) 7, 075005, [arXiv:2105.06189](#) [9]

RESUMEN

El Modelo Estándar de la física de partículas (SM), a pesar de ser una teoría notablemente exitosa, deja muchas preguntas sin respuesta. Entre estas cuestiones abiertas se destacan el origen de las masas de los neutrinos, la explicación de la asimetría bariónica observada en el universo y la naturaleza de la materia oscura. A lo largo de los años, se han propuesto varias extensiones mínimas del SM, aunque hasta la fecha no se ha obtenido ninguna pista concreta sobre esta “nueva física”.

Esta tesis explora la fenomenología de extensiones no mínimas del SM que abordan las tres cuestiones mencionadas, centrándose en un sector oscuro rico y en mecanismos novedosos de generación de masas de neutrinos. Además, se investigan marcos unificados que pueden proporcionar un origen común para los neutrinos, la asimetría bariónica y la materia oscura. La tesis se presenta como un compendio de artículos publicados, y está estructurada en tres partes. En la parte introductoria, se revisan los marcos y mecanismos más comúnmente empleados más allá del Modelo Estándar (BSM), y por tanto sirve como material de contexto para los trabajos de investigación presentados en la segunda parte, que contiene los artículos publicados. La tercera parte consiste en un breve resumen de los objetivos, la metodología y los resultados de la tesis en español.

En la Parte I: Introducción, se revisan brevemente el contenido, la estructura y el desarrollo histórico del modelo estándar, así como sus problemas abiertos en el Capítulo 1. Se discute el marco de teoría de campos efectiva y su complementariedad con la construcción de modelos. Se destaca un apartado interesante titulado “Una receta para construir modelos”, donde se exponen los pasos clave para desarrollar un modelo de nueva física. Los tres capítulos siguientes están dedicados, respectivamente, a las masas de los neutrinos, la asimetría bariónica del universo y la materia oscura.

En el Capítulo 2, se revisan las oscilaciones de neutrinos y las mediciones del sector de neutrinos, resaltando las preguntas abiertas en el campo de la física de neutrinos. A continuación, se analiza el operador de Weinberg y los mecanismos seesaw estándar para la generación de masas de neutrinos. También se discuten variantes a baja escala del mecanismo seesaw tipo I y modelos radiativos de masas de neutrinos.

En el Capítulo 3, se presentan las observaciones y evidencias de la asimetría materia-antimateria. Se argumenta por qué dicha asimetría debe generarse dinámicamente.

Se revisan algunos mecanismos dinámicos de generación de asimetría, centrándose especialmente en la leptogénesis térmica, que está vinculada al mecanismo seesaw para masas de neutrinos. Se analiza en detalle el caso de leptogénesis con neutrinos pesados mano derecha —leptogénesis vainilla, se deriva la asimetría CP y su cota superior, las ecuaciones de Boltzmann relevantes y el factor de conversión de los procesos de esfalerones. También se revisan escenarios más allá de la leptogénesis vainilla.

Las observaciones y evidencias de la materia oscura se presentan en el Capítulo 4. Se discuten las propiedades generales que debería tener una partícula candidata a materia oscura. Luego se presenta el mecanismo de congelación térmica, el caso de las partículas masivas débilmente interactivas (WIMPs) y sus búsquedas experimentales. Dado el resultado nulo de las búsquedas de WIMPs, se exploran escenarios más allá, que incluyen el mecanismo de congelación por interacción débil (freeze-in) y las partículas masivas débilmente interactivas (FIMPs), el marco de sectores oscuros multicomponente y la materia oscura asimétrica. Finalmente, se revisa una amplia gama de candidatos a materia oscura en la sección denominada “zoológico de DM”.

La Parte II consiste en copias de artículos de investigación publicados en revistas con revisión por pares. Al presentar un tema en la Parte I, se menciona cómo está relacionado con los trabajos de investigación realizados durante el doctorado. Por tanto, estos trabajos están debidamente citados en el contenido de la tesis. Las investigaciones se pueden clasificar en tres categorías: Sector de Neutrinos, Sector Oscuro y Marcos Unificados.

En lo que respecta al sector de neutrinos, uno de los enfoques principales son los modelos seesaw extendidos. En particular, extensiones del seesaw tipo I con un sector oscuro ampliado permiten explicar las masas de los neutrinos mediante el mecanismo seesaw, la asimetría bariónica mediante leptogénesis térmica, mientras que el candidato a materia oscura reside en el sector oscuro conectado. También se discuten nuevos modelos seesaw que surgen en presencia de múltiples escalares y fermiones adicionales. Se investigan además modelos de materia oscura asimétrica y multicomponente, donde las asimetrías juegan un papel crucial en determinar la abundancia relictas, así como la naturaleza asintótica del candidato a materia oscura. Se analiza el papel de un sector oscuro rico y débilmente acoplado en la cosmología de neutrinos.

Un resumen y discusión de estos trabajos de investigación, junto con las conclusiones, se presentan en el Capítulo 5. Finalmente, se incluyen tres apéndices con las derivaciones detalladas de la lagrangiana efectiva del seesaw, los diagramas de lazos en leptogénesis y el formalismo de las ecuaciones de Boltzmann.

RESUM

El Model Estàndard de la física de partícules (SM), tot i ser una teoria extraordinàriament exitosa, deixa moltes preguntes sense resposta. Entre estes qüestions obertes destaquen l'origen de les masses dels neutrins, l'explicació de l'asimetria bariònica observada a l'univers i la naturalesa de la matèria fosca. Al llarg dels anys, s'han proposat diverses extensions mínimes del SM, però fins ara no s'ha obtingut cap indicatiu clar d'esta "nova física".

Aquesta tesi explora la fenomenologia d'extensions no mínimes del SM que aborden les tres qüestions esmentades, centrant-se en un sector fosc ric i en mecanismes innovadors de generació de massa de neutrins. A més, s'investiguen marcs unificats que poden proporcionar un origen comú per als neutrins, l'asimetria bariònica i la matèria fosca. La tesi es presenta com un compendi d'articles publicats i està estructurada en tres parts. En la part introductòria, es revisen els marcs i mecanismes més utilitzats més enllà del Model Estàndard (BSM), i per tant serveix com a material de context per als treballs d'investigació presentats en la segona part, que conté els articles publicats. La tercera part consisteix en un breu resum dels objectius, la metodologia i els resultats de la tesi en castellà.

A la Part I: Introducció, es revisa breument el contingut, l'estructura i el desenvolupament històric del model estàndard, així com els seus problemes oberts al Capítol 1. Es discuteix el marc de teoria de camps efectiva i la seua complementarietat amb la construcció de models. Cal destacar un apartat interessant titulat "Una recepta per a construir models", on s'exposen els passos clau per a desenvolupar un model de nova física. Els tres capítols següents es dediquen, respectivament, a les masses dels neutrins, l'asimetria bariònica de l'univers i la matèria fosca.

Al Capítol 2, es revisen les oscil·lacions de neutrins i les mesures del sector de neutrins, destacant les preguntes obertes dins del camp de la física de neutrins. A continuació, s'analitza l'operador de Weinberg i els mecanismes seesaw estàndard per a la generació de masses de neutrins. També es discuteixen variants a baixa escala del mecanisme seesaw de tipus I i models radiatius de massa de neutrins.

Al Capítol 3, es presenten les observacions i evidències de l'asimetria matèria-antimatèria. Es defensa per què aquesta asimetria hauria de generar-se dinàmicament. Es revisen alguns mecanismes dinàmics de generació d'asimetria, centrant-se especialment en

la leptogènesi tèrmica, que està vinculada al mecanisme seesaw per a la massa dels neutrins. S'analitza en detall el cas de leptogènesi amb neutrins destres pesats — leptogènesi vainilla, es deriva l'asimetria CP i el seu límit superior, les equacions de Boltzmann rellevants i el factor de conversió dels processos de esfalerons. També es revisen escenaris més enllà de la leptogènesi vainilla.

Les observacions i evidències de la matèria fosca es presenten al Capítol 4. Es discuteixen les propietats generals que hauria de tindre una partícula candidata a matèria fosca. A continuació es presenta el mecanisme de congelació tèrmica, el cas de les partícules massives feblement interactuants (WIMPs) i les seues recerques experimentals. Davant dels resultats nuls d'aquestes recerques, s'exploren escenaris alternatius, com ara el mecanisme de congelació per interacció feble (freeze-in) i les partícules massives feblement interactuants (FIMPs), el marc de sectors foscos multicomponents i la matèria fosca asimètrica. Finalment, es revisa una àmplia gamma de candidats a matèria fosca en la secció coneguda com a “zoo de DM”.

La Part II conté còpies dels articles d'investigació publicats en revistes amb revisió per parells. Quan s'introdueix un tema a la Part I, s'indica com està relacionat o ha motivat els treballs d'investigació realitzats durant el doctorat. Per tant, aquests treballs estan degudament citats al contingut de la tesi. Les investigacions es poden classificar en tres categories: Sector de Neutrins, Sector Fosc i Marcs Unificats.

Pel que fa al sector dels neutrins, un dels punts centrals són els models seesaw estesos. En particular, les extensions del seesaw tipus I amb un sector fosc ampliat permeten explicar les masses dels neutrins mitjançant el mecanisme seesaw, l'asimetria bariònica mitjançant leptogènesi tèrmica, mentre que el candidat a matèria fosca es troba al sector fosc connectat. També es discuteixen nous models seesaw que apareixen en presència de múltiples multiplets escalars i fermiònics. Es continuen investigant models de matèria fosca asimètrica i multicomponent, on les asimetries tenen un paper clau a l'hora de determinar l'abundància relíctica, així com la naturalesa asimptòtica del candidat a matèria fosca. També s'analitza el paper d'un sector fosc ric i feblement connectat en la cosmologia dels neutrins.

Un resum i discussió d'aquests treballs, junt amb les conclusions, es presenta al Capítol 5. Finalment, s'inclouen tres apèndixs amb les derivacions detallades de la lagrangiana efectiva del seesaw, els diagrames de llaç en leptogènesi i el formalisme de les equacions de Boltzmann.

PREFACE

The Standard Model of particle physics (SM), despite being a remarkably successful theory, leaves many questions unanswered. Notably among these open questions are the origin of neutrino masses, explanation for the observed baryon asymmetry of the universe and the nature of dark matter. Over the years in the past, several minimal extensions of the SM have been proposed, though we have not had any hint of this ‘new physics’.

This thesis explores the phenomenology of non-minimal extensions of the SM that address the three issues mentioned above by focusing on rich dark sector and novel mechanisms of neutrino mass generation. Additionally, it explores unified frameworks that can provide a common origin for neutrinos, baryon asymmetry and dark matter. It is presented as a *compendium* of published articles, and is structured in three parts. In the introductory part, we review the commonly employed beyond the standard model (BSM) frameworks, and mechanisms. Hence, it serves the background material for the research works presented in the second part containing the published articles. The third part is a brief summary of the objectives, methodology, and results of the thesis in Spanish.

In Part I: Introduction, we first briefly review the contents, structure and historical development of the standard model, as well as its open problems in Chapter 1. We discuss the effective field theory framework and how it is complementary to model building. An interesting highlight is ‘A recipe for model building’ that outlines the key steps of cooking a new physics model. The next three chapters are all dedicated to the neutrino masses, baryon asymmetry of the universe and dark matter, respectively.

In Chapter 2, we review neutrino oscillations and measurements from the neutrino sector, and highlight the open questions in the field of neutrino physics. We then discuss the Weinberg operator and the standard seesaw mechanisms for neutrino mass generation. Low-scale variants of the type-I seesaw mechanism and radiative neutrino mass models are also discussed.

In Chapter 3, we first present the observation and evidences for the matter-antimatter asymmetry. We then argue why such an asymmetry should be dynamically generated. Some dynamical mechanisms for generating such an asymmetry are reviewed, focusing particularly on thermal leptogenesis which is linked to the seesaw mechanism for

neutrino masses. We discuss in detail the case of leptogenesis with heavy right-handed neutrinos–Vanilla leptogenesis, derive the CP asymmetry and the upper bound on it, the relevant Boltzmann equations, and the sphaleron conversion factor. We further review the scenarios beyond vanilla leptogenesis.

The observations and evidences for dark matter are presented in Chapter 4. We discuss the generic features that a particle dark matter candidate should possess. We then discuss the thermal freeze-out mechanism, the case of Weakly Interacting Massive Particles (WIMPs) and their experimental searches. Given the null results from WIMP searches, we discuss the scenarios beyond WIMPs, these include the thermal freeze-in mechanism and Feebly Interacting Massive Particles (FIMPs), the framework of multi-component dark sectors and asymmetric dark matter. We finally review a wide range of dark matter candidates in the ‘DM Zoo’ section.

Part II consists of copies of research articles published in peer-reviewed journals. When a topic is introduced in Part I, we mention how it is linked or motivated the research works carried out during the course of PhD. Hence, these research works have been thoroughly cited in the thesis contents. The research works can be categorised into three categories: Neutrino Sector, Dark Sector and Unified frameworks.

Concerning the neutrino sector, a key focus is on extended seesaw models. In particular, extensions of type-I seesaw with an extended dark sector can account for neutrino masses via seesaw mechanism, baryon asymmetry is explained via thermal leptogenesis, while the dark matter candidate resides in the connected dark sector. We also discuss new seesaw models arising in the presence of additional scalar and fermion multiplets. We further investigate asymmetric and multi-component dark matter models, where asymmetries play a crucial role in shaping the relic abundance, as well as in determining the asymptotic nature of the dark matter candidate. The role of a feebly-connected rich dark sector in determining neutrino cosmology is examined.

A summary and discussion of these research works along with concluding remarks is presented in Chapter 5. Finally, we provide three appendices with detailed derivations of the seesaw effective Lagrangian, loop diagrams in leptogenesis, and the formalism of Boltzmann equations.

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As I write these acknowledgements after finishing this thesis, I cannot fail to look back at this incredible (and not so linear) journey into research and academia, a journey that began even before I started my PhD. I have many wonderful people I met along the way to thank for it.

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I would also like to thank Thomas Schwetz for the opportunity to carry out a research stay at KIT and for introducing me to the field of neutrino cosmology. I had a great learning experience thanks to his guidance and attention to detail. I greatly acknowledge the contribution of my many collaborators, without whom I could not have brought the projects to completion. I would also like to thank the members of the SOM group for various useful discussions and meetings, especially to Nuria, for being my tutor and all her help with the bureaucracy.

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Part I: Introduction

THE STANDARD MODEL AND BEYOND

In the minuscule fraction of time that modern humans have existed compared to the age of the universe, we have made remarkable progress in understanding the fundamental nature of the world around us, from the tiniest building blocks to the largest cosmic structures, spanning scales beyond imagination. Starting from ancient philosophical ideas, such as the Indian philosopher Kanada's concept of four indestructible *Paramanu* and Democritus's 5th century BCE idea of uncuttable and infinite *Atomos*, we have transitioned to a rigorous scientific framework. This framework is embodied in the Standard Model (SM) of Particle Physics, which describes all known elementary particles and fundamental forces except gravity. With the discovery of the Higgs boson in 2012, the SM now stands tall with 17 fundamental particles. Its predictions have been verified through numerous experiments, making it one of the most successful and comprehensive theories in physics.

Much like Mendeleev's 19th-century periodic table organized chemical elements based on shared properties and predicted undiscovered elements, particle physicists throughout the 20th century have classified elementary particles into families, and predicted new particles later confirmed experimentally. Today, physicists continue to predict new particles that will form an expanded Standard Model, aiming to resolve phenomena the current SM cannot explain. The 'Extension' in the thesis title precisely refers to this approach of going beyond our current formulation to incorporate new particles and theories, without causing any harm to the model that we already have, as it has been experimentally tested to a great precision.

In this chapter, we briefly review the contents, structure and historical development of the SM, as well the issues that plague it and therefore necessitate the existence of new physics Beyond the Standard Model (BSM).

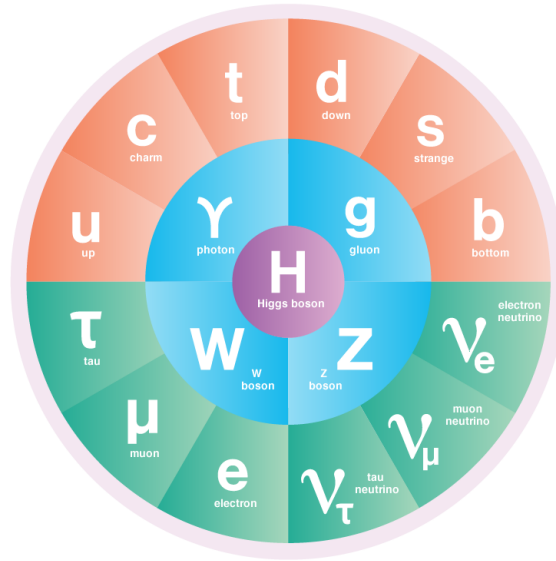


Figure 1.1: The Standard Model of Particle Physics: the fermions are shown in green (leptons) and red (quarks), the gauge bosons in blue and the Higgs boson in purple. Image courtesy: Symmetry magazine. Illustration by Sandbox Studios, Chicago.

1.1 A brief overview

If you are reading a thesis on ‘Non-Minimal Extensions’, you are probably aware of what the Standard Model of particle physics is. However, for completeness and to grasp how far we have come, it is worth mentioning the key highlights and historical developments. In order to understand these concepts in great detail, the reader is advised to refer to these books [10,11].

Based on the theoretical framework of Quantum Field Theory (QFT) with the gauge group of $SU(3)_C \times SU(2)_L \times U(1)_Y$, the Standard Model of particle physics consists of 17 fundamental particles. These include 12 spin-1/2 fermions, 4 different kinds of vector spin-1 bosons and 1 scalar spin-0 boson known as the Higgs boson (see Fig. 1.1) plus the corresponding antiparticles. These describe the matter content of our universe as well as three out of four fundamental forces: the strong force associated with the gauge symmetry $SU(3)_C$, and the weak and the electromagnetic force described by $SU(2)_L \times U(1)_Y$. In the following subsections, these are described in detail.

1.1.1 Particle content and gauge structure

Simply looking at Fig. 1.1, one may think that for a theory describing fundamental forces and elementary particles, it looks quite complex. Though if you tell this to particle physicists in the mid 20th century who were discovering new particles more often (look for ‘Particle zoo’), the SM is indeed a neat classification. In order to understand it, we have to look at the *players* of the game: the ‘particles’ and the *rules* of the game: the ‘gauge structure’.

The matter content of the Universe consists of fermions that follow Fermi-Dirac statis-

tics. In the SM, the fermions are divided into two categories: quarks and leptons, which are further classified in three generations (or families), i.e., they share the same properties but vary in mass and flavour. The matter that we see around us is made up of fermions belonging to the first generation (lightest) only.

Leptons: The leptons take part in all interactions except the strong interactions. The charged leptons (carrying an electric charge of $-1e$) are: electron, muon and tau, each grouped into one generation with their corresponding neutral leptons: electron neutrino, muon neutrino and tau neutrino. The heavier charged leptons decay rapidly into electrons and neutrinos. The electron, the first discovered lepton by J.J. Thompson in 1897, was incorporated into Niels Bohr's atomic theory (1910s). He proposed that these charged particles orbit around the nucleus, which was later found out to be made up of protons and neutrons.

Quarks: It was later discovered that hadrons such as protons and neutrons are not elementary but rather composite particles made up of quarks, thanks to the deep inelastic scattering (DIS) experiment at SLAC in 1968. However, the quark model had already been independently proposed by Murray Gell-Mann (see the 'Eightfold way' inspired from the Eightfold path in Buddhism [12]) and George Zweig in the early 1960s while studying the properties of hadrons. The quarks are the only elementary particles in SM that take part in all the four fundamental interactions. Further, they are also distinguished because of their six different flavours: up and down quarks (1st generation), charm and strange quarks (2nd generation), and top and bottom quarks (3rd generation). The most well-known baryons: protons and neutrons that are made of up and down quarks only.

Gauge bosons: Let us now have a look at the gauge structure of the SM: $SU(3)_C \times SU(2)_L \times U(1)_Y$. The SM is a gauge theory built on certain symmetry principles, or in other words, the physics remains invariant under these local transformations on the fields, known as gauge transformations¹. However, in order to keep the theory invariant under these transformations, one is forced to introduce the gauge fields which act as mediators of fundamental forces. As stated above, the SM describes three out of four fundamental forces/interactions, these are the strong force, the weak force and the electromagnetic force. Each of these interactions can be thought of as being mediated or facilitated by the exchange of their associated gauge bosons: gluons, $W^\pm$, Z bosons and photons.

The electromagnetic force and QED: Chronologically, the electromagnetic interactions were the first discovered of the three (gravitational force was discovered much earlier). The description was based on a simple $U(1)$ gauge theory known as Quantum Electrodynamics (QED). Its development starting in the 1920s, with the formulation of quantum theory by Dirac and later refinement by Feynman, Schwinger and Tomonaga

¹ The current usage of the term 'gauge invariance' is attributed to Weyl in 1929, see Ref. [13] for historical roots of gauge principles.

(for which they won the Nobel Prize in 1965) [14–16]. QED interactions are mediated by the massless photon A^μ (an outcome of gauge invariance), with the associated charge being the electric charge.

The strong force and QCD: These are the strongest of all four interactions, they involve quarks and are mediated by massless gluons. Just like quarks, ‘Gluon’ was also coined by Gell-Mann for acting like an adhesive to bind quarks in 1962. In 1979, they were finally detected. Unlike other SM particles, quarks do not exist freely and combine to form composite particles such as baryons (qqq) and mesons ($\bar{q}q$), collectively known as hadrons. This implied the existence of a new quantum number, called the ‘colour’ charge, where each quark species comes in three different colours. Since we do not observe free quarks, it is postulated that all the hadrons are colourless, i.e., colour singlet combinations. This phenomenon is known as colour confinement, and it prohibits the observability of free quarks carrying colour. The theory of strong interactions is thus governed by Quantum Chromodynamics (QCD), and is based on the non-Abelian gauge group $SU(3)_C$. It determines the interactions between the 6 flavoured quarks and the 8 gluons G_a^μ , ($a = 1, 2, \dots, 8$), corresponding to the generators of the group.

QCD was built on the template of QED and the Yang-Mills theory [17], where the colour charge of the former is analogous to the electric charge of the latter². However, due to the non-Abelian structure, we also obtain the interactions among the gluons themselves belonging to the adjoint representation of the colour group.

The weak force and unification: Now we come to the weak force (responsible for radioactive decay) mediated by the massive mediators: the charged $W^\pm$ bosons and the neutral Z boson, and therefore is a short range interaction. These gauge bosons were among the last to be discovered (CERN, 1983). Unlike the two cases above, the weak force was quite puzzling not just because of massive mediators, but especially due to the discovery of parity violation in the Wu experiment (1956). This discovery confirmed that the SM differentiates between the left and the right-handed fields, making it a chiral gauge theory. It was found that the left-handed fields Q, L corresponding to quarks and leptons should transform as doublets (the simplest group allowing for doublets is $SU(2)$). Meanwhile, the right-handed fields: u_R, d_R and e_R , corresponding to up-type, down-type quarks and charged leptons, respectively, should transform as singlets, leading to $SU(2)_L$. You will notice that there is no term corresponding to right-handed neutrinos, which can be taken as a hint of new physics (more on this later!).

Moreover, given that $SU(2)$ group has three generators, and we have three associated gauge bosons (the non-Abelian structure will lead to self-interactions among the bosons in this case as well), such a choice was well motivated. However, it was not

² Since we observed hadrons and not free quarks, historically the baryons were assigned integer electric charges. Upon discovering their constituents, the quarks were thus assigned fractional electrical charges (the only ones to carry such charges in SM).

enough to explain the observed weak phenomenon, where you have a neutral boson (Z) along with charged bosons ($W^\pm$) and $m_W \neq m_Z$. It was found that the weak interactions can be described by the using the product group $SU(2)_L \times U(1)_Y$ with their generators corresponding to weak isospin (T_1, T_2, T_3) and weak hypercharge Y and the associated gauge fields W_i^μ with $i = 1, 2, 3$ and B^μ , respectively. This structure yields four bosons: $W^\pm, Z, \gamma$ and unifies the weak interactions with the electromagnetic ones to yield the **Electroweak (EW) Theory** [18–20] for which Salam, Glashow and Weinberg won the Nobel Prize in 1979.

The Higgs boson: Finally, we come to the last piece of the SM particle content which was discovered in 2012, the Higgs boson, the only scalar boson in the theory. The Higgs boson is responsible for giving masses to other SM particles due to their interactions with the Higgs field. This is known as the **Higgs mechanism** and it deals with the spontaneous breaking of electroweak symmetry (see below). The long awaited discovery of the Higgs boson, proposed in the works of Higgs, Brout, Englert, Guralnik, Hagen and Kibble [21–25] by the ATLAS and CMS experiments at CERN [26, 27] was a major milestone in the field of particle physics.

1.1.2 The SM Lagrangian

The gauge sector contains the associated gauge fields of $SU(2)_L \times U(1)_Y$ that are W_i^μ with $i = 1, 2, 3$ and B^μ , and the gluons G_a^μ , ($a = 1, 2, \dots, 8$) of $SU(3)_C$ as mentioned above. The corresponding Lagrangian is

$$-\mathcal{L}_{\text{kin,gauge}} = \frac{1}{4}G_{\mu\nu}^a G_{\mu\nu}^a + \frac{1}{4}W_{\mu\nu}^i W_{\mu\nu}^i + \frac{1}{4}B_{\mu\nu} B_{\mu\nu}, \quad (1.1)$$

where the field strength tensors are given by

$$\begin{aligned} W_{\mu\nu}^i &= \partial_\mu W_\nu^i - \partial_\nu W_\mu^i + g\epsilon^{ijk}W_\mu^j W_\nu^k, \\ B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu, \\ G_{\mu\nu}^a &= \partial_\mu G_\nu^a - \partial_\nu G_\mu^a + g_s f^{abc}G_\mu^b G_\nu^c, \end{aligned} \quad (1.2)$$

where $i, j, k = 1, 2, 3$ and $a, b, c = 1, 2, \dots, 8$; g_s and g are the couplings associated with the $SU(3)_C$ and $SU(2)_L$, respectively. ϵ^{ijk} and f^{abc} are respectively the structure constants of $SU(3)_C$ and $SU(2)_L$. They satisfy the commutation relations $[\sigma_i, \sigma_j] = 2i\epsilon^{ijk}\sigma_k$, $[\lambda_a, \lambda_b] = if^{abc}\lambda_c$, where $\sigma(\lambda)$ are the Pauli (Gell-Mann) matrices corresponding to the generators of $SU(2)_L$ ($SU(3)_C$).

Coming to the Lagrangian associated with the fermions, we have

$$\mathcal{L}_{\text{kin,fermions}} = \sum_\alpha i\bar{\psi}^\alpha \gamma^\mu D_\mu \psi_\alpha, \quad (1.3)$$

where $\psi = Q, u_R, d_R, L, e_R$ denotes the fermion fields and γ^μ are the Dirac matrices,

Table 1.1: The fermionic and bosonic content of the SM and their charge assignments under the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$.

Fields				$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
Quarks	$Q_1 = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$Q_2 = \begin{pmatrix} c_L \\ s_L \end{pmatrix}$	$Q_3 = \begin{pmatrix} t_L \\ b_L \end{pmatrix}$	3	2	1/6
	u_R	c_R	t_R	3	1	2/3
	d_R	s_R	b_R	3	1	-1/3
Leptons	$L_1 = \begin{pmatrix} \nu_L^e \\ e_L \end{pmatrix}$	$L_2 = \begin{pmatrix} \nu_L^\mu \\ \mu_L \end{pmatrix}$	$L_3 = \begin{pmatrix} \nu_L^\tau \\ \tau_L \end{pmatrix}$	1	2	-1/2
	e_R	μ_R	τ_R	1	1	-1
Bosons		G^a		8	1	0
		W^i		1	3	0
		B		1	1	0
		H		1	2	1/2

with the covariant derivative

$$D_\mu = \partial_\mu - ig_s \frac{\lambda_a}{2} G_\mu^a - ig \frac{\sigma_i}{2} W_\mu^i - ig' Y B_\mu. \quad (1.4)$$

Here, g' is the coupling associated with $U(1)_Y$. Note how the covariant derivative term takes into account the interaction of matter (fermions) with the gauge bosons. Further, only those fermions that are charged under a gauge group take part in the interactions associated with that particular group. The charge assignments of the particles under the gauge groups are given in Table 1.1.

Importantly enough, you can spot that the mass terms are missing! (Also, you can ask where exactly is the Higgs?) Strange at first, it is not surprising as the gauge symmetry forbids writing such a mass term for the gauge bosons. This worked for photons and gluons, but not for massive W, Z that were predicted from the Fermi theory for beta decay. Moreover, given the different transformation for left and right-handed fermions, we do not have fermion mass terms as well which have the form $m_\psi(\psi_L^\dagger \psi_R + \psi_R^\dagger \psi_L)$. This shortcoming can be tackled only if this gauge symmetry is somehow broken. This is exactly where the Higgs field will come into play.

The Higgs mechanism

In order to accommodate the fermion and gauge boson masses, one must introduce the concept of Spontaneous Symmetry Breaking (SSB). It is associated with a symmetric Lagrangian with a non-symmetric vacuum. The idea is analogous to the alignment of spins of the ground state of a ferromagnet in some arbitrary direction, even though the entire system is rotationally invariant. Therefore, in a system with degenerate minima, the symmetry is said to have been spontaneously broken if one of the minima is arbitrarily chosen.

In the SM, the electroweak symmetry is spontaneously broken (EW SSB) via the Higgs

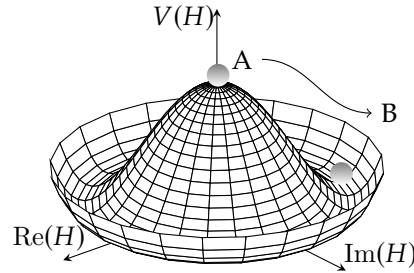


Figure 1.2: The Higgs mechanism: the Mexican hat shaped SM Higgs potential with infinitely many degenerate minima. The Higgs field is initially in a symmetric state (at point A) though this is not the true vacuum. It arbitrarily rolls down to the true vacuum $\langle H \rangle = v/\sqrt{2}$ choosing one of the infinite degenerate minima (at point B).

mechanism: where the Higgs field takes a non-zero Vacuum Expectation Value (VEV), breaking the electroweak symmetry to that of remnant QED

$$SU(2)_L \times U(1)_Y \xrightarrow{SSB} U(1)_{\text{QED}}. \quad (1.5)$$

Let us see how it works (Kibble talks about the history of electroweak symmetry breaking in Ref. [28]). We need to introduce a complex $SU(2)$ scalar doublet $H \equiv (H^+, H^0)^T$ governed by the Lagrangian

$$\mathcal{L}_{\text{Higgs}} = (D_\mu H)^\dagger D^\mu H + \mu^2 H^\dagger H - \lambda (H^\dagger H)^2, \quad (1.6)$$

where D_μ concerns only the electroweak part of eq. (1.4). Here $\mu^2 > 0$ will lead to a non-zero VEV when you try to minimize the scalar potential³

$$\frac{\partial V(H)}{\partial H} = 0 \Rightarrow \langle H \rangle = \sqrt{\mu^2/2\lambda} \equiv v/\sqrt{2}, \quad (1.7)$$

where $v \simeq 246$ GeV is the measured SM Higgs VEV extracted from G_F , the Fermi constant [29], breaking the symmetry of the potential, see Fig. 1.2. Now, for a perturbative approach, one can expand the scalar field about this particular vacuum: $\langle H \rangle \equiv (0, v/\sqrt{2})^T$. Since we introduced a complex scalar field (corresponding to four real scalar fields, say $\theta_{1,2,3}$ and h), the fluctuations can be parameterized as [25]

$$H(x) = \frac{v + h(x)}{\sqrt{2}} \exp \left[i \frac{\sigma_i}{2} \theta_i(x)/v \right] \begin{pmatrix} 0 \\ 1 \end{pmatrix}, \quad (1.8)$$

where the σ_i are once again the Pauli matrices ($i = 1, 2, 3$), and the exponential term in eq. (1.8) takes into account the $SU(2)$ transformations. Three out of the four scalar fields (θ_i) correspond to the three massless would-be Goldstone bosons which are eaten by the three gauge bosons ($W^\pm, Z$) to become their longitudinal degree of freedom. We are thus left with h ,⁴ the physical Higgs boson with its mass given by $m_h^2 =$

³ Note that $|H| = 0$ does not correspond to the true minimum.

⁴ One can choose $\theta_i(x) = 0$, so that only the physical h field remains, this is known as the unitary gauge.

$2\lambda v^2$ (you can obtain this by plugging the above equation in the scalar potential). The discovered Higgs boson's mass was found to be 125 GeV [26,27]. Note that the Higgs boson needs to be electrically neutral, because $U(1)_{\text{QED}}$ remains an exact symmetry of the vacuum.

Gauge boson masses: Focusing on the first term of the scalar Lagrangian in eq. (1.6) and the electroweak part of the covariant derivative in eq. (1.4), we get (in the unitary gauge)

$$\begin{aligned} (D_\mu H)^\dagger D^\mu H &= \left| \left(-ig \frac{\sigma_i}{2} W_\mu^i - ig' Y B_\mu \right) H \right|^2 \\ &= \frac{1}{8} \left| \begin{pmatrix} gW_\mu^3 + g'B_\mu & g(W_\mu^1 - iW_\mu^2) \\ g(W_\mu^1 + iW_\mu^2) & -gW_\mu^3 + g'B_\mu \end{pmatrix} \begin{pmatrix} 0 \\ v+h \end{pmatrix} \right|^2. \end{aligned} \quad (1.9)$$

Let us define $W^\pm = (W^1 \mp W^2)/\sqrt{2}$ and look at terms only proportional to v (the terms with h will give us interactions of the gauge bosons with the Higgs boson)

$$(D_\mu H)^\dagger D^\mu H \rightarrow \frac{g^2 v^2}{4} W_\mu^+ W^{\mu-} + \frac{1}{8} v^2 (W_\mu^3, B_\mu) \begin{pmatrix} g^2 & -gg' \\ gg' & g'^2 \end{pmatrix} \begin{pmatrix} W_\mu^3 \\ B_\mu \end{pmatrix}. \quad (1.10)$$

The first term above can be clearly identified as the mass of the charged $W^\pm$ boson:

$$m_W = g v/2. \quad (1.11)$$

The second term on the other hand corresponds to the mixing between W_μ^3 and B_μ fields. In order to obtain the gauge boson masses, we need to diagonalize the mixing matrix. The diagonalization of a mass matrix is among the key tools in particle physics models, something that we will use a lot more in the upcoming chapters. Coming back to the matrix, one can see that one of the eigenvalues is zero, which corresponds to the massless photon, while the other one is non-zero and gives us the Z-boson mass

$$m_A = 0, \quad m_Z = \frac{1}{2} v \sqrt{g^2 + g'^2}. \quad (1.12)$$

By defining $g'/g \equiv \tan \theta_W$ with θ_W being the weak mixing angle, we can write the mass eigenstates in terms of the diagonalizing matrix as

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W \\ -\sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} B_\mu \\ W_\mu^3 \end{pmatrix}, \quad (1.13)$$

corresponding the photon and Z boson fields. Comparing the masses of $W^\pm$ and Z bosons, we have $m_W/m_Z = \cos \theta_W$. Hence, in an elegant way, we have obtained the massive gauge bosons, as well as a massless one corresponding to the unbroken $U(1)_{\text{QED}}$ subgroup. Furthermore, if you now express the covariant derivative (see eq. (1.4)) in terms of the mass eigenstates, you can derive the relation between the

electric charge (QED) and the weak charges as $Q = T_3 + Y$. It also leads to the relation $e = g \sin \theta_W$.

Fermion masses: The same scalar field that we introduced above for inducing the gauge boson masses also has the correct quantum numbers (see Table 1.1) to couple to the left-handed fermion doublets and the right-handed singlets (so that they form a singlet). This leads us to the Yukawa Lagrangian

$$\mathcal{L}_{\text{Yukawa}} = -Y_L^{\alpha\beta} \bar{L}_\alpha H e_\beta - Y_u^{\alpha\beta} \bar{Q}_\alpha \tilde{H} u_R - Y_d^{\alpha\beta} \bar{Q}_\alpha H d_R + \text{H.c.}, \quad (1.14)$$

where α, β denote the flavour indices and $\tilde{H} \equiv i\sigma_2 H^*$. Note that due to the hypercharge assignments, there is no term that couples lepton to quarks with the Higgs. We only have the Higgs coupling to just the leptons or the quarks. Among them, it is possible to couple to different generations; hence we use the flavour indices. This implies that $Y^{\alpha\beta}$ are 3×3 matrices, which are not diagonal (we will discuss the flavour structure and mixing in the next section).

After EWSSB as discussed above, we can substitute $H = (0, (v + h)/\sqrt{2})^T$, which will result in the following

$$\mathcal{L}_{\text{Yukawa}} = -\frac{v + h}{\sqrt{2}} \left(Y_e \bar{e}_L e_R + Y_u \bar{u}_L u_R + Y_d \bar{d}_L d_R + \text{H.c.} \right), \quad (1.15)$$

where we have suppressed the flavour indices, rather one can read them as:

$$e = (e, \mu, \tau), \quad u = (u, d, s), \quad d = (d, s, b).$$

The terms proportional to v can be identified as the mass terms associated with the fermions: $M_f = Y_f v / \sqrt{2}$

$$-\mathcal{L}_{\text{mass,fermions}} = \bar{e}_L M_e e_R + \bar{u}_L M_u u_R + \bar{d}_L M_d d_R + \text{H.c.} \quad (1.16)$$

We see that all the fermion masses are proportional to same Higgs VEV, therefore the hierarchy in the masses of different fermion generations that we see is a consequence of the hierarchical Yukawa couplings, see Fig. 1.3. Once again, notice the absence of any mass term for the neutrinos, as there are no right-handed neutrino counterparts like the quarks and charged leptons.

Flavour structure

In the previous section, we saw how masses emerge due to the Higgs mechanism. However, we still have not determined the fermion mass eigenstates that would lead us to the observed fermion masses. Once again, we need to diagonalize the associated mass matrix, i.e., $M_f = Y_f v / \sqrt{2}$, due to the non-diagonal Yukawa couplings. In order to diagonalize the fermion mass matrix, we need to first rotate the involved fields (which

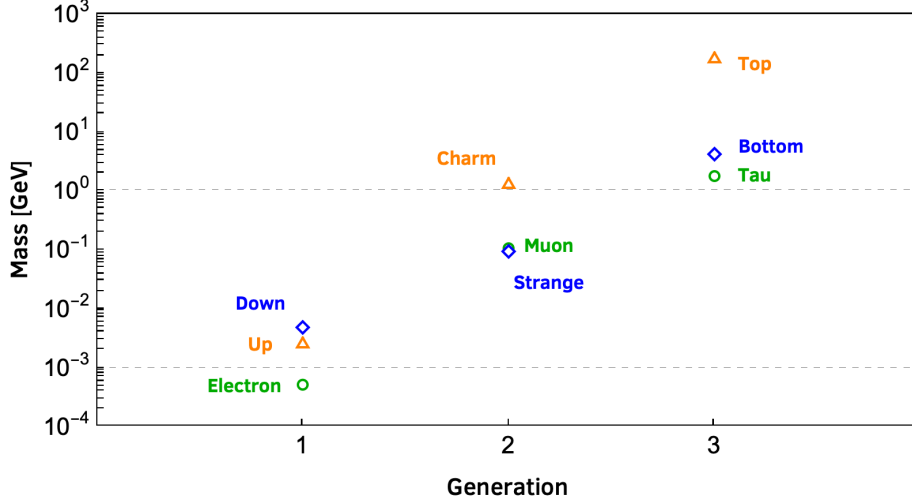


Figure 1.3: Fermion mass spectrum in the SM. The top quark with $m_t \sim 172$ GeV is the heaviest particle in the SM.

are currently in the *gauge basis*). We can perform the following rotation on the fields seen in eq. (1.15)

$$\begin{aligned} u'_L &\rightarrow \mathcal{U}_L u_L, & d'_L &\rightarrow \mathcal{D}_L d_L, & e'_L &\rightarrow \mathcal{E}_L e_L, \\ u'_R &\rightarrow \mathcal{U}_R u_R, & d'_R &\rightarrow \mathcal{D}_R d_R, & e'_R &\rightarrow \mathcal{E}_R e_R, \end{aligned} \quad (1.17)$$

where $\mathcal{U}, \mathcal{D}, \mathcal{E}$ are unitary matrices and will diagonalize the mass matrices. After applying the transformations in eq. (1.15) (focusing only on the terms $\propto v$ which give us the masses), we get

$$-\mathcal{L}_{\text{mass,fermions}} = \bar{e}'_L \mathcal{M}_e e'_R + \bar{u}'_L \mathcal{M}_u u'_R + \bar{d}'_L \mathcal{M}_d d'_R + \text{H.c.}, \quad (1.18)$$

where $\mathcal{M}_f$ are the mass diagonal matrices given by $\mathcal{M}_f = \mathcal{F}_L^\dagger M_f \mathcal{F}_R$ with $\mathcal{F} = (\mathcal{U}, \mathcal{D}, \mathcal{E})$, i.e., $\mathcal{M}_u = \text{diag}(m_u, m_c, m_t)$ and so on. Though the theory was written in the gauge basis, it is this mass basis that is of physical consequence and used for calculations.

The next question that arises is about the consequences of these transformations. For this, we need to go back to the kinetic terms associated with the fermions in eq. (1.3) which give us the interactions of fermions with the electroweak bosons post SSB

$$\begin{aligned} \mathcal{L}_{CC} &= \frac{g}{\sqrt{2}} W_\mu^+ (\bar{u}_L \gamma^\mu d_L + \bar{\nu}_L \gamma^\mu e_L) + \text{H.c.}, \\ \mathcal{L}_{NC} &= e Q_f A_\mu \bar{f} \gamma^\mu f + e Z_\mu \bar{f} \gamma^\mu \left[\frac{T_3^f}{\sin \theta_W \cos \theta_W} \frac{1}{2} (1 - \gamma_5) - Q_f \tan \theta_W \right] f, \end{aligned} \quad (1.19)$$

where f are the fermion fields. The first equation corresponds to the charged currents as they involve the W bosons whereas the second one corresponds to the neutral currents involving Z, A .

Once, we rotate to the mass eigenstates using eq. (1.17), in the quark sector, the $\mathcal{L}_{CC}$ takes the form

$$\mathcal{L}_{CC}^{\text{quarks}} = \frac{g}{\sqrt{2}} W_{\mu}^{+} \bar{u}'_L \gamma^{\mu} V_{CKM} d'_L + \text{H.c.}, \quad (1.20)$$

where $V_{CKM} \equiv \mathcal{U}_L^{\dagger} \mathcal{D}_L$ is a complex unitary matrix called the Cabibbo-Kobayashi-Mashikawa (CKM) matrix or the quark mixing matrix [30,31]

$$V_{CKM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix}. \quad (1.21)$$

This implies that the charged current interactions can change quark flavour at the tree level, with the strength of the flavour changing interactions determined by the elements of the CKM matrix above. The CKM being a 3×3 unitary matrix has 3 angles and 6 phases, five out of these can be removed by field redefinitions. Hence, at the end for three generations, we only have 3 angles ($\theta_{12}, \theta_{13}, \theta_{23}$) and one phase δ_{13} (see the standard parameterization of CKM matrix [29]). The latter is the only complex phase in the SM and a source of CP violation. Historically, Cabibbo's mixing matrix involved only 2 generations (1963). It was due to the observation of this CP violation in the Kaon systems that the existence of third generation quarks was predicted, and the matrix was extended by Kobayashi and Mashikawa (1973) before the discovery of bottom and top quarks.

Further, you can see that in the case of neutral currents, there is no flavour changing at the tree level. This is because the corresponding product of unitary rotation matrices is diagonal in gauge/flavour basis. Second, in the absence of neutrino masses, there is no such leptonic flavour mixing matrix because one can choose a transformation where the mass and gauge eigenstates can coincide (for example, by redefining the neutrino flavours). Hence, we do not have Lepton Flavour Violation (LFV) in the SM. Of course, as we will see later, in the presence of neutrino masses, we will obtain the lepton mixing matrix, called Pontecorvo-Maki-Nagakawa-Sakata (PMNS) matrix [32,33].

Putting it all together

Now, we have all the ingredients for the game, the rules and players; after putting it all together, the final SM renormalizable Lagrangian is given by

$$\mathcal{L}_{SM} = \mathcal{L}_{\text{kin,gauge}} + \mathcal{L}_{\text{kin,fermions}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yukawa}} \quad (1.22)$$

which are defined in eqs. (1.1), (1.3), (1.6) and (1.14). After electroweak symmetry breaking, we obtain the gauge boson and fermion masses, as well as their interactions with each other and the Higgs. The parameters that it contains are the 6 quark masses, 3 charged lepton masses (the fermion Yukawas can be determined from there), 4 CKM

parameters, 3 gauge couplings, the Higgs boson mass and VEV, and a strong CP violating phase (see below) for a total of 19 parameters. Additionally, the standard model also has **accidental symmetries** (as in the sense that they were not imposed by hand but emerge from the SM gauge structure and particle content) such as baryon and lepton number (LN) conservation. Furthermore, there is also a residual $SU(2)$ symmetry in the Higgs potential after SSB called the custodial symmetry. The custodial symmetry is exact in the absence of Yukawas and hypercharges, which leads to the following relation at the tree level [34,35]:

$$\rho = \frac{m_W^2}{M_Z^2 \cos^2 \theta_W} = 1. \quad (1.23)$$

The current global fits of Electroweak Precision Observables (EWPO) set $\rho = 1.00031 \pm 0.00019$ [29]. We observe $\rho \approx 1$ as a consequence of the accidental custodial symmetry.⁵ Therefore, the ρ parameter measures the deviations from this symmetry and provides constraints on new physics models. Note that these accidental symmetries may not be respected at higher scales.

Despite being complex, the standard model successfully explains a wide range of observed phenomena and accommodates well our present understanding of the strong and electroweak interactions. Yet, there remain challenges that it does not accommodate as well as unexplained patterns (if one seeks a deep understanding) as we discuss in the next section.

1.2 What does the SM fail to explain?

The SM has explained, up to a great accuracy, most of the observed low-energy electroweak phenomenon ($\sim O(100 \text{ GeV})$). These include the astonishing agreement between the predicted and experimentally measured value of the anomalous magnetic moment of the electron, the prediction and subsequent discoveries of Z, W bosons, third generation quarks (in particular the top mass), along with that of the Higgs bosons. These are some great accomplishments in the field of physics.

Along with Λ CDM, the standard model of cosmology which is based on the cosmological principle and the General Theory of Relativity (GTR), explains large scale cosmological and astrophysical observations. We have gathered a great amount of information about our universe over the years. Still, there are a lot of unsolved mysteries.⁶

The fields of particle physics and cosmology are rather interdisciplinary, and a fundamental theory should be able to describe all scales small and large (i.e., towards a ‘A Theory of Everything’). Therefore, we seek explanations for both the observed cosmological phenomenon from a particle viewpoint by looking at some issues within the

⁵ The custodial symmetry is not preserved when radiative corrections are included, i.e., $\rho \neq 1$.

⁶ The Λ CDM model has its own set of challenges which are not included in this section, for a recent review see Ref. [36]. Some of these will be discussed in subsequent chapters.

SM. In other words, we would like to unify the two standard models. Let us have a look at some of the problems that the SM fails to explain:

1. **Neutrino masses:** We saw that the neutrinos in the SM are massless, and they do not have any right-handed counterparts like the other fermions. After the observation of neutrino flavour oscillations that show that different flavours mix in the leptonic sector, (initially proposed by Pontecorvo in 1957 [37]), we now know that neutrinos have a very tiny mass (sub-eV scale), about 6 orders of magnitude below the electron mass. This is one of the most clear evidences of new physics, as there is no explanation within the SM for neutrino masses.
2. **Matter-antimatter asymmetry:** Everything that we see around is made up of matter, whereas antimatter exists in trace amounts. This unbalance termed as the 'Baryon asymmetry of the universe' is the reason of our existence. Such an asymmetry needs to be dynamically generated after the Big Bang, though, we do not know the mechanism behind it. In principle, the SM has all the ingredients to dynamically generate an asymmetry. However, these are not enough to explain the observed asymmetry, and we need to again extend the SM.
3. **Dark matter:** We know that in Λ CDM, the energy budget of the universe contains roughly 25% non-luminous and non-baryonic matter (5 times more than the visible matter that makes up everything that we see), aptly termed as Dark Matter (DM). Accommodating a DM candidate clearly requires extending the SM content. Despite significant evidence for dark matter, we have not been able to detect it yet and establish either its mass or nature.
4. **Dark energy:** The 70% of the energy budget of the universe is made up of dark energy due to which the universe undergoes an accelerated expansion, another unexplained phenomenon in the SM. As this acceleration is driven by a component with negative pressure, one of the simplest explanations is to assume a constant energy density characterized by the cosmological constant Λ (introduced by Einstein in 1917). Explaining the small value of Λ within a quantum field theory framework leads to a much larger value (120 orders of magnitude of difference!). There are other explanations such as dynamic dark energy [38, 39]. Recent data release from the DESI collaboration [40] hints at dynamical dark energy, with many interesting implications [41, 42].
5. **Gravity:** Let us not miss the elephant in the room: the SM offers no explanation for the fundamental force of gravity, and is incompatible with GTR. Unifying gravity with other interactions in the SM leads to non-renormalizable infinities. Coming up with a quantum theory of gravity (for example string theory or loop quantum gravity) is one of the most important theoretical challenges that physicists are working on. Similar to other gauge bosons, spin-2 mediators of gravity called *gravitons* have been proposed, which cannot be currently accounted for in

the SM gauge structure.

This thesis focuses on the first three items of the list above, namely neutrino masses, baryon asymmetry of the universe and dark matter. These are discussed in more detail in Chapters 2, 3 and 4, respectively. The problems mentioned above are not exhaustive, there are other issues with the SM as well. The ones mentioned above highlight certain shortcomings, particularly when it comes to what we have already observed. Moreover, on theoretical grounds, i.e., when trying to seek a deeper understanding of what we know, we have:

1. **Flavour puzzle:** We do not know the reason behind the flavour structure of quarks and leptons in the SM: why do we have three generations, why do we have the hierarchy in their masses and the mixing pattern. In fact, all fermion masses go as free parameters in the SM, and we believe there is an underlying structure yet to be unveiled to reduce the number of these free parameters and offer a rationale behind the observed distinct values. Note that the constraints on new physics are quite stringent, coming from the flavour sector.
2. **Strong CP problem:** In the QCD Lagrangian involving gluons and quarks, we can write the following CP violating term (consistent with all the SM gauge groups):

$$\mathcal{L} \supset -\frac{n_f g_s^2 \theta_{\text{QCD}}}{32\pi^2} G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}. \quad (1.24)$$

The θ parameter above is constrained by the neutron electric dipole moment to be tiny:

$$\bar{\theta} = \theta_{\text{QCD}} + \arg[\det(M)] \lesssim 10^{-10}, \quad (1.25)$$

where M is the quark mass matrix. Such a small value, in a way, represents a fine-tuning problem as there is no symmetry in the SM that would lead to this parameter being extremely tiny (usually $O(1)$ parameters are natural in the t'Hooft sense [43]). Thus, even though there is no reason for the CP symmetry to be necessarily conserved in the strong interactions, the strong sector seems to conserve it, and we do not observe such violation in experiments. This is known as the 'Strong-CP problem'. One of the most popular solutions is to impose an additional $U(1)$ symmetry called the **Peccei-Quinn symmetry** which is spontaneously broken, leading to a light pseudoscalar called the *axion* and setting $\theta_{\text{QCD}} \rightarrow 0$ [44–47]. Axions are very popular DM candidates these days (more on this later).

3. **Hierarchy problem:** We see from all the discussion above that new physics is required, however, this new physics scale (Λ_{NP}) is unknown. The Higgs mass in the SM is quite sensitive to high scales (there is no symmetry protecting its mass) and can receive large quadratic corrections depending on Λ_{NP} . Therefore, the observed mass of 125 GeV can again be seen as unnatural or fine-tuned in

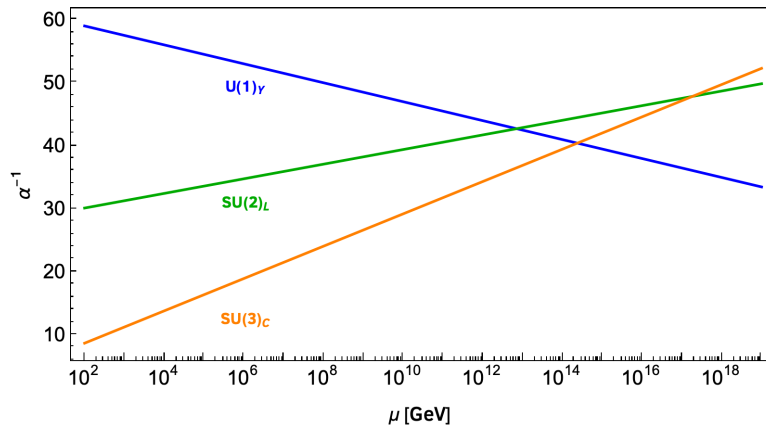


Figure 1.4: Running of the gauge couplings (α^{-1}) in the SM with the energy scale μ . It can be seen that orange and green curves coincide at $\mu \sim 10^{16}$ GeV; however, the three of them never exactly coincide together.

presence of this new physics, as it would require the bare mass term to be adjusted in such a way that one can get rid of these large quadratic corrections. This is called the hierarchy problem, based purely on theoretical grounds that depend on the Λ_{NP} .⁷

4. **Gauge coupling unification:** As mentioned above, we have three different gauge couplings corresponding to $SU(3)_C \times SU(2)_L \times U(1)_Y$. These couplings also vary with energy due to renormalization, also known as *running* with energy. It is believed that the SM gauge group should unify together to one single gauge symmetry characterized by one coupling constant, leading to the structure of a more fundamental theory termed as Grand Unified Theory (GUT). In the SM, the gauge couplings should unify at very high energy scales corresponding to 10^{16} GeV, however, they do not. We show the running of the three gauge couplings, $\alpha_i \equiv g_i^2/4\pi$, in the SM in Fig. 1.4.

5. **Charge quantization:** Remember that quarks carry fractional electric charges, in particular, in units of $1/3$ (see Table 1.1). We do not have any explanation for this assignment, or why fermion charges should come in discrete units of the electron's charge. This quantization maybe considered ad-hoc, though, it can also be linked to high-scale theory from which it can arise naturally.

It should be reiterated that the list above is not exhaustive. For example, there are problems such as the existence of Landau poles in QFT, which is linked with quantum triviality (see Refs. [48, 49]). These arise completely from a theory building perspective.

⁷ Note that we do have the Planck energy scale at 10^{19} GeV set by gravity, which in principle makes this problem worse.

1.3 Physics beyond the SM

So far, we have reviewed the Standard Model, acknowledged its successes, and recognized its shortcomings. We have conclusive evidence that it remains incomplete and requires extensions to address the questions it cannot currently answer. Such extensions are referred to as Beyond the Standard Model (BSM) theories, which aim to resolve one or more of the SM's unresolved issues. The term New Physics (NP) is also used interchangeably with BSM; however, BSM refers more specifically to theoretical frameworks, while NP serves as a broader umbrella term for any deviations from our current understanding of nature.

In general, such extensions involve enlarging the SM particle or symmetry content to involve new interactions, new couplings, and new quantum numbers to make a BSM model. However, this must not harm the SM we know and have tested already. With this idea of keeping the main predictions of SM almost intact, one has some creative liberty to implement new ideas and frameworks with the approach of looking outward. The final goal is not only to create models that can resolve the challenges but also the ones that can be tested with our current experimental capabilities, or rather probed in the future. Hence, the synergy between the theorists and experimentalists is vital for this quest for new physics. Identifying the correct extension(s) that bring us closer to a more complete understanding of nature remains a challenging task, though the finite set of experimental results serves as a guiding beacon for model builders.

Given the electroweak scale, it was expected that some new physics could lie at the TeV scale—a range currently accessible at the LHC, operating at approximately 13 TeV. **Supersymmetry** models (see Ref. [50] for a review) were predicted to produce a wide range of signatures within the reach of LHC. However, no such evidence has been found over the years, constraining parts of its viable parameter space. On the other hand, if the scale of new physics is much higher than the TeV scale, then it would remain beyond the reach of collider based experiments (it would also aggravate the hierarchy problem). In this case, one needs to look for indirect hints and turn to the cosmic laboratory, i.e., search for the implications of BSM dynamics in the early universe.

Since the SM describes phenomena at distance scales larger than $O(\text{TeV}^{-1})$, it is regarded as an Effective Field Theory (EFT). In this context, the SM is understood as a low-energy manifestation of a more fundamental theory, involving particles whose characteristic mass scales Λ_{NP} are much higher than those of SM. Let us review this concept below.

1.3.1 Effective Field Theories

One of the remarkable features of nature is the separation of energy scales, which allows us to treat physical processes occurring at different scales independently. This separation is the reason nuclear physics and particle physics exist as distinct disci-

plines. Nuclear physicists can carry out calculations and make predictions without needing to account for the substructure of elementary particles. In the context of QFTs, EFTs provide a framework to describe low-energy phenomena (or physics at any given scale) without requiring detailed knowledge of the full high-energy theory.

At low energies, the effects of heavier particles or unknown new physics can be described by higher dimensional (non-renormalizable) operators. These operators are suppressed by the scale of new physics (Λ_{NP}) or the cut-off scale of the theory. In other words, this implies that we have decoupled from the high energy scale, and the heavy particle will have negligible effects on the low energy observables. Hence, the SM itself can be viewed as an EFT valid at the electroweak scale, with its predictions and calculations in agreement with the experimental results. The SM effective Lagrangian can be written as

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \frac{c_5}{\Lambda_{\text{NP}}} \mathcal{O}^{d=5} + \frac{c_6}{\Lambda_{\text{NP}}^2} \mathcal{O}^{d=6} + \frac{c_7}{\Lambda_{\text{NP}}^3} \mathcal{O}^{d=7} + \dots, \quad (1.26)$$

where the c_i are dimensionless coefficients. It can be seen that the higher the dimension of the operator, the more suppressed it is. The scale Λ_{NP} corresponds to the mass scale of the heavy fields that are said to be *integrated out* (calculation of path integral only over the heavy fields). Therefore, we only have the light fields remaining in the effective Lagrangian. For example, consider a theory consisting of two real scalars $\phi(x)$ and $\Phi(x)$ with masses m_ϕ and M_Φ respectively, with $m_\phi \ll M_\Phi$. The Lagrangian can be written as

$$\mathcal{L}_{\text{full}}[\phi, \Phi] = \mathcal{L}_{\text{free}}[\phi] + \mathcal{L}_{\text{free}}[\Phi] + \mathcal{L}_{\text{int}}[\phi, \Phi] \quad (1.27)$$

where $\mathcal{L}_{\text{free}}$ includes the kinetic and mass terms for the respective fields and $\mathcal{L}_{\text{int}}$ describes the interaction between the two fields. The effective Lagrangian can be defined by integrating out the heavy scalar field Φ as [51]:

$$\begin{aligned} e^{iS_{\text{eff}}} &= \exp \left[i \int d^4x \mathcal{L}_{\text{eff}}(x) \right] = \int D\Phi D\Phi^\dagger \exp \left[i \int d^4x \mathcal{L}_{\text{full}}[\phi(x), \Phi(x)] \right] \\ &= \exp \left[i \int d^4x \mathcal{L}_{\text{full}}[\phi(x), \Phi(x) = 0] + \mathcal{O}\{\phi(x)\} \right], \end{aligned} \quad (1.28)$$

where the second term corresponds to the operator obtained after integrating out the heavy field Φ . Hence, it can be seen that the effective Lagrangian at low energies does not consist of the heavy fields. Once the integration has been carried out, and we have the complete description of the full theory and the EFT, we can determine the involved coefficients by *matching* both the high-scale and low-scale theories so that we recover the low energy predictions/observables.

The EFT framework has become an extremely powerful tool for BSM, and is complementary to top-down approaches, see Refs. [52–54] for a review. One can say that

simplified models are the ones that lie in between EFTs (systematic, general) and well-grounded theories describing the full UV-complete model (predictive, complex).

As we will see in the next chapter, the EFT framework will come in handy when discussing the neutrino mass generation mechanism. In particular, we will discuss the unique dimension-5 operator made up of the SM Higgs and lepton doublets, called the Weinberg operator [55]. This operator violates lepton number by two units (which being an accidental symmetry of the SM need not be respected by higher dimensional operators).

1.3.2 A recipe for model building

Having set the background, it is now time that we discuss the approach and path that BSM model builders take when extending the SM. Let us see how to cook a BSM model:

1. **[Food item] Select a problem to work on:** We have seen several shortcomings of the SM. The first step is to target one or more of these issues—whether they stem from unexplained phenomena or theoretical inconsistencies. You may choose to address multiple problems within a single model. This is analogous to deciding what dish you want to cook and eat.
2. **[Cuisine] Choose your approach:** This serves as your guiding principle for building the model, much like choosing a cuisine— Indian, Italian, Spanish, etc. The most common approach is based on symmetry arguments, as the SM itself is built upon symmetry principles. Enlarging the symmetry group can lead to unification schemes or entire classes of models, especially within supersymmetric frameworks. Alternatively, you can adopt a bottom-up perspective using the EFT framework. Other guiding ideas include *minimality* and *naturalness*, depending on your taste. This thesis focuses on the non-minimal approach.
3. **[Ingredients] Add new particles; extend the symmetries:** Next step is to extend the particle content and symmetry group of the SM. You can add new fermions or scalars, the symmetries can either be continuous (local or global) or discrete. The gauge bosons will be determined whether the symmetry is gauged or not.
4. **[Cooking I] Construct the Lagrangian:** Once you have made the additions, one needs to write all renormalizable terms allowed by the symmetries (both old and new). If you use the EFT approach, you will have to write the higher dimensional operators as well. Next, you have to choose if you want the symmetry to be spontaneously broken (i.e., new gauge bosons, the new scalar field takes a VEV) or be intact depending on the problem you are trying to solve.
5. **[Tasting and adjusting] Consistency check:** With a model in hand, it's essential to check for theoretical consistency. This includes constraints from perturbativity, unitarity, anomaly cancellation, vacuum stability, and sensitivity to fine-tuning (if naturalness is a guiding principle). If problems arise, you may need to adjust your ingredients—perhaps add more 'salt'—to fix the issues.

6. **[Cooking II] Phenomenology:** At this stage, your model is only half-cooked. Now you must explore its phenomenological implications. This involves studying new interactions (scatterings, decays, etc.), calculating the mass spectrum and performing any necessary diagonalization, and computing contributions from loop diagrams.
7. **[Serving] Predictions:** The final step is serving your model, you need to first constrain it with experimental constraints (current and future sensitivities), for example: electroweak precision tests, bounds from astrophysics and cosmology, to map the allowed parameter space. You should make predictions about the unique signatures of your model that can be tested either with our current set-up or in the future.

This general flow is commonly followed in constructing BSM models and has also been employed in developing the models presented in Part II of this thesis. Next, we will dig deeper into the issues of neutrino masses, baryon asymmetry and dark matter in the following chapters.

NEUTRINO MASSES

The elusive and mysterious nature of neutrinos has garnered the interest of physicists ever since their existence was conjectured by Pauli in 1930. Neutrino interactions are so feeble that they rarely interact with the surrounding matter, despite being present in great abundance in the universe. They take part in charged current interactions, where their flavour is determined from the easier-to-detect charged lepton's flavour. It took 26 years after being postulated for neutrinos to be discovered in the Cowan and Reines experiment (1956) [56]. They observed electron antineutrinos ($\bar{\nu}_e$) from a nuclear reactor by utilizing the inverse beta decay process ($\bar{\nu}_e + p \rightarrow n + e^+$); they aptly called this 'Project Poltergeist'.

As we saw in the last chapter, neutrinos are the only fermions that are considered to be massless in the SM. This is due to the absence of their right-handed counterparts, consequently prohibiting the Dirac mass. However, we now have evidence for tiny neutrino masses from several ongoing experiments on solar, atmospheric and reactor neutrinos that have confirmed the phenomenon of neutrino oscillations, indicating a mixing between flavour and mass eigenstates [57–60]. These measurements hint that neutrino masses are 6 orders of magnitude smaller than the electron mass, making the observation of neutrino masses not only a direct indication of BSM physics but also leading us to the question: Why are neutrino masses so tiny? In this chapter, we will briefly review neutrino oscillations and the parameters involved in them, open questions in neutrino physics and neutrino mass models.

2.1 Oscillations and measurements

In the quark sector, the flavour and mass eigenstates do not coincide as both up and down quarks are massive. If neutrinos were to carry a tiny mass, the same mismatch between the mass and flavour eigenstates would allow them to change their flavour while propagating. This is termed as 'neutrino oscillations'. It was first predicted by

Pontecorvo [37], who thought of neutrino-antineutrino oscillations analogous to the ones in neutral Kaons. Later, it was expanded by Maki, Nagakawa and Sakata [33] to include the muon neutrinos (discovered in 1962), and describe flavour oscillations.

In 1968, Bahcall and Davis's Homestake experiment designed to count neutrinos from the fusion in the sun only observed one third of the prediction; this was termed as the 'solar neutrino problem'. Further, in the GUTs that we discussed in the last chapter, it was predicted that the protons could decay into antineutrinos and meson. An experiment called Kamioka Neutron Decay Experiment (Kamiokande) was designed to look for these decays. The experiment did not find any signal of the proton decay, however, it observed fewer ν_μ than expected. In 1996, the experiment was upgraded to Super-K, and gave us the first evidence for atmospheric neutrino oscillations a couple of years later. In 2001, the Sudbury Neutrino Observatory (SNO) in Canada published their results for solar neutrino oscillations, solving the problem. Finally, KamLAND reported the evidence of antineutrino disappearance from reactor neutrinos, confirming neutrino oscillations from various sources [61]. Takaaki Kajita and Art McDonald from Super-K and SNO were finally awarded the Nobel Prize in 2015 for the discovery of neutrino oscillations. These evidences confirmed that neutrinos are indeed massive.

Hence, we also have a unitary mixing matrix in the lepton sector (analogous to the CKM matrix), called the PMNS matrix

$$U_{\text{PMNS}} = \begin{pmatrix} U_{e1} & U_{e2} & U_{e3} \\ U_{\mu1} & U_{\mu2} & U_{\mu3} \\ U_{\tau1} & U_{\tau2} & U_{\tau3} \end{pmatrix}. \quad (2.1)$$

The flavour and mass eigenstates related by $\nu_\alpha = U_{\alpha i} \nu_i$, with $\alpha = \{e, \mu, \tau\}$ and $i = \{1, 2, 3\}$. Of course, if neutrinos were massless, this matrix would be diagonal as we discussed before.¹ Just like the CKM, we can characterize this matrix by 3 mixing angles and one CP-violating Dirac phase δ for three generations. Additionally, as neutrinos are neutral elementary particles, they can very well be their own antiparticles, i.e., they could be Majorana fermions (discussed below). Therefore, in addition to the four parameters above, we also need two additional physical Majorana phases ($\alpha_{1,2}$) for Majorana neutrinos. Hence, the full transformation matrix is given by $U_{\text{PMNS}} = \mathcal{R} \times \text{diag}(e^{i\alpha_1/2}, e^{i\alpha_2/2}, 1)$, where

$$\mathcal{R} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix}, \quad (2.2)$$

with $c_{ij} \equiv \cos \theta_{ij}$, $s_{ij} \equiv \sin \theta_{ij}$.

¹ Additionally, the presence of mixing in the leptonic sector also implies charged lepton flavour violation (LFV), for example, processes such as $l_\alpha \rightarrow l_\beta \gamma$ with $\alpha, \beta = e, \mu, \tau$, which are absent in the SM. These processes that thus would appear in several BSM models are severely constrained, with the strongest bounds typically coming from $\text{Br}(\mu \rightarrow e \gamma) < 4.2 \times 10^{-13}$ (90% CL) [62].

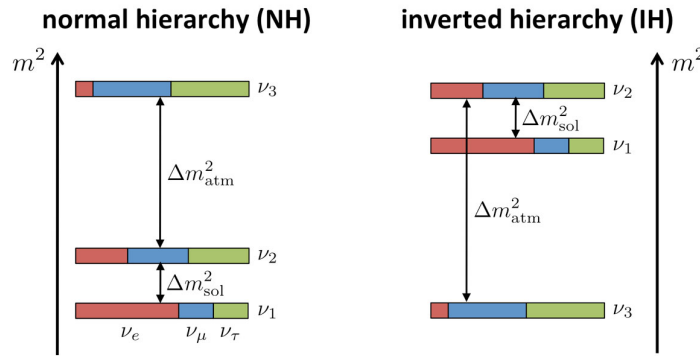


Figure 2.1: Two possible choices for neutrino mass ordering/hierarchy. The different colours correspond to the three lepton flavours, with their respective lengths proportional to the mixing parameters. Image courtesy: JUNO Collaboration.

Let us now discuss the oscillation probabilities. Considering first just the two-flavoured scenario as in the MNS paper [33], the flavour eigenstates can be written in terms of two mass eigenstates and one mixing angle θ

$$\begin{aligned} \nu_\alpha &= \nu_1 \cos \theta + \nu_2 \sin \theta, \\ \nu_\beta &= \nu_2 \cos \theta - \nu_1 \sin \theta. \end{aligned} \quad (2.3)$$

The probability (in vacuum) that a neutrino in a particular flavour state remains in the same flavour state after travelling a distance L (in km) is given by

$$P(\nu_\alpha \rightarrow \nu_\alpha) = 1 - \sin^2 2\theta \sin^2 \left(\frac{1.27 \Delta m^2 L}{E} \right), \quad (2.4)$$

where E (taken in GeV) is the neutrino beam energy, and $\Delta m^2_{21} = m^2_2 - m^2_1$ (in units of eV^2). If neutrinos are produced in a flavour eigenstate α at the source, and one detects fewer ν_α at the detector, the deficit is determined by the probability above. Such experiments/events are termed as 'disappearance'.

Alternatively, if a beam of ν_α is produced but one detects ν_β instead in an experiment, these are called as 'appearance' events. The probability that flavour state changes can be simply found by $P(\nu_\alpha \rightarrow \nu_\beta) = 1 - P(\nu_\alpha \rightarrow \nu_\alpha)$.

We can now extend this to the three flavoured case in the SM by using the PMNS matrix as

$$\begin{aligned} P(\nu_\alpha \rightarrow \nu_\beta) &= \delta_{\alpha\beta} - 4 \sum_{i>j} \text{Re}(U_{\alpha i}^* U_{\alpha j} U_{\beta i} U_{\beta j}^*) \sin^2 \left(\frac{\Delta m_{ij}^2 L}{4E} \right) \\ &\quad + 2 \sum_{i>j} \text{Im}(U_{\alpha i}^* U_{\alpha j} U_{\beta i} U_{\beta j}^*) \sin^2 \left(\frac{\Delta m_{ij}^2 L}{2E} \right). \end{aligned} \quad (2.5)$$

These equations show that measurements from neutrino oscillation experiments are

NuFIT 6.0 (2024)					
IC19 without SK atmospheric data		Normal Ordering ($\Delta\chi^2 = 0.6$)		Inverted Ordering (best fit)	
		bfp $\pm 1\sigma$	3σ range	bfp $\pm 1\sigma$	3σ range
	$\sin^2 \theta_{12}$	$0.307^{+0.012}_{-0.011}$	$0.275 \rightarrow 0.345$	$0.308^{+0.012}_{-0.011}$	$0.275 \rightarrow 0.345$
	$\theta_{12}/^\circ$	$33.68^{+0.73}_{-0.70}$	$31.63 \rightarrow 35.95$	$33.68^{+0.73}_{-0.70}$	$31.63 \rightarrow 35.95$
	$\sin^2 \theta_{23}$	$0.561^{+0.012}_{-0.015}$	$0.430 \rightarrow 0.596$	$0.562^{+0.012}_{-0.015}$	$0.437 \rightarrow 0.597$
	$\theta_{23}/^\circ$	$48.5^{+0.7}_{-0.9}$	$41.0 \rightarrow 50.5$	$48.6^{+0.7}_{-0.9}$	$41.4 \rightarrow 50.6$
	$\sin^2 \theta_{13}$	$0.02195^{+0.00054}_{-0.00058}$	$0.02023 \rightarrow 0.02376$	$0.02224^{+0.00056}_{-0.00057}$	$0.02053 \rightarrow 0.02397$
	$\theta_{13}/^\circ$	$8.52^{+0.11}_{-0.11}$	$8.18 \rightarrow 8.87$	$8.58^{+0.11}_{-0.11}$	$8.24 \rightarrow 8.91$
	$\delta_{\text{CP}}/^\circ$	177^{+19}_{-20}	$96 \rightarrow 422$	285^{+25}_{-28}	$201 \rightarrow 348$
	$\frac{\Delta m_{21}^2}{10^{-5} \text{ eV}^2}$	$7.49^{+0.19}_{-0.19}$	$6.92 \rightarrow 8.05$	$7.49^{+0.19}_{-0.19}$	$6.92 \rightarrow 8.05$
	$\frac{\Delta m_{3\ell}^2}{10^{-3} \text{ eV}^2}$	$+2.534^{+0.025}_{-0.023}$	$+2.463 \rightarrow +2.606$	$-2.510^{+0.024}_{-0.025}$	$-2.584 \rightarrow -2.438$
IC24 with SK atmospheric data		Normal Ordering (best fit)		Inverted Ordering ($\Delta\chi^2 = 6.1$)	
		bfp $\pm 1\sigma$	3σ range	bfp $\pm 1\sigma$	3σ range
	$\sin^2 \theta_{12}$	$0.308^{+0.012}_{-0.011}$	$0.275 \rightarrow 0.345$	$0.308^{+0.012}_{-0.011}$	$0.275 \rightarrow 0.345$
	$\theta_{12}/^\circ$	$33.68^{+0.73}_{-0.70}$	$31.63 \rightarrow 35.95$	$33.68^{+0.73}_{-0.70}$	$31.63 \rightarrow 35.95$
	$\sin^2 \theta_{23}$	$0.470^{+0.017}_{-0.013}$	$0.435 \rightarrow 0.585$	$0.550^{+0.012}_{-0.015}$	$0.440 \rightarrow 0.584$
	$\theta_{23}/^\circ$	$43.3^{+1.0}_{-0.8}$	$41.3 \rightarrow 49.9$	$47.9^{+0.7}_{-0.9}$	$41.5 \rightarrow 49.8$
	$\sin^2 \theta_{13}$	$0.02215^{+0.00056}_{-0.00058}$	$0.02030 \rightarrow 0.02388$	$0.02231^{+0.00056}_{-0.00056}$	$0.02060 \rightarrow 0.02409$
	$\theta_{13}/^\circ$	$8.56^{+0.11}_{-0.11}$	$8.19 \rightarrow 8.89$	$8.59^{+0.11}_{-0.11}$	$8.25 \rightarrow 8.93$
	$\delta_{\text{CP}}/^\circ$	212^{+26}_{-41}	$124 \rightarrow 364$	274^{+22}_{-25}	$201 \rightarrow 335$
	$\frac{\Delta m_{21}^2}{10^{-5} \text{ eV}^2}$	$7.49^{+0.19}_{-0.19}$	$6.92 \rightarrow 8.05$	$7.49^{+0.19}_{-0.19}$	$6.92 \rightarrow 8.05$
	$\frac{\Delta m_{3\ell}^2}{10^{-3} \text{ eV}^2}$	$+2.513^{+0.021}_{-0.019}$	$+2.451 \rightarrow +2.578$	$-2.484^{+0.020}_{-0.020}$	$-2.547 \rightarrow -2.421$

Figure 2.2: Latest results of the global fit of three-flavour oscillations from the NuFIT collaboration [66]. The lower (upper) section includes (does not include) the tabulated χ^2 atmospheric data from the Super-K and IceCube collaborations [67, 68].

sensitive to the mixing angles and the mass-squared differences. However, they do not provide any information about the absolute neutrino mass scale. So far, all three mixing angles have been measured: the solar mixing angle θ_{12} , the atmospheric mixing angle θ_{23} , and the smallest—and most recently measured—angle, θ_{13} , determined from reactor neutrino experiments [63–65].

We also know two mass splittings: the positive solar mass-squared difference, $\Delta m_{21}^2 \equiv \Delta m_{\text{sol}}^2 \sim \mathcal{O}(10^{-5}) \text{ eV}^2$, and the atmospheric mass-squared difference, $\Delta m_{3l}^2 \equiv \Delta m_{\text{atm}}^2 \sim \mathcal{O}(10^{-3}) \text{ eV}^2$, whose sign remains unknown. This ambiguity leaves us with two possible neutrino mass orderings (see Fig. 2.1):

- **Normal Ordering (NO):** $m_1 \lesssim m_2 < m_3$
- **Inverted Ordering (IO):** $m_3 < m_1 \lesssim m_2$

These are also known as hierarchies. Note that $\Delta m_{3l}^2 = \Delta m_{31}^2$ for NO and $\Delta m_{3l}^2 = \Delta m_{32}^2$ for IO. Depending on the mass ordering, the global fits for the different neutrino parameters differ; in Fig. 2.2, we show the latest results from the NuFIT collaboration. [66].

2.1.1 Laboratory versus Cosmology

Even though these experiments and measurements have established the existence of tiny neutrino masses, they only confirm that at least two neutrinos are massive. We still have no information about the mass of the lightest neutrino, which is key to determining the absolute neutrino mass scale. However, by assuming $m_{\text{lightest}} \rightarrow 0$, we can place lower bounds on the sum of neutrino masses using results from oscillation experiments:

$$\sum_{i=1}^3 m_i = \begin{cases} m_0 + \sqrt{\Delta m_{21}^2 + m_0^2} + \sqrt{\Delta m_{31}^2 + m_0^2} > 0.058 \text{ eV} & \text{(NO)} \\ m_0 + \sqrt{|\Delta m_{32}^2| + m_0^2} + \sqrt{|\Delta m_{32}^2| - \Delta m_{21}^2 + m_0^2} > 0.098 \text{ eV} & \text{(IO)}. \end{cases} \quad (2.6)$$

There are experiments designed to search for the absolute neutrino mass scale or the effective electron neutrino mass. One such experiment is KATRIN, which seeks to measure the mass of the electron neutrino through precise observations of the endpoint spectrum in tritium beta decay. The latest experimental result gives us the following upper bound [69]:

$$m_\beta = \left(\sum_{i=1}^3 |U_{ei}|^2 m_i^2 \right)^{1/2} < 0.45 \text{ eV (90\% CL)}. \quad (2.7)$$

KATRIN is designed to achieve a final sensitivity goal of 0.2 eV. Project 8, another experiment designed to measure the absolute mass scale (under development) aims to lower this limit to 0.04 eV [70]. Moreover, if neutrinos are Majorana particles as mentioned above, there could be a process called Neutrinoless Double Beta Decay ($0\nu\beta\beta$) ($0\nu\beta\beta$): $(A, Z) \rightarrow (A, Z + 2) + 2e^-$ as an interesting experimental consequence. The amplitude of this process is sensitive to the effective Majorana mass,

$$m_{\beta\beta} = \left| \sum_{i=1}^3 U_{ei}^2 m_i \right|, \quad (2.8)$$

which also involves the Majorana phases. Such a process has not yet been observed, meaning we still lack experimental confirmation of the Majorana nature of neutrinos, i.e., that neutrinos are their own antiparticles. Several ongoing and upcoming experiments are actively searching for this signature, including KamLAND-ZEN [71], CUORE [72], LEGEND [73], NEXT [74], and the planned nEXO experiment [75]. Currently, the most stringent limits on the effective Majorana mass come from KamLAND-ZEN [71]:

$$m_{\beta\beta} < 0.028 - 0.122 \text{ eV (90\% CL)}, \quad (2.9)$$

where the uncertainty majorly stems from the calculation of nuclear matrix elements.

The experiments mentioned above fall under the category of *laboratory* experiments.

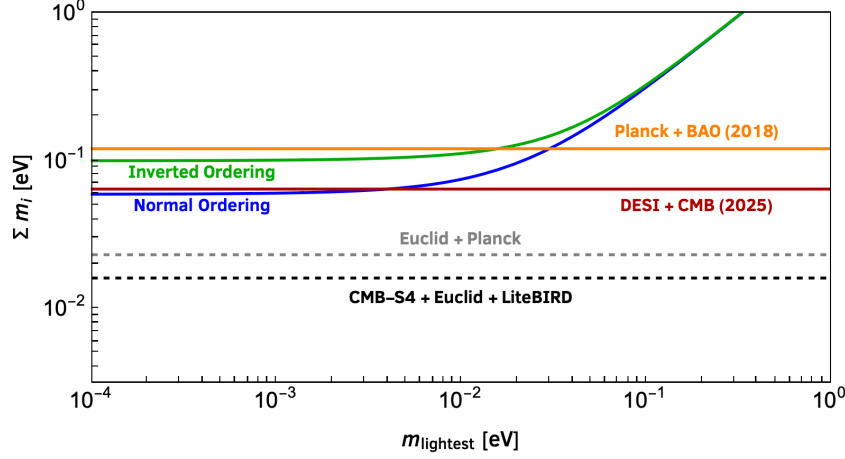


Figure 2.3: Comparison of the bounds on the sum of neutrino masses depending on the lightest neutrino mass from laboratory (using NuFIT data without including SK data [66]) and cosmology. The dashed lines correspond to the sensitivities of the future experiments CMB-S4 and Euclid [79].

However, massive neutrinos also leave an imprint on cosmological observables, for instance, by contributing to the total energy density of the universe and influencing the formation and distribution of large-scale structures (see Ref. [76] for a review). As a result, we also obtain constraints on the sum of neutrino masses from cosmological observations, which are significantly more stringent than those from laboratory experiments. These bounds, however, depend on the underlying cosmological model (typically Λ CDM) and the datasets used. For example, we have the following upper bounds on the sum of neutrino masses from cosmology:

$$\sum_{i=1}^3 m_i < \begin{cases} 0.12 \text{ eV (95\%CL)} & \text{Planck CMB+BAO [77]} \\ 0.064 \text{ eV(95\%CL)} & \text{DESI2025 + CMB [78]} \end{cases} \quad (2.10)$$

These upper bounds are largely in tension with the bounds from the neutrino oscillation data shown above in eq. (2.6). If one takes the latest DESI bound at face value, we see that IO is already excluded at 2σ level, and the constraint is approaching the limit for NO. Furthermore, the bounds in eq. (2.6) are derived under the assumption that the lightest neutrino mass is negligible. If the lightest neutrino has a non-zero mass, the tension becomes even more pronounced. Future experiments such as Euclid and CMB-S4 are expected to tighten the cosmological bound further, potentially down to ~ 0.02 eV [79,80]. The constraints are compared in Fig. 2.3. Consequently, existing cosmological bounds would make any compelling measurements of finite neutrino masses from our current laboratory experiments impossible. This discord is what we term as ‘Laboratory vs. Cosmology’.

However, these cosmological bounds are not directly sensitive to neutrino masses themselves, but rather to their contribution to the total energy density. As a result, they can be relaxed in non-standard cosmological scenarios, collectively referred to as *large neutrino mass cosmologies*. In Ref. [1], we present a model for relaxing the cosmological

neutrino mass bound. In this model, we utilized the seesaw mechanism (discussed below) and the presence of an extended dark sector, which we will discuss in Chapter 4.

2.1.2 What is still not known?

Despite the tremendous success of neutrino oscillation experiments in the past few decades, the discussion above has already highlighted a few gaps in our understanding of the neutrino sector:

1. **Mass ordering:** Unlike the charged lepton mass spectrum, we do not know the neutrino mass spectrum. Precise measurement of the sign of Δm_{3l}^2 would help us determine the mass ordering. This is one of the main goals of the JUNO experiment in China (expected to start taking data in 2025) [81,82]. Experiments under construction such as DUNE [83,84] and Hyper-Kamiokande [85] also aim to resolve the mass ordering issue.
2. **Absolute mass scale:** We are probing the absolute neutrino mass scale in experiments such as KATRIN as well as from $0\nu\beta\beta$ experiments. Experiments currently running and in the future, for example, Project 8 and PTOLEMY [86] have projected sensitivities $< O(0.1 \text{ eV})$.
3. **Nature of neutrinos:** Whether neutrinos are Dirac or Majorana particles continues to be a pressing issue. We hope to obtain a positive signal from the $0\nu\beta\beta$ experiments mentioned above. The implications of whether neutrinos are Dirac or Majorana will be discussed in detail in the next section.
4. **CP violation:** Our current knowledge of CP violation in the leptonic sector is still incomplete. The measurement of CP violating phase δ_{CP} has large uncertainties (see Fig. 2.2) which do not conclusively rule out the CP conservation in the leptonic sector. Experiments such as T2K [87] and NOvA [88] are currently taking data to determine strong constraints on δ_{CP} .
5. **Sterile neutrinos:** Apart from the SM neutrinos that we call *active*, BSM models also predict the existence of *sterile* neutrinos. These sterile neutrinos do not take part in the weak interactions like the active ones but can mix with them and contribute to oscillations. Currently, we do not know whether sterile neutrinos exist or what their mass scale might be. Their masses could range widely from the eV scale up to the GUT scale, depending on the theoretical motivation. The minimal model containing a single eV-scale sterile neutrino is disfavored by cosmology and MicroBooNE results [89], but searches for heavier sterile neutrinos are ongoing. These belong to the category known as Heavy Neutral Leptons (HNLs).

2.2 Dirac or Majorana?

All SM fermions, except neutrinos, carry an electric charge. This charge allows us to distinguish particles from their antiparticles, which carry equal but opposite charges. Moreover, these charged fermions are Dirac fermions, characterized by a Dirac mass

term $m_D \bar{\psi}_R \psi_L$ that connects left- and right-handed fields and conserves lepton number. In total, a Dirac fermion has four degrees of freedom: two for particle and antiparticle, and two for their helicities.

Neutrinos differ in two important ways: they are electrically neutral, and right-handed neutrinos are absent in the SM. Hence, it is possible that they are Majorana fermions and are their own antiparticles satisfying the Majorana condition: $\psi = \psi^c \equiv C \bar{\psi}^T$, where C is the charge conjugation operator [90]. Unlike Dirac fermions, Majorana fermions only have two degrees of freedom. The corresponding Majorana mass term $m_i \bar{\psi}_i^c \psi_i$ with ($i = L, R$) also connects a left-handed field to its CP conjugated counterpart (right-handed), but now violates lepton number by two units. This fundamental ambiguity raises the crucial question: Are neutrinos Dirac or Majorana particles? The answer has profound implications for the mechanism behind their tiny masses, the symmetry structure of the SM, and the origin of the baryon asymmetry in the universe.

To examine the neutrino mass terms, let us first add right-handed sterile neutrinos ν_R to the SM and write all the possible combinations:

$$-\mathcal{L}_{m_\nu} = m_D \bar{\nu}_R \nu_L + \frac{1}{2} m_L \bar{\nu}_L^c \nu_L + \frac{1}{2} m_R \bar{\nu}_R^c \nu_R + \text{H.c.} \quad (2.11)$$

Here, the first term corresponds to the Dirac mass term whereas the other two represent Majorana mass terms. Note that we do have the m_L term even in the absence of ν_R ; however, such a term is not gauge invariant and requires non-renormalizable operators (in particular, the Weinberg operator). Therefore, adding the right-handed field is necessary only to generate the Dirac mass term.

If neutrinos are purely Dirac particles, the Majorana mass term m_R must be forbidden by imposing an $U(1)$ lepton number symmetry. However, this means that lepton number becomes an imposed symmetry rather than an accidental one—a choice that is less appealing on theoretical grounds. Another challenge with Dirac neutrinos is explaining their tiny sub-eV masses. In the Standard Model, Dirac fermions acquire mass through the Higgs mechanism; neutrinos would do the same after EWSSB. Given the EW breaking scale of $\mathcal{O}(100 \text{ GeV})$, generating sub-eV neutrino masses requires extremely small Yukawa couplings, $Y_\nu \sim \mathcal{O}(10^{-12})$, which are about six orders of magnitude smaller than the electron's Yukawa coupling. Such fine-tuning is generally viewed as unattractive and suggests the existence of a deeper underlying mechanism.²

Now consider the Majorana case, where the Dirac mass term is forbidden. Majorana mass terms violate lepton number, but this is not problematic within an EFT framework, where accidental symmetries do not need to be strictly respected. Moreover, the right-handed neutrinos ν_R are sterile—they do not participate in weak interactions or the Higgs mechanism—and can thus have a wide range of masses. Interestingly, a

² One way to motivate Dirac neutrino masses and have suppressed Yukawas is to introduce extra dimensions, see Refs. [91, 92].

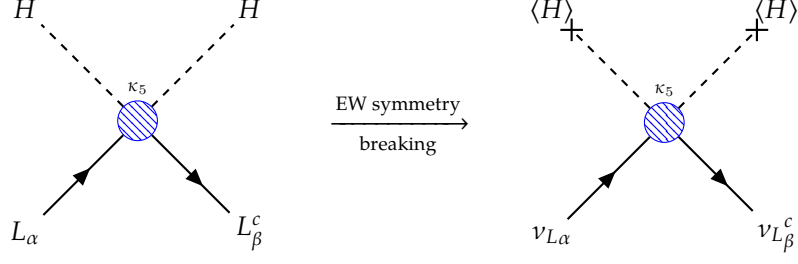


Figure 2.4: Majorana mass generation from the Weinberg operator after spontaneous symmetry breaking. Here, κ_5 is the effective coupling and α, β denote flavour indices.

very large mass of the right-handed neutrino can lead to very small active neutrinos via the so called **seesaw mechanism** that we review below. As has been already been pointed out above, the cleanest signature of Majorana neutrinos is $0\nu\beta\beta$, arising from the Majorana mass induced $\nu - \bar{\nu}$ mixing.

For these reasons, the Majorana scenario is often considered more *elegant*. Of course, one can argue about the notion of beauty and aesthetics as guiding principles in particle physics which have served us well in reaching our current formulation of the theory. Going forward, we will focus on Majorana neutrinos and see how the Majorana mass terms can be generated.

2.3 The Weinberg operator and the Seesaw mechanism

Let us recall the EFT approach discussed in Sec. 1.3.1 involving higher dimensional operators ($d \geq 5$). The Weinberg operator is the unique dimension-5 operator ($d = 5$) which can generate tiny masses for neutrinos after electroweak symmetry breaking. The operator is compatible with the SM gauge group and can be written as: $\mathcal{L}_5 \sim \kappa_5 L_L L_L H H$, where L and H are the SM lepton and Higgs doublets with hypercharge $-1/2$ and $1/2$ respectively. The coefficient $\kappa_5 \sim c_5/\Lambda_{NP}$ is the effective coupling suppressed by the high scale $\Lambda_{NP} \gg v_H$, where $\langle H \rangle = v_H = 246$ GeV is the SM Higgs VEV, and c_5 is a dimensionless coefficient. The operator in its full form is given by [55]

$$-\mathcal{L}_5 = \frac{1}{2} \kappa_5^{\alpha\beta} (\bar{L}_\alpha \tilde{H})(H^\dagger \tilde{L}_\beta) + \text{H.c.}, \quad (2.12)$$

where we include the flavour indices in the dimensionful coefficient κ_5 . After EWSSB, the Higgs acquires a VEV and generates a Majorana mass term for neutrinos

$$\mathcal{L}_5 \supset \frac{1}{2} m_\nu^{\alpha\beta} \bar{\nu}_L^c \nu_L + \text{H.c.} \implies m_\nu^{\alpha\beta} = c_5^{\alpha\beta} \frac{v_H^2}{\Lambda_{NP}}. \quad (2.13)$$

Assuming $c_5 \sim 1$ and plugging in $v_H = 246$ GeV, in order to get $m_\nu \sim \mathcal{O}(0.01$ eV), we get $\Lambda_{NP} \sim \mathcal{O}(10^{15}$ GeV), close to the GUT scale (a surprising coincidence!). The L and H fields in the Weinberg operator due to their $SU(2)$ properties can be contracted in three different ways to form a gauge singlet ($2 \otimes 2 = 3 \oplus 1$). However, only one of

these contractions is truly independent, as the others are related to it through Fierz transformations.

$$(\bar{L}_\alpha \tilde{H})(H^\dagger \tilde{L}_\beta) = \frac{1}{2}(\bar{L}_\alpha \vec{\sigma} \tilde{L}_\beta)(H^\dagger \vec{\sigma} \tilde{H}) = -(\bar{L}_\alpha \vec{\sigma} \tilde{H})(H^\dagger \vec{\sigma} \tilde{L}_\beta), \quad (2.14)$$

where $\vec{\sigma} = (\sigma_1, \sigma_2, \sigma_3)$. These correspond to the UV completions of the operator known as tree level seesaws: Type I - fermion singlets [93–96], Type II - scalar triplets [97–100]) and Type III - fermion triplets [101], respectively. These tree-level seesaw mechanisms thus offer a glimpse into high-scale physics through the observation of tiny neutrino masses. However, the scale suggested by the order-of-magnitude estimate above is far beyond the reach of current and foreseeable collider experiments, making direct probes extremely challenging. Nonetheless, additional suppression from the dimensionless coefficient c_5 can lower Λ_{NP} , providing motivation for *low-scale seesaw* models. We discuss these models in detail below.

While the Weinberg operator is unique in the SM, the situation changes in the presence of new scalar fields in BSM frameworks. In principle, one can construct new ‘Weinberg-like’ dimension-5 operators involving these additional scalars and the SM lepton doublets. Since these are dimension-5 operators, one can construct them with up to one or two new scalar multiplets, for example: $LLH\Phi_i, LL\Phi_i\Phi_j$ with $i, j = 1, 2$. Once these new scalars take a VEV, we will obtain light neutrino masses proportional to the new VEV. The VEVs of such new scalars are constrained to be $\lesssim \mathcal{O}(1 \text{ GeV})$ from electroweak precision measurements. Such small VEVs can be induced by the SM Higgs VEV, which effectively results in $d > 5$ operators for neutrino masses, thereby allowing for a further reduction of the associated new physics scale [102, 103]. This was the motivation behind the work carried out in Ref. [6] by the author, where we extensively discussed the new operators and their phenomenology.

2.4 The standard seesaw models

In the section above, we discussed how the different contractions of the fields involved in the Weinberg operator lead to three distinct UV completions at tree level—the standard seesaws: type I/II/III. Here, we provide a brief review of the most commonly employed tree-level seesaw models. The seesaw models beyond these usual seesaws obtained from UV completions of the new-Weinberg like operators were studied in Ref. [7].

2.4.1 Type I

In Type-I seesaw (SSI), the SM is extended by n heavy Right-Handed Neutrinos (RHNs) which are fermion singlets.³ The gauge-invariant Lagrangian for the type-I seesaw can

³ Usually $n = 3$ to generate three active neutrino masses, however, the minimal models contain $n = 2$ as we have currently measured only two mass splittings and the lightest neutrinos can be taken to be

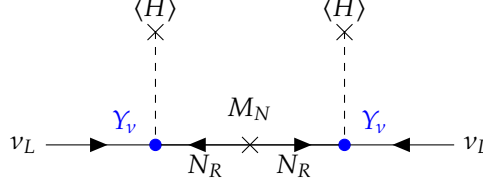


Figure 2.5: Schematic representation of the type-I see saw mechanism to generate tiny Majorana mass $m_\nu \approx v_H^2 Y_\nu M_N^{-1} Y_\nu^T$ for the SM neutrinos.

be written as $\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{N_R}$, where

$$\mathcal{L}_{N_R} = \bar{N}_R i \not{\partial} N_R - \left(\frac{1}{2} \bar{N}_R^c M_N N_R + Y_\nu \bar{L} \tilde{H} N_R + \text{H.c.} \right), \quad (2.15)$$

where Y_ν is a $3 \times n$ Yukawa matrix and M_N is a complex symmetric $n \times n$ matrix. We can choose a basis where M_N is real and diagonal without loss of generality, hence, we will work using that. A detailed derivation to obtain the effective Lagrangian after integrating out the heavy N_R fields is provided in Appendix A. After the EWSSB, we can write the neutrino mass Lagrangian as (also see Fig. 2.5)

$$-\mathcal{L}_{\text{mass}} = \bar{\nu}_L m_D N_R + \frac{1}{2} \bar{N}_R^c M_N N_R + \text{H.c.} \quad (2.16)$$

Here, we define $m_D \equiv Y_\nu v_H / \sqrt{2}$. Making use of $\bar{\nu}_L m_D N_R = \bar{N}_R^c m_D^T \nu_L^c$, the above equation can be rewritten as

$$-\mathcal{L}_{\text{mass}} = \frac{1}{2} \begin{pmatrix} \bar{\nu}_L & \bar{N}_R^c \end{pmatrix} \begin{pmatrix} 0 & m_D \\ m_D^T & M_N \end{pmatrix} \begin{pmatrix} \nu_L^c \\ N_R \end{pmatrix} + \text{H.c.} = \frac{1}{2} \bar{\chi}_R^c \mathcal{M}_\nu \chi_R + \text{H.c.} \quad (2.17)$$

where $\chi_R \equiv (\nu_L^c, N_R)^T$. For n right-handed neutrinos along with the three active ones, the mass matrix can be diagonalized by a $(3+n) \times (3+n)$ unitary matrix $U_0^T \mathcal{M}_\nu U_0 = \mathcal{M}_\nu^{\text{diag}}$, where $\mathcal{M}_\nu^{\text{diag}} = \text{Diag}[m_1, m_2, m_3, M_1, M_2, \dots, M_n]$ with mass eigenvalues m_i ($i = 1, 2, 3$) and M_j ($j = 1, 2, \dots, n$). The diagonalizing matrix U_0 is given by [104]

$$U_0 = W T = \begin{pmatrix} U_L & V \\ S & U_H \end{pmatrix} = \begin{pmatrix} (1 - \frac{1}{2}\eta)U_\nu & m_D^* (M_N^{-1})^* U_R \\ -M_N^{-1} m_D^T U_\nu & (1 - \frac{1}{2}\eta')U_R \end{pmatrix}, \quad (2.18)$$

where U_L, V, S and U_H are 3×3 , $3 \times n$, $n \times 3$ and $n \times n$ matrices, respectively, which need not be unitary. U_L is nothing but the neutrino mixing matrix that we discussed before, which has now acquired a non unitarity correction, $1 - \eta/2$. For $M_N \gg m_D$, the mass matrix can be first diagonalized by the matrix W , taking the $(3+n) \times (3+n)$

massless. From a theoretical perspective, the $n = 3$ is more appealing because of anomaly cancellations in a gauged $U(1)_{B-L}$ model.

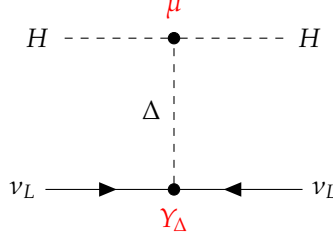


Figure 2.6: Feynman diagram for the type-II see saw mechanism to generate tiny Majorana mass $m_\nu = 2Y_\Delta v_\Delta$ for the SM neutrinos.

neutrino mass matrix to the block diagonal form

$$W^T \begin{pmatrix} 0 & m_D \\ m_D^T & M_N \end{pmatrix} W = \begin{pmatrix} M_{\text{light}} & 0 \\ 0 & M_{\text{heavy}} \end{pmatrix} \quad (2.19)$$

with

$$W \simeq \begin{pmatrix} 1 - \frac{1}{2} m_D^* (M_N^* M_N)^{-1} m_D^T & m_D^* (M_N^{-1})^* \\ -M_N^{-1} m_D^T & 1 - \frac{1}{2} (M_N)^{-1} m_D^T m_D^* (M_N^{-1})^* \end{pmatrix} \quad (2.20)$$

In the second step, $T = \text{Diag}[U_\nu, U_R]$ further diagonalizes M_{light} and M_{heavy} corresponding to the light and heavy neutrino sector, respectively. Comparing eq. (2.18) and eq. (2.20), values of the parameters characterizing non-unitarity, η and η' can be inferred. Hence, from eq. (2.19) and eq. (2.20), we get the well-known seesaw expressions⁴

$$M_{\text{light}} = m_\nu \simeq -m_D M_N^{-1} m_D^T, \quad M_{\text{heavy}} \simeq M_N. \quad (2.21)$$

The Type-I seesaw is by far the most popular model for generating neutrino masses—and for good reasons. It involves the addition of right-handed singlet fermions, which is a natural extension of the Standard Model, as all charged fermions already have their right-handed singlet counterparts. Furthermore, these singlets can be easily embedded into broader high-scale frameworks such as left-right symmetric models [100, 106] or GUTs (for example, $SO(10)$) [107].

2.4.2 Type II

When the L and H fields in the Weinberg operator are contracted in a way that their product transforms as an $SU(2)$ scalar triplet, the corresponding UV complete model is known as Type-II seesaw (SSII). In this model, a scalar triplet Δ with $Y = 1$ is added

⁴ Note that the Yukawa couplings Y_ν that determine m_D , need to be parameterized in order to fit the data from neutrino oscillations; there are several options for this parameterization, the most common being the Casas-Ibarra parameterization [105], which we use in the next chapter.

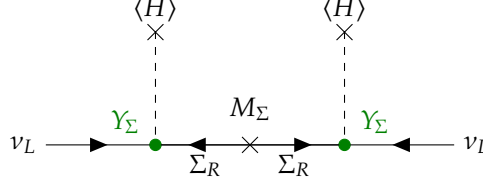


Figure 2.7: Feynman diagram for the type-III see saw mechanism to generate tiny Majorana mass $m_\nu \simeq v_H^2 Y_\Sigma M_\Sigma^{-1} Y_\Sigma^T$ for the SM neutrinos

to the SM, where

$$\Delta = \frac{1}{\sqrt{2}} \sigma_i \Delta_i = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}, \quad (2.22)$$

with $\Delta^{++} = (\Delta_1 - i\Delta_2)/\sqrt{2}$, $\Delta^+ = \Delta_3$ and $\Delta^0 = (\Delta_1 + i\Delta_2)/\sqrt{2}$. The gauge-invariant Lagrangian for the type-II seesaw relevant for neutrino masses is given by

$$\mathcal{L}_\Delta = -Y_\Delta \bar{L}^c \Delta L + \mu \tilde{H}^\dagger \Delta^\dagger H - M_\Delta^2 \text{Tr}[\Delta \Delta^\dagger] + \text{H.c.}, \quad (2.23)$$

where Y_Δ is a flavour symmetric matrix and μ is a dimensionful parameter. Note that lepton number is violated by two units due to the simultaneous presence of both the Yukawa as well as the μ term. After EWSSB, the neutral component of the scalar triplet also acquires a VEV (which can be included from the SM Higgs VEV, see the discussion on induced VEVs in Ref. [6])

$$\langle \Delta^0 \rangle \equiv v_\Delta \simeq \mu \frac{v_H^2}{2M_\Delta^2}. \quad (2.24)$$

We can then generate the neutrino masses from the Yukawa term as

$$m_\nu = 2Y_\Delta v_\Delta = Y_\Delta \mu \frac{v_H^2}{M_\Delta^2}. \quad (2.25)$$

The corresponding Feynman diagram is shown in Fig. 2.6. As Δ is a $SU(2)$ triplet, its VEV is constrained to be less than $\mathcal{O}(1 \text{ GeV})$ from the ρ parameter. Due to this, the scale of M_Δ may be smaller than the type-I case where $m_\nu \propto v_H^2$. Furthermore, the scale can be suppressed by the LN-violating μ parameter, whose small values are technically natural in the sense of 't Hooft [43], as setting $\mu \rightarrow 0$ enhances the symmetry of the model. For example, taking $Y_\Delta \sim 1$, $\mu \sim \mathcal{O}(\text{keV})$, we can get $m_\nu \sim 0.01 \text{ eV}$ for $M_\Delta \sim \mathcal{O}(100 \text{ TeV})$.

2.4.3 Type III

This scenario corresponds to the UV completion of the Weinberg operator when L and H are contracted to form an $SU(2)$ fermion triplet. In Type-III seesaw (SSIII), we extend the SM by adding n right-handed fermion triplets Σ_R with $Y = 0$, analogous to SSI. As

they are triplets, they can be written as

$$\Sigma_R = \frac{1}{2}\sigma_i\Sigma_i = \begin{pmatrix} \Sigma_R^0/\sqrt{2} & \Sigma_R^+ \\ \Sigma_R^- & -\Sigma_R^0/\sqrt{2} \end{pmatrix}, \quad (2.26)$$

where $\Sigma_R^\pm = (\Sigma_{R,1} \mp i\Sigma_{R,2})/\sqrt{2}$. The Lagrangian governing neutrino masses is given by

$$\mathcal{L}_\Sigma = \sqrt{2} Y_\Sigma \bar{\Sigma} \tilde{H}^\dagger L - \frac{1}{2} \text{Tr}[\bar{\Sigma}_R M_\Sigma \Sigma_R^c] + \text{H.c.}, \quad (2.27)$$

where Y_Σ is a complex $3 \times n$ matrix, and M_Σ is a $n \times n$ matrix, which can be taken as diagonal and real. After EWSSB, we can write the neutrino mass Lagrangian as

$$-\mathcal{L}_{\text{mass}} = \frac{1}{2} \begin{pmatrix} \bar{\nu}_L & \bar{\Sigma}_R^{0c} \end{pmatrix} \begin{pmatrix} 0 & m_D \\ m_D^T & M_\Sigma \end{pmatrix} \begin{pmatrix} \nu_L^c \\ \Sigma_R^0 \end{pmatrix}, \quad (2.28)$$

where $m_D = Y_\Sigma v_H/\sqrt{2}$. For $M_\Sigma \gg m_D$, the mass matrix can be diagonalized to get the light neutrino masses

$$m_\nu = -m_D M_\Sigma^{-1} m_D^T, \quad (2.29)$$

just like SSI. The associated Feynman diagram is shown in Fig. 2.7. The key difference, however, lies in the presence of charged components in the fermion triplet, which can mix with the charged leptons and give rise to interesting signatures, such as lepton flavour violation. Furthermore, the direct production of these triplets at colliders implies $M_\Sigma \geq O(100 \text{ GeV})$ [108, 109].

For a detailed discussion on novel seesaw models involving multiple $SU(2)$ fermionic multiplets, see the author's work in Ref. [7].

2.5 Low-scale models

As BSM models are scrutinised by their predictability and testability, it is crucial to develop scenarios that yield testable predictions. While the standard seesaws discussed above offer elegant explanations for the sub-eV scale of neutrino masses, they are difficult to probe at current collider experiments that probe TeV-scale physics. Moreover, the high scale associated with seesaw models may introduce a hierarchy problem [110, 111], as we discussed in Section 1.2.

Although nature is under no obligation to reveal new physics at the TeV scale, it remains a compelling target. Building low-scale models with $\Lambda_{\text{NP}} \sim O(\text{TeV})$ is therefore a worthwhile effort: either we discover new phenomena or we place meaningful constraints on the model's parameter space. In this section, we focus on TeV-scale extensions of the type-I seesaw and on radiative models, where neutrino masses are generated at the loop level.

2.5.1 Inverse and linear seesaw

These models are TeV-scale extensions of SSI (can be extended for SSIII), where in addition to the n right-handed neutrinos N_R carrying LN +1, m singlet fermions S with LN -1 are added. The key idea is to decouple the scale of LN breaking from the mass scale of the heavy neutrinos. The most general Lagrangian relevant for generating neutrino masses in this framework can be written as

$$-\mathcal{L}_\nu = Y_N \bar{L} \tilde{H} N_R + Y_S \bar{L} \tilde{H} S + \overline{N_R^c} M_R S + \frac{1}{2} \overline{N_R^c} M_N N_R + \frac{1}{2} \overline{S^c} \mu S + \text{H.c.}, \quad (2.30)$$

where Y_S, Y_N are the $3 \times m$ and $3 \times n$ Yukawa matrices, μ is the $m \times m$ Majorana mass matrix and M_R is a Dirac mass matrix. Here, we will fix $M_N = 0$.⁵ After EWSSB, this becomes

$$-\mathcal{L}_{\text{mass}} = \bar{\nu}_L m_D N_R + \bar{\nu}_L m_S S + \overline{N_R^c} M_R S + \frac{1}{2} \overline{S^c} \mu S + \text{H.c.}, \quad (2.31)$$

where $m_D = Y_N v_H / \sqrt{2}$ as in SSI, $m_S = Y_S v_H / \sqrt{2}$. Hence, we write the full mass matrix in the (ν_L^c, N_R, S) basis as

$$-\mathcal{L} = \frac{1}{2} \begin{pmatrix} \bar{\nu}_L & \overline{N_R^c} & \overline{S^c} \end{pmatrix} \begin{pmatrix} 0 & m_D & m_S \\ m_D^T & 0 & M_R \\ m_S^T & M_R^T & \mu \end{pmatrix} \begin{pmatrix} \nu_L^c \\ N_R \\ S \end{pmatrix} + \text{H.c.}, \quad (2.32)$$

The above mass matrix must be diagonalized to obtain the mass eigenstates. However, its structure allows for two special cases that simplify the diagonalization process.

Inverse seesaw: Consider the case where $m_S = 0$ in eq. (2.32), so we have

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & 0 & M_R \\ 0 & M_R^T & \mu \end{pmatrix}, \quad (2.33)$$

which can be block-diagonalized by a $(3 + n + m) \times (3 + n + m)$ unitary matrix to give the effective light neutrino mass matrix

$$m_\nu \simeq m_D (M_R^T)^{-1} \mu (M_R)^{-1} m_D^T. \quad (2.34)$$

Now, note that the Majorana mass term for S, μ , violates lepton number, hence, it is natural to take it to be quite small (a non-zero μ can be generated via radiative corrections). Therefore, the light neutrinos are not only doubly-suppressed by m_D/M_R ($m_D \ll M_R$) but also by the small μ parameter. In the limit of conserved lepton number: $\mu \rightarrow 0 \Rightarrow m_\nu \rightarrow 0$. This is called the Inverse Seesaw (ISS) mechanism [112, 113]. Note that the

⁵ M_N being a Majorana mass term is forbidden by the $U(1)_L$ symmetry. However, it can still take a small non-zero value like μ term. We take the former to be 0 because its contribution to the active neutrino masses will be sub-dominant compared to the latter.

minimal model to generate two light neutrino masses requires $m = n = 2$. Taking $O(1)$ Yukawas and $\mu \sim 0.1$ keV, one can get $m_\nu \sim 0.01$ eV for $M_R \sim \mathcal{O}(\text{TeV})$, which can be tested at the colliders, and other signals can come from sizeable m_D/M_R mixing.⁶

Linear seesaw: Now consider the case $\mu = 0$ and $m_S \neq 0$ in eq. (2.32), leading to the following neutral fermion mixing matrix

$$\begin{pmatrix} 0 & m_D & m_S \\ m_D^T & 0 & M_R \\ m_S^T & M_R^T & 0 \end{pmatrix}. \quad (2.35)$$

In the limit $M_R \gg m_D \gg m_S$, the above matrix can be diagonalized to give

$$m_\nu = m_D(m_S M_R^{-1})^T + (m_S M_R^{-1})m_D^T, \quad (2.36)$$

where the neutrino masses are now linearly proportional to m_D instead of $\sim m_D^2$ in SSI. Therefore, this model is called Linear Seesaw (LSS) [115,116]. In this case, the minimal model requires $m = n = 1$. The light neutrino mass scale can be obtained for a TeV-scale M_R for suppressed values of the LN-violating parameter m_S . A similar mass matrix occurs in our seesaw models [7], which are characterized by two distinct Yukawa couplings $y_{1,2}$ corresponding to the scalar multiplets $\Phi_{1,2}$. In the limit of hierarchical Yukawas, say $y_1 \gg y_2$, we can have TeV scale seesaw as well as a rich phenomenology from Φ_1 .

2.5.2 Radiative models

All the models discussed above generate neutrino masses at tree level. However, it is also possible to generate tiny Majorana masses entirely (or dominantly) at the loop level in a class of models known as radiative models. These models are appealing because the loop suppression factors, along with mass ratios of particles running in the loop, naturally explain the sub-eV neutrino masses. At the same time, they predict new particles that may be accessible at current or near-future experiments. These models often allow for sizeable interactions, leading to a richer phenomenology. Some well-known radiative neutrino mass models are the Zee-Babu model [117–119] and the scotogenic model [120,121]. In this section, we focus on the scotogenic model, which not only explains neutrino masses but also provides a viable dark matter candidate, thereby linking two key BSM phenomena. For a detailed review of radiative neutrino mass models, see Ref. [122].

⁶ A variant of the inverse seesaw termed ‘Double seesaw’ has the same mass matrix structure but with $\mu \gg m_D, M_R$ and $m_D \ll M_R^2/\mu$; however, this does not lead to a low-scale scenario [113,114].

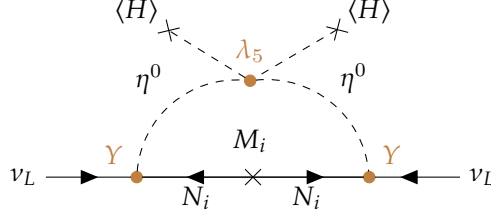


Figure 2.8: Neutrino mass generation in the scotogenic model at one loop, where η^0 corresponds to $\eta_{R,I}$.

The Scotogenic model

Named after the Greek word for darkness, *scotos*, the Scotogenic Model (ScM) offers an explanation for small neutrino masses through their interactions with dark matter. The SM is extended by three (or two in the minimal case) right-handed fermion singlets N_i along with an inert scalar doublet $\eta = (\eta^+, \eta^0)^T$. All these new fields are odd under a discrete $\mathbb{Z}_2$ symmetry, while the SM fields are even. The $\mathbb{Z}_2$ discrete symmetry is commonly used in dark matter model building, since it ensures that the lightest $\mathbb{Z}_2$ -odd particle is stable and a viable dark candidate. The Yukawa Lagrangian for the model is given by

$$-\mathcal{L}_{\text{ScM}} = Y_{\alpha i} \bar{L}_\alpha \tilde{\eta} N_i + \frac{1}{2} \overline{N_i^c} M_i N_i + \text{H.c.}, \quad (2.37)$$

and the scalar potential is given by

$$\begin{aligned} V(\eta, H) = & m_1^2 H^\dagger H + m_2^2 \eta^\dagger \eta + \frac{1}{2} \lambda_1 (H^\dagger H)^2 + \frac{1}{2} \lambda_2 (\eta^\dagger \eta)^2 + \lambda_3 (H^\dagger H) (\eta^\dagger \eta) \\ & + \lambda_4 (H^\dagger \eta) (\eta^\dagger H) + \left[\frac{1}{2} \lambda_5 (H^\dagger \eta)^2 + \text{H.c.} \right], \end{aligned} \quad (2.38)$$

where Y_ν is a 3×3 matrix, M_N is the heavy Majorana mass matrix, and all the couplings in the scalar potential can be chosen to be real. Note that because of the $\mathbb{Z}_2$ -odd charge, the lepton doublets couple to N only when η is present (instead of H in SSI). For $m_1^2 < 0$ and $m_2^2 > 0$, we will have $\langle H \rangle = v_H / \sqrt{2}$ and $\langle \eta^0 \rangle = 0$. This implies that the EW symmetry is spontaneously broken but the $\mathbb{Z}_2$ symmetry remains intact, leading to $m_\nu = 0$ at tree-level. However, the λ_5 term in the scalar potential can induce small Majorana masses for neutrinos at the loop level, via the diagram shown in Fig 2.8. The mass spectrum of the scalars after EW SSB in the model is given by

$$\begin{aligned} m_h^2 &= \lambda_1 v_H^2, & m_{\eta^\pm}^2 &= m_2^2 + \lambda_3 \frac{v_H^2}{2}, \\ m_{\eta_R}^2 &= m_2^2 + (\lambda_3 + \lambda_4 + \lambda_5) \frac{v_H^2}{2}, & m_{\eta_I}^2 &= m_2^2 + (\lambda_3 + \lambda_4 - \lambda_5) \frac{v_H^2}{2}, \end{aligned} \quad (2.39)$$

where h is the SM Higgs boson. Note the mass splitting between the two η components, $m_{\eta_R}^2 - m_{\eta_I}^2 \propto \lambda_5 v_H^2$. The light neutrino mass matrix gets contributions from this exchange

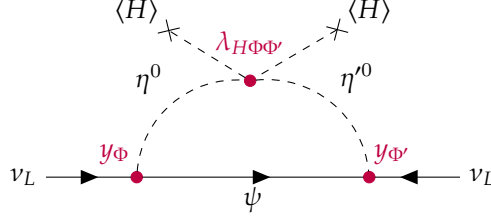


Figure 2.9: Neutrino mass generation in the generalized scotogenic model at one loop.

of η_R and η_L , so we have

$$(\mathcal{M}_\nu)_{\alpha\beta} = \sum_i Y_{\alpha i} Y_{\beta i} \frac{M_i}{32\pi^2} \left[\frac{m_{\eta_R}^2}{m_{\eta_R}^2 - M_i^2} \ln \left(\frac{m_{\eta_R}^2}{M_i^2} \right) - \frac{m_{\eta_L}^2}{m_{\eta_L}^2 - M_i^2} \ln \left(\frac{m_{\eta_L}^2}{M_i^2} \right) \right]. \quad (2.40)$$

We can make use of $m_{\eta_R}^2 - m_{\eta_L}^2 = \lambda_5 v^2 \ll m_0^2 \simeq m_2^2 + (\lambda_3 + \lambda_4) v_H^2 / 2$ to write

$$(\mathcal{M}_\nu)_{\alpha\beta} = \frac{\lambda_5 v_H^2}{32\pi^2} \sum_i Y_{\alpha i} Y_{\beta i} \frac{M_i^2}{m_0^2 - M_i^2} \left[1 - \frac{M_i}{m_0^2 - M_i^2} \ln \left(\frac{m_0^2}{M_i^2} \right) \right]. \quad (2.41)$$

Here, we see that the neutrino masses can be suppressed if $\lambda_5 \ll 1$. This is once again natural, because the model contains a global $U(1)_L$ symmetry, which is enhanced by setting $\lambda_5 \rightarrow 0$. Assuming hierarchies in the mass scales M_i, m_0 , the above expression can be further simplified. For instance, a TeV-scale M_i can produce $m_\nu \sim 0.01$ eV for $\lambda_5 Y_\nu^2 \sim \mathcal{O}(10^{-9})$. Depending on the masses of η and N_i , the lightest of them can be a DM candidate (taking $\lambda_4 \pm \lambda_5 > 0$, see eq. (2.39)). Given their fermionic and scalar nature, the dark matter phenomenology differs depending on the candidate. For instance, in the case of fermionic dark matter, there are constraints from LFV observables. In contrast, for scalar DM, we have constraints on λ_5 from the DM relic abundance and direct detection experiments, which we discuss in Chapter 4.

Generalized scotogenic model

An interesting variant of the scotogenic model introduces Dirac mediators along with new scalar doublets. All these new fields are charged under a new global $U(1)$ symmetry. This is called the Generalised Scotogenic Model (GScM), see Refs. [123–126].

Let us consider a model based on a global $U(1)_{\text{DM}}$ symmetry, where one Dirac fermion ψ and two scalar doublets Φ, Φ' (with hypercharges 1/2 and -1/2, respectively) are added to the SM. All the new fields are charged +1 under an unbroken $U(1)_{\text{DM}}$ symmetry, so $\langle \Phi \rangle = \langle \Phi' \rangle = 0$. The Lagrangian is given by [125]

$$\mathcal{L}_{\text{GScM}} = i\bar{\psi}\not{\partial}\psi - m_\psi\bar{\psi}\psi - \left(y_\Phi^\alpha \bar{\psi} \tilde{\Phi}^\dagger L_\alpha + (y_{\Phi'}^\alpha)^* \bar{\psi} \tilde{\Phi}'^\dagger \tilde{L}_\alpha + \text{H.c.} \right), \quad (2.42)$$

where the Yukawa couplings y 's are three-component vectors. The scalar potential is

a bit more complex than the ScM due to the presence of two new scalars

$$\begin{aligned}
V_{H,\Phi,\Phi'} = & -m_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 + m_\Phi^2 \Phi^\dagger \Phi + \lambda_\Phi (\Phi^\dagger \Phi)^2 + m_{\Phi'}^2 \Phi'^\dagger \Phi' + \lambda_{\Phi'} (\Phi'^\dagger \Phi')^2 \\
& + \lambda_{H\Phi} (H^\dagger H) (\Phi^\dagger \Phi) + \lambda_{H\Phi'} (H^\dagger H) (\Phi'^\dagger \Phi') + \lambda_{\Phi\Phi'} (\Phi^\dagger \Phi) (\Phi'^\dagger \Phi') \\
& + \lambda_{H\Phi,2} (H^\dagger \Phi) (\Phi^\dagger H) + \lambda_{H\Phi',2} (H^\dagger \tilde{\Phi}') (\tilde{\Phi}'^\dagger H) + \lambda_{\Phi\Phi',2} (\Phi^\dagger \tilde{\Phi}') (\tilde{\Phi}'^\dagger \Phi) \\
& + \lambda_{H\Phi\Phi'} [(H^\dagger \tilde{\Phi}') (H^\dagger \Phi) + \text{H.c.}] .
\end{aligned} \tag{2.43}$$

After EWSSB, neutrino masses are generated at the one-loop level by the diagram shown in Fig. 2.9. The scalar fields mix among themselves, leading to the physical states: SM Higgs boson h , two neutral scalars η^0, η'^0 and two charged scalars η^+, η'^+ . Their masses given by

$$m_{\eta^+} = \sqrt{m_\Phi^2 + \frac{1}{2} \lambda_{H\Phi} v_H^2}, \quad m_{\eta'^0}^2 \simeq m_{\eta^+}^2 + \frac{1}{2} \lambda_{H\Phi,2} v_H^2, \tag{2.44}$$

and $m_{\eta'^+, \eta'^0}$ can be obtained by replacing the terms with their primed counterparts. The mixing angle θ is given by $\tan 2\theta = 2c/(a' - a)$ where

$$c = -\frac{1}{2} \lambda_{H\Phi\Phi'} v_h^2, \quad a = m_\Phi^2 + \frac{1}{2} v_h^2 (\lambda_{H\Phi} + \lambda_{H\Phi,2}), \tag{2.45}$$

and a' is obtained by substituting the masses and coupling with their primed counterparts. The light neutrino mass matrix can thus be written as

$$(\mathcal{M}_\nu)_{\alpha\beta} = \frac{\sin 2\theta m_\psi}{32\pi^2} (y_\Phi^\alpha y_{\Phi'}^\beta + y_{\Phi'}^\alpha y_\Phi^\beta) L(m_{\eta^0}, m_{\eta'^0}, m_\psi). \tag{2.46}$$

While this model can also feature either fermion or scalar dark matter (scalar DM is generally disfavored compared to fermion DM), it exhibits phenomenologically distinct signatures from the scotogenic model due to the presence of two different Yukawa couplings. For instance, the Yukawa combinations that contribute to the light neutrino mass differ from those involved in charged lepton flavour violation processes. Moreover, in the minimal version of the model, only one new fermion is required to generate masses for two light neutrinos.

We make use of the ISS, ScM and GScM models in our work in Ref. [7] to build low-scale variants of the novel seesaw models.

2.5.3 Collider searches

Here, we briefly discuss the implications of these low-scale seesaw models at colliders as they serve as a big motivation behind them, see Ref. [127] for a review. Let us start with the right-handed Majorana neutrinos N_R ,⁷ which can be produced at the colliders

⁷ Right handed neutrinos are a part of a broader search program corresponding to heavy neutral leptons (HNLs), which may or may not be linked to neutrino mass generation, see Ref. [128] for a recent status report.

and give signals like [129–132]

$$\begin{aligned}
 pp &\rightarrow W^{*\pm} \rightarrow l^\pm N_i \rightarrow l^\pm l^\pm jj, \\
 pp &\rightarrow W^{*\pm} W^{*\pm} \rightarrow l^\pm l^\pm jj, \\
 pp &\rightarrow W^{*\pm} \gamma^* \rightarrow l^\pm N_i jj
 \end{aligned} \tag{2.47}$$

where j corresponds to jets, and there is no missing transverse energy. These LN-violating processes are smoking gun signals of Majorana neutrinos (in a way analogous to $0\nu\beta\beta$). Such processes allow us to constrain the $|V_{lN}|^2 - M_N$ parameter space ($l = e, \mu, \tau$), where the former characterizes the mixing. On the other hand, for Dirac-type heavy neutrinos ($\Delta L = 0$), we have: $pp \rightarrow W^{*\pm} \rightarrow l^\pm N_i \rightarrow l^\pm l^\pm l^\pm + \cancel{E}_T$ [133, 134], the so-called golden trilepton channel with the missing transverse energy, $\cancel{E}_T$. This channel is relevant also in inverse seesaw models with pseudo-Dirac neutrinos.

Next, we discuss the $SU(2)_L$ triplet scalars and fermions of SSII and SSIII, respectively. Being triplets, they can be directly produced at the LHC via their gauge interactions, unlike the singlet fermions. For Type-II seesaw, these processes are [133, 135]

$$\begin{aligned}
 pp &\rightarrow Z^*/\gamma^* \rightarrow \Delta^\pm \Delta^\mp, \Delta^{\pm\pm} \Delta^{\mp\mp}, \\
 pp &\rightarrow W^{*\pm} \rightarrow \Delta^{\pm\pm} \Delta^\mp, \Delta^{\pm\pm} W^\mp, \\
 pp &\rightarrow W^{*\pm} W^{*\pm} \Delta^\pm \Delta^\pm.
 \end{aligned} \tag{2.48}$$

There is a unique signature associated with SSII, where the doubly-charged scalar $\Delta^{\pm\pm}$ can decay to a pair of same-sign dileptons or gauge bosons: $\Delta^{\pm\pm} \rightarrow l^\pm l^\pm, W^\pm W^\pm$. The branching ratios to the leptonic channel or the bosonic channel depend on the VEV v_Δ and the neutrino mass matrix flavour structure. For $v_\Delta \lesssim \mathcal{O}(100 \text{ keV})$, the decays to dileptons dominate, which is a clean signature (free from the SM background) of LNV. Such signals can also be exploited in the case of new scalar multiplets, allowing us to place bounds on them from LHC searches, as we show in Ref. [6]. Currently, the lower limit on doubly charged scalars is 1080 GeV [136].

Finally, for the case of fermion triplets in SSIII, we have [133, 137]

$$\begin{aligned}
 pp &\rightarrow Z^*/\gamma^* \rightarrow \Sigma^\pm \Sigma^\mp, \\
 pp &\rightarrow Z^* \rightarrow \Sigma^\pm l^\mp, \Delta^{\pm\pm} W^\mp, \\
 pp &\rightarrow W^{*\pm} \Sigma^\pm \Sigma^0, l^\pm \Sigma^0.
 \end{aligned} \tag{2.49}$$

Note that the neutral component Σ^0 is similar to N_R , and has similar decay modes $\Sigma^0 \rightarrow l^\pm W^\mp, Z\nu_L(\bar{\nu}_L), H\nu_L(\bar{\nu}_L)$. The current lower bounds on M_Σ are 790 GeV [108] and 880 GeV [109] from ATLAS and CMS, respectively.

3

THE BARYON ASYMMETRY OF THE UNIVERSE

Almost everything we see and interact with is made of matter, while antimatter remains elusive and is more commonly associated with science fiction—famously featured in Dan Brown’s *Angels and Demons*, where canisters of antimatter are stolen from CERN to create a fictional antimatter bomb. When such a canister is opened, it would annihilate upon contact with matter, releasing an immense amount of energy. In reality, we know antimatter is extremely rare because we do not observe matter-antimatter annihilations at any significant scale in our daily lives.

That said, small amounts of antimatter constantly reach Earth as cosmic rays or are produced in the decays of energetic particles originating from outer space. Interestingly, bananas emit trace amounts of antimatter due to the presence of the isotope potassium-40. Noticeable quantities are also produced in medical imaging procedures and in particle accelerator experiments.

This naturally leads to the question: what if antimatter were not so rare and existed in equal amounts to matter? The short answer is that life as we know it would not exist. Our very existence depends on the excess of matter over antimatter. Yet, the origin of this excess remains an open question in particle physics. Understanding the observed matter-antimatter asymmetry—or more precisely, the Baryon Asymmetry of the Universe (BAU), as introduced in Chapter 1—requires a viable theoretical explanation.

We will see that the SM has all the ingredients to dynamically generate an asymmetry, but not in enough amounts to match the observations. In this chapter, we review the observational evidences that have established a baryon asymmetric universe, and then explore mechanisms that could generate it dynamically. A key focus will be on thermal leptogenesis, which links the BAU to the neutrino sector.

3.1 Observations and evidences

We believe that shortly after the Big Bang, the universe contained equal amounts of baryons and antibaryons—i.e., it began in a symmetric state. If the universe continued to be in a symmetric state, the matter-antimatter annihilations would render it empty. One can argue, what if it is symmetric, and there are regions in the visible universe composed entirely of antimatter, just like our observable region is made entirely of matter. We know that this is not the case because we would have detected intense gamma-ray emissions from the boundaries where matter and antimatter meet and annihilate.

Alternatively, one might argue that the universe started in an asymmetric state. However, this scenario would require a highly fine-tuned initial condition to match current observations. Moreover, any such primordial asymmetry would likely have been diluted by cosmic inflation, further weakening this argument.

These considerations strongly support the idea that the observed baryon asymmetry was dynamically generated from an initially symmetric universe. Mechanisms that can produce such an asymmetry fall under the umbrella of baryogenesis. The baryon asymmetry is quantitatively expressed as [29]:

$$\eta \equiv \left. \frac{n_B - n_{\bar{B}}}{n_\gamma} \right|_0 = (6.04 \pm 0.12) \times 10^{-10} \quad (3.1)$$

$$Y_{\Delta B} \equiv \left. \frac{n_B - n_{\bar{B}}}{s} \right|_0 = (8.58 \pm 0.17) \times 10^{-11}. \quad (3.2)$$

Here, the first definition is in terms of the baryon-to-photon ratio and the second one is in terms of yield, $Y \equiv n/s$ (number density upon entropy) to include the effect of expansion of the universe, with $\Delta B = B - \bar{B}$. The number densities of baryons, antibaryons and photons are denoted by $n_B, n_{\bar{B}}$ and n_γ , respectively. The subscript 0 denotes the present time. It is easy to see that the two definitions are related as

$$Y_{\Delta B} = \frac{n_\gamma}{s} \eta. \quad (3.3)$$

Now, we discuss the astrophysical and cosmological evidences against the presence of a large amount of antimatter in the universe, and how η is evaluated.

3.1.1 Antimatter searches

There are two ways to search for antimatter: direct and indirect detection. Direct detection involves searching for antimatter or antinuclei in the cosmic rays, it has been found that the upper limit on the flux of antiprotons to protons is $\Phi_{\bar{p}}/\Phi_p < 2 \times 10^{-4}$ [138]. The smallness of this ratio is attributed to the antiprotons being only the byproducts of the reaction initiated by primary cosmic ray proton collisions, such as $p + p \rightarrow \bar{p} + 3p$. Similarly, there are upper bounds on anti-Helium to Helium flux, $\Phi_{\bar{\text{He}}^4}/\Phi_{\text{He}^4} < \mathcal{O}(10^{-6})$

[139].¹ Both these limits indicate that antimatter is rare. On the other hand, the same conclusion can be drawn from indirect search experiments, which aim to detect the products of matter-antimatter annihilations. The observed spectrum of cosmic diffuse γ -rays imposes an upper limit on such annihilations and tells us that the size of the matter domain is equal to the scale of our visible universe [141].

3.1.2 Big Bang Nucleosynthesis

In the early universe, when the universe was still very hot and dense, i.e., at temperatures around $T \sim 0.1 - 1$ MeV (just within a few seconds after the Big Bang), nuclear reactions were fast enough for protons and neutrons to fuse into light nuclei such as deuterium (D) and helium-4 (He^4). This is known as Big Bang Nucleosynthesis (BBN), see Refs. [142, 143] for a review.

At $T \gg 1$ MeV, the neutron-proton conversions ($n + \nu_e \rightarrow p + e^-$) mediated by the weak interactions were in equilibrium ($\Gamma > H$, where Γ is the interaction rate and H is the Hubble expansion of the universe). These conversions freeze-out when $\Gamma < H$ at $T \sim 0.8$ MeV, and the neutrinos also decouple from the thermal plasma slightly before $T \sim 2$ MeV.

After freeze-out, most free neutrons quickly bound into He^4 nuclei via intermediate formation of deuterium. The synthesis of He^4 is governed by two key processes: $\text{D}(p, \gamma)\text{He}^3$ and $\text{He}^3(\text{D}, p)\text{He}^4$, both of these depend on the baryon to photon ratio η . Specifically, the number densities of the incoming species scale as η and η^2 , respectively.

Therefore, BBN allows us to calculate the primordial abundance of light elements such as D, He^3 , He^4 and Li^7 from a given value of η , and vice versa: measuring the abundances precisely allows us to constrain η .² A large value of η would suggest a smaller freeze out abundance of D and He^3 . Further, such larger values of η would imply that the relative photon number is smaller, leading to an earlier production of He^4 and subsequently higher abundance. The dependence of the light element abundances on η are shown in Fig. 3.1. Several observations lead to the following range of η (95% CL), consistent with primordial abundances [143]:

$$5.8 \times 10^{-10} \leq \eta \leq 6.5 \times 10^{-10}, \quad 0.021 \leq \Omega_B h^2 \leq 0.024. \quad (3.4)$$

Alternatively, since $\Omega_B h^2 = 3.66 \times 10^7 \eta$, and Planck has measured it to be 0.0024(2) [77], one can determine η and subsequently predict the light element abundances, as

¹ Recent results from AMS-02 have observed some anti-Helium 4 events, though the detection is yet to be confirmed. If confirmed, this would be an exciting signal present in several BSM models, specially DM annihilations, see Ref. [140].

² Note that the dependence of Li^7 is complicated, as the two key production processes have opposite dependencies on η . Moreover, BBN predicts a Li^7 abundance that is about three times larger than what is observed in old, metal-poor stars. This is known as the ‘Lithium problem/discrepancy’, see Ref. [144].

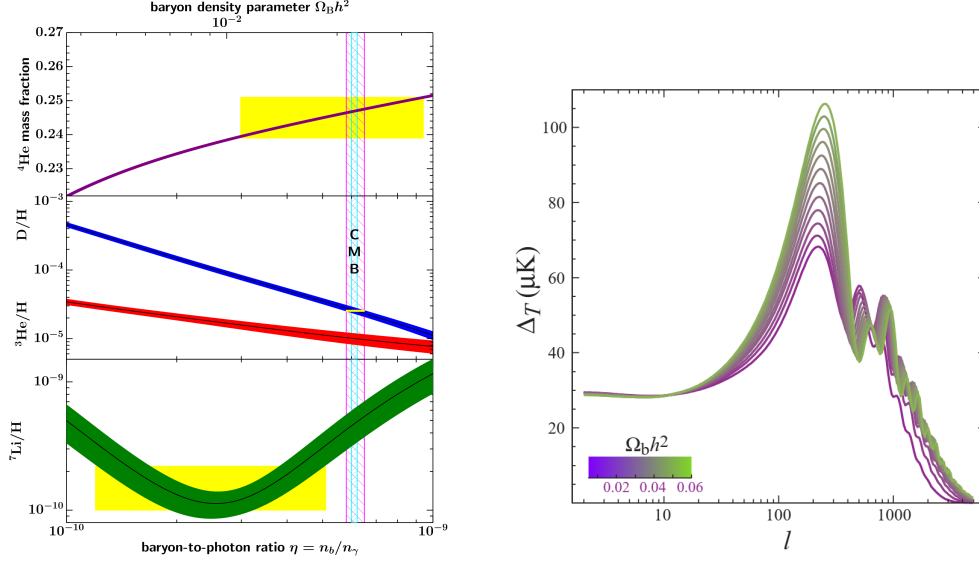


Figure 3.1: The predicted abundances of He^4 , D , He^3 and Li^7 from BBN. Yellow boxes indicate the observed values. The narrow vertical band shows the baryon density inferred from the CMB. Taken from [29]. Right: Sensitivity of the CMB acoustic spectrum to the physical baryon density $\Omega_b h^2$, varied around a fiducial model with $\Omega_b h^2 = 0.02$. The horizontal axis l corresponds to the multipole moment. Taken from [145].

described above.

3.1.3 Cosmic Microwave Background Radiation

About 300,000 years after the Big Bang, when the temperature had dropped to around $O(0.1 \text{ eV})$, electrons and protons combined to form neutral hydrogen—a period known as recombination. The Cosmic Microwave Background (CMB) formed shortly after recombination, as the thermal bath of the universe became transparent to the relic photons, which have been propagating freely ever since. The value of the baryon asymmetry can be inferred from the anisotropies in the CMB, which encode information about primordial density fluctuations.

The primordial density fluctuations lead to the formation of gravitational potential wells, with the baryon falling into it getting compressed and heated. This hot baryon fluid expanded in all directions, radiating photons until it eventually cooled down.³ As the radiation pressure decreased, gravitational forces regained dominance, initiating another phase of compression. This cyclical interplay between gravity and radiation pressure led to Baryon Acoustic Oscillations (BAO) in the baryon–photon fluid. These oscillations became imprinted in the CMB as a pattern of acoustic peaks and troughs—the defining features of the temperature anisotropy spectrum.

Measurements of these anisotropies, particularly by WMAP, provide an estimate of the

³ Prior to recombination, the thermal plasma is approximated as a photon-baryon fluid and its dynamics is governed by fluid mechanics and plasma physics.

cosmic baryon-to-photon number density ratio as [77]:

$$5.92 \times 10^{-10} \leq \eta \leq 6.16 \times 10^{-10}, \quad 0.02214 \leq \Omega_B h^2 \leq 0.02252 \quad (3.5)$$

The sensitivity of the acoustic spectrum to the physical baryon density $\Omega_B h^2$ is shown in Fig. 3.1. Notice the remarkable coincidence between the values of η in eqs. (3.4) and (3.5), from two different periods of the early universe at $t \gtrsim 1$ s and $t \sim 3.8 \times 10^5$, respectively. This highlights the success of Big Bang cosmology.

3.2 Baryogenesis

Having established the evidences for a baryon asymmetry of $O(10^{-10})$, as well as for its dynamical origin, we will briefly review some well-known baryogenesis mechanism, see Refs. [146, 147] for reviews on the topic. However, before that, let us discuss the famous Sakharov conditions for the dynamical generation of such an asymmetry.

3.2.1 Sakharov's conditions

In 1967, Andrei Sakharov wrote his seminal article on the baryon asymmetry of the universe and proposed three necessary conditions for Baryogenesis [148]. No text on baryogenesis is complete without mentioning these three conditions:

1. **Baryon number violation:** In order to evolve from a symmetric state ($\Delta B = 0$) to an antisymmetric state ($\Delta B \neq 0$), the first and foremost condition is presence of interactions which do not conserve baryon number.
2. **C and CP violation:** If C (charge conjugation) and CP (charge conjugation times parity) are not violated in addition to baryon number, no net asymmetry can be generated. This is because, in the absence of such violations, baryons and antibaryons would be produced at the same rate, unless a preference exists for specific final states, such as decays favouring baryons over antibaryons.
3. **Departure from thermal equilibrium:** In thermal equilibrium, asymmetries in non-conserved quantum numbers cannot be generated due to the principle of detailed balance. The chemical potential (μ) of particles with such non-conserved quantum numbers is 0. Moreover, CPT invariance ensures that particles and antiparticles have equal masses, leading to equal phase space densities for baryons and antibaryons. As a result, forward and backward reactions occur at the same rate, preventing any net asymmetry. A process is said to depart from equilibrium when its interaction rate falls below the Hubble expansion rate of the universe.

Interestingly, all the above-mentioned conditions are present in the SM, yet it still fails to account for the observed baryon asymmetry. First, in the SM, both baryon number B and lepton number L are conserved at the classical level, although they are accidental symmetries—unlike electric charge conservation ($U(1)_{\text{QED}}$). However, at the non-perturbative level, there exist interactions (sphalerons) that, after accounting for the

axial anomaly, violate $B + L$ while preserving $B - L$ [149]. These baryon number violating processes are present in the SM and are sufficiently fast in the early universe—we will review them in detail later.

Second, we know that the weak interactions violate both C and CP due to the presence of a complex phase in the CKM matrix. However, this CP violation is found to be of the order of 10^{-20} , which is 10 orders of magnitude too small to explain the observed baryon asymmetry [150, 151].

Finally, while a departure from thermal equilibrium could occur during the electroweak phase transition, this transition in the SM is a smooth crossover. Successful baryogenesis requires a strong first-order phase transition, which in the SM would necessitate a lighter Higgs boson, now ruled out experimentally [152].

3.2.2 Candidates

As we see that the SM alone cannot incorporate the three Sakharov's conditions: we require new physics beyond the SM to incorporate them all. Below, we discuss some famous extensions and modifications to the SM:

Electroweak baryogenesis: An appealing scenario that remains closely tied to the SM is electroweak baryogenesis (EW baryogenesis) [152–154], where the asymmetry is generated during the Electroweak Phase Transition (EWPT). However, for the EWPT to be strongly first-order (FOPT), any viable model of EW baryogenesis must involve modifications to the SM scalar potential that alter the nature of the transition. Additionally, such models require new sources of CP violation, since the one present in the SM is far too small. These requirements can be satisfied, for example, in Two Higgs Doublet Model (2HDM) or in the Minimal Supersymmetric Standard Model (MSSM) [155–157].

GUT baryogenesis: We saw that GUTs are motivated by the unification of the strong and electroweak forces. In these theories (for instance, $SU(5)$ or $SO(10)$), quarks and leptons are embedded in the same multiplet, and a new gauge boson X mediates interactions between them simultaneously [107]. If the decays of X violate CP and occur out of equilibrium, they can generate a net baryon asymmetry [158]. Furthermore, the non-observation of proton decay places a lower bound on the mass of the decaying boson. However, the simplest GUT models conserve $B - L$, and the rapid $B + L$ violating sphaleron processes in the early universe would wash out any preexisting asymmetry unless $B - L$ is also broken [159].

Affleck-Dine mechanism: In supersymmetric models, a non-zero baryon asymmetry can be generated through the coherent oscillations of a scalar field, denoted ϕ , which carries a baryon or lepton number—this is known as the Affleck-Dine mechanism [160]. Initially, when $H \gg M_\phi$, the scalar field is frozen at a large value $\phi = \phi_0 \gg 0$, and no baryon asymmetry is present. As the universe expands and $H < M_\phi$, the field

Table 3.1: Various baryogenesis mechanisms and how they meet the Sakharov conditions.

Model	B violation	CP violation	Out of equilibrium
EW baryogenesis	EW sphalerons	Bubble wall dynamics	FOPT
GUT baryogenesis	Heavy gauge boson X	Complex couplings	X decays
Affleck-Dine	Scalar ϕ	Scalar potential	ϕ oscillations
Leptogenesis	Majorana neutrinos N_i , EW sphalerons	Complex Yukawas	N_i decays

begins to oscillate, generating a non-zero baryon number. These scalar fields are the supersymmetric partners of quarks and leptons, with baryon number and CP violation arising from small parameters in the scalar potential. Eventually, the scalar field decays into SM particles, thereby transferring its baryon or lepton number to the SM thermal bath. Being a high-scale mechanism, it is free from electroweak constraints.

Leptogenesis: As the name suggests, this mechanism involves generating an initial lepton asymmetry which is then processed into a baryon one via the same sphalerons that have been mentioned above. First proposed by M. Fukugita and T. Yanagida in 1986 [161], leptogenesis models have become one of the most popular models for the BAU due to their strong connection with neutrino physics, particularly in seesaw mechanisms. We will discuss these connections in detail below; for a comprehensive review, see Ref. [162].

3.3 Leptogenesis

Let us focus on the type-I seesaw of Sec. 2.4.1, with the Lagrangian

$$\mathcal{L}_{N_i} = \bar{N}_i i \not{\partial} N_i - \left(\frac{1}{2} \bar{N}_i^c M_i N_i + y_{\alpha i} \bar{L}_\alpha \tilde{H} N_i + \text{H.c.} \right). \quad (3.6)$$

Looking at the Lagrangian above, we can see how the Sakharov conditions can be met:

1. **Baryon number violation:** Since the singlet neutrinos (RHNs) satisfy the Majorana condition $N_i = N_i^c$, their mere presence will lead to L -violation due to the mass and Yukawa terms present in the Lagrangian. The electroweak sphalerons relate this LNV with BNV.
2. **CP Violation:** The new Yukawa couplings (y) in the neutrino sector can generally be complex, providing a source of CP violation.
3. **Departure from thermal equilibrium:** Since the interactions of N with the SM leptons occur only through the Yukawa term, and given that the heavy neutrinos have a large mass scale to generate sub-eV neutrino masses, these RHNs can decay out of equilibrium. This happens when their decay rate $\Gamma(N \rightarrow lH)$ becomes smaller than the Hubble expansion.

Therefore, we see that the neutrino mass generation mechanism has the right ingredients for generating an asymmetry in lepton number.

Leptogenesis is governed by the values of the Yukawa couplings, which determine the rates of N_i decays and inverse decays. In the early universe, at temperatures $T \gg M_i$, a thermal population of N_i can be produced through inverse decays such as $LH \rightarrow N_i$ and $\overline{LH} \rightarrow N_i$. These interactions can be sufficiently rapid to maintain N_i in thermal equilibrium. However, at later times, when $T < M_i$, these processes fall out of equilibrium. The RHNs then decay completely, and their number density becomes Boltzmann suppressed.

A CP asymmetry, $\varepsilon \propto \Gamma(N \rightarrow LH) - \Gamma(N \rightarrow \overline{LH})$, arises in the decays due to the complex Yukawa couplings, leading to a net lepton asymmetry. For a hierarchical heavy neutrino spectrum ($M_1 \ll M_2 \ll M_3$), only the asymmetry generated from the lightest neutrino, say N_1 , is significant. This is because the asymmetries generated in the decays of N_2 and N_3 at $T \gg M_1$ can be washed out by N_1 mediated interactions, such as $LH \leftrightarrow N_1 \leftrightarrow \overline{LH}$. This scenario is known as thermal leptogenesis. Once the lepton asymmetry is generated, it is partially converted into a baryon asymmetry via electroweak sphaleron processes. The baryon asymmetry can then be expressed as

$$Y_{\Delta B} = C_{\text{sph}} Y_{\Delta L} = C_{\text{sph}} \varepsilon \delta Y_{N_1}^{\text{eq}}(T \gg M_1), \quad (3.7)$$

where C_{sph} is the sphaleron conversion factor, $Y_{N_1}^{\text{eq}} \equiv n^{\text{eq}}/s = 135\zeta(3)/(4\pi^4 g_*)$ is the equilibrium N_1 yield (n is the number density defined later in eq. (3.46)) with g_* , the total number of degrees of freedom of the thermal plasma. $\delta \leq 1$ is the efficiency parameter which characterizes the washout of the generated lepton asymmetry ($\delta = 1$ corresponds to maximal efficiency). The value of δ is determined from solving the Boltzmann equations for the involved number densities. In the sections below, we will discuss in detail the different parameters entering the relation above: the CP asymmetry ε in Sec. 3.3.1, the efficiency factor δ in Sec. 3.3.2, and C_{sph} in Sec. 3.3.3.

For thermal leptogenesis to work, the temperature of the thermal bath must be higher than the mass of the lightest RHN for their thermal production, thus requiring the reheating temperature $T_{\text{rh}} > M_1$. Given that the seesaw mechanism typically sets $M_1 \gtrsim \mathcal{O}(10^9 \text{ GeV})$ (as we discuss later), this implies a very high reheating temperature. Such high temperatures pose challenges in supersymmetric theories due to the overproduction of gravitinos, the supersymmetric partners of gravitons [163, 164]. This issue can be resolved if the RHNs are produced non-thermally. For example, the RHNs are produced in the decays of a heavy scalar field (which could be the inflaton), and then they decay promptly [165–167]. This mechanism is known as **non-thermal leptogenesis**.

While the discussion above focused on type-I seesaw leptogenesis, leptogenesis can also occur in other seesaw frameworks. Although the core principles remain the same, qualitative differences arise depending on the model. For instance, the triplet scalars and fermions present in type-II and type-III seesaw models possess gauge interactions absent in type-I [168]. Below, we review the common features shared by all leptogen-

esis scenarios, using type-I seesaw as a representative example.

3.3.1 CP asymmetry

In the type-I seesaw framework, the lepton asymmetry arises from the CP-violating, out-of-equilibrium decays of heavy Majorana neutrinos N_i . This CP asymmetry is defined as

$$\varepsilon_i \equiv \sum_{\alpha} \frac{\Gamma(N_i \rightarrow L_{\alpha} + H) - \Gamma(N_i \rightarrow \bar{L}_{\alpha} + \bar{H})}{\Gamma(N_i \rightarrow L_{\alpha} + H) + \Gamma(N_i \rightarrow \bar{L}_{\alpha} + \bar{H})}, \quad (3.8)$$

where α is the lepton flavor index, $i = 1, 2, 3$, and summation over $SU(2)_L$ indices is implicit. This expression shows that to generate a net CP asymmetry, the heavy neutrino must have a decay preference for leptons over antileptons (or vice versa).

However, CPT invariance forbids any asymmetry at tree level since the decay rates for $N \rightarrow LH$ and $N \rightarrow \bar{L}\bar{H}$ are identical at that order. Consequently, one must consider one-loop corrections, as depicted in Fig. 3.2. The CP asymmetry arises from the interference between the tree-level and one-loop amplitudes, which can be made explicit by expressing the decay rate as

$$\Gamma(N_i \rightarrow L_{\alpha} H) = \int \left| \mathcal{M}_{\text{tree}}^2 + \mathcal{M}_{\text{loop}} \mathcal{M}_{\text{tree}}^{\dagger} + \mathcal{M}_{\text{loop}}^{\dagger} \mathcal{M}_{\text{tree}} + \dots \right|, \quad (3.9)$$

where the integral is taken over the phase space. Expressing the amplitudes in terms of the Yukawa couplings, y , and the kinematic factors, I , which include the phase space of the final states as well as the internal loop momentum due to the exchange of N_j (and L_{β}) in the decay of N_i (see Fig. 3.2), we obtain

$$\begin{aligned} \Gamma(N_i \rightarrow L_{\alpha} H) &= |y_{\alpha i}|^2 I_{\text{tree}} + y_{\alpha i}^* y_{\alpha j} y_{\beta j} y_{\beta i}^* I_{\text{loop}} + y_{\alpha i} y_{\alpha j}^* y_{\beta j}^* y_{\beta i} I_{\text{loop}}^* + \mathcal{O}(y^6) \\ \Gamma(N_i \rightarrow \bar{L}_{\alpha} \bar{H}) &= |y_{\alpha i}|^2 I_{\text{tree}} + y_{\alpha i} y_{\alpha j}^* y_{\beta j}^* y_{\beta i} I_{\text{loop}} + y_{\alpha i}^* y_{\alpha j} y_{\beta j} y_{\beta i}^* I_{\text{loop}}^* + \mathcal{O}(y^6). \end{aligned} \quad (3.10)$$

For $I_{\text{loop}} = 0$, we get $\Gamma(N \rightarrow LH) = \Gamma(N \rightarrow \bar{L}\bar{H})$, and thus $\varepsilon = 0$. Substituting the expression above in eq. (3.8), we get

$$\begin{aligned} \varepsilon_i &= \frac{1}{\Gamma_{\text{total}}} (\mathbf{Y} - \mathbf{Y}^*) (I_{\text{loop}} - I_{\text{loop}}^*) = \frac{1}{\Gamma_{\text{total}}} 2i \text{Im}(\mathbf{Y}) 2i \text{Im}(I_{\text{loop}}) \\ &= -\frac{4}{\Gamma_{\text{total}}} \text{Im}(\mathbf{Y}) \text{Im}(I_{\text{loop}}), \end{aligned} \quad (3.11)$$

where $\mathbf{Y} \equiv y_{\alpha i}^* y_{\alpha j} y_{\beta j} y_{\beta i}^*$ and $\Gamma_{\text{total}} = \Gamma(N_i \rightarrow L_{\alpha} + H) + \Gamma(N_i \rightarrow \bar{L}_{\alpha} + \bar{H}) = 2 |y_{\alpha i}|^2 I_{\text{tree}}$.

A few comments are in order. We immediately see that the decay amplitudes require at least two different weak phases and two different strong phases. The weak phases arise from the complex Yukawa couplings, necessitating at least two generations of N . The strong phases originate from the loop factors; specifically, if the intermediate particles

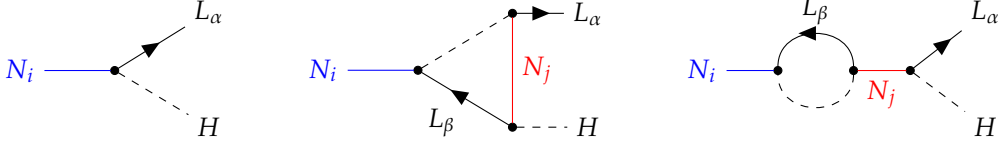


Figure 3.2: Tree level and one loop diagrams for the decay $N_i \rightarrow L_\alpha H$ in type-I leptogenesis. The CP asymmetry is calculated from the interference of the tree level diagram (left) with the vertex correction (center) and the self-energy diagrams (right).

in the loop can go on-shell, the Feynman amplitude acquires an imaginary part due to unitarity and the optical theorem. This latter condition is easily satisfied since the decaying heavy neutrino is much heavier than the SM lepton and Higgs doublets. As evident from eq. (3.11), we require only the imaginary part of the interference of the loop and tree diagrams.⁴

Hence, we have seen that the CP asymmetry in the heavy RHN decays is obtained from the imaginary part of the interference between the tree and one-loop amplitudes. An explicit derivation of the contribution from the vertex and the self-energy corrections to the CP asymmetry ε_i has been provided in Appendix B. Here, we write the final result

$$\varepsilon_i = \frac{1}{8\pi} \sum_{j \neq i} \frac{\text{Im}[y_{\alpha i}^* y_{\alpha j} (y^\dagger y)_{ij}]}{(y^\dagger y)_{ii}} \left[f_v \left(\frac{M_j^2}{M_i^2} \right) + f_s \left(\frac{M_j^2}{M_i^2} \right) \right], \quad (3.12)$$

where

$$f_v(x) = \sqrt{x} \left[1 - (1+x) \ln \left(\frac{1+x}{x} \right) \right]; \quad f_s(x) = \frac{\sqrt{x}}{1-x}, \quad (3.13)$$

are the loop factor associated with the vertex correction and the self-energy diagrams, respectively.⁵ We can obtain the CP asymmetry in the flavour-blind scenario (summed over flavour α) as

$$\varepsilon_i = \frac{1}{8\pi} \sum_{j \neq i} \frac{\text{Im}[(y^\dagger y)_{ij}^2]}{(y^\dagger y)_{ii}} \left[f_v \left(\frac{M_j^2}{M_i^2} \right) + f_s \left(\frac{M_j^2}{M_i^2} \right) \right]. \quad (3.14)$$

Further, in the case of a hierarchical spectrum where only the lightest RHN N_1 decays are considered, the loop factors can be simplified as

$$f_v \left(\frac{M_j^2}{M_1^2} \right) + f_s \left(\frac{M_j^2}{M_1^2} \right) \simeq -\frac{3M_1}{2M_j} \quad \text{for } M_j^2 \gg M_1^2, \quad (3.15)$$

⁴ A simple way to compute this imaginary part is to use Cutkosky's cutting rules [169], which helps in isolating and computing the imaginary part of an amplitude.

⁵ Note that for $M_j \simeq M_i$, the denominator of $f_s \rightarrow 0$, or in other words, the contribution of the self-energy diagram will be maximal. This feature is utilized with a careful treatment in resonant leptogenesis scenarios [170–172].

and the total CP asymmetry generated in N_1 decays becomes

$$\varepsilon_1 \simeq -\frac{3}{16\pi} \sum_{j=2,3} \frac{\text{Im}[(y^\dagger y)_{1j}^2]}{(y^\dagger y)_{11}} \left(\frac{M_1}{M_j} \right). \quad (3.16)$$

This is the simplest approximation for the CP asymmetry, applicable to so-called *vanilla* leptogenesis, which assumes a hierarchical spectrum for M_i , a flavour-blind approach, and thermal production of the heavy RHNs.

Upper bound on CP asymmetry

In the previous chapter, we mentioned how the Yukawas of the seesaw model can be related to the neutrino oscillation data using the Casas-Ibarra (CI) parametrization [105]. Using the latter, it is possible to derive an upper bound on the CP asymmetry produced in N_1 decays. Consider the light neutrino masses in SSI

$$m_\nu = y M_N^{-1} y^T v^2 / 2, \quad (3.17)$$

where M_N is the RHN mass matrix and $v = 246$ GeV is the Higgs VEV. Let us choose a diagonal basis for the RHNs, so we have $M_N = \text{diag}(M_1, M_2, M_3) \equiv D_M$. On the other hand, we can go to the mass basis of light neutrinos using the PMNS matrix of eq. (2.1): $U^T m_\nu U = \text{diag}(m_1, m_2, m_3) \equiv D_m$, to write

$$D_m = \frac{v^2}{2} U^T y D_M^{-1} y^T U = \frac{v^2}{2} U^T y D_M^{-1/2} D_M^{-1/2} y^T U. \quad (3.18)$$

Multiplying the above equation by $D_m^{-1/2}$ from both the left and right sides, we get:

$$\mathbf{1} = \left(\frac{v}{\sqrt{2}} D_M^{-1/2} y^T U D_m^{-1/2} \right)^T \left(\frac{v}{\sqrt{2}} D_M^{-1/2} y^T U D_m^{-1/2} \right) = \Omega^T \Omega, \quad (3.19)$$

where $\Omega = (D_M^{-1/2} y^T U D_m^{-1/2}) v / \sqrt{2}$ is a complex orthogonal matrix. Now, we can write the coupling matrix y in terms of Ω as:

$$y = \frac{\sqrt{2}}{v} U^* D_m^{1/2} \Omega^T D_M^{1/2}. \quad (3.20)$$

This is known as the CI-parametrization for the Yukawa couplings, using which we can write eq. (3.16) as

$$\varepsilon_1 \simeq -\frac{3}{16\pi} \frac{M_1}{v^2} \left(\frac{\sum_j m_j^2 \text{Im}[\Omega_{1j}^2]}{\sum_j m_j |\Omega_{1j}|^2} \right). \quad (3.21)$$

Using the orthogonality condition of the Ω matrix i.e., $\sum_j \Omega_{1j}^2 = 1$, we get

$$|\varepsilon_1| \lesssim \frac{3}{16\pi} \frac{M_1}{v^2} (m_3 - m_1) = |\varepsilon_1| \lesssim \frac{3}{16\pi} \frac{M_1 (\Delta m_{\text{atm}}^2)^{1/2}}{v^2}, \quad (3.22)$$

where in the last step we assume normal ordering. We can use this to translate it into a lower bound for M_1 using eq. (3.7)

$$M_1 \gtrsim Y_{\Delta_B} \left(C_{\text{sph}} \frac{3}{16\pi} \frac{(\Delta m_{\text{atm}}^2)^{1/2}}{v^2} \delta Y_N^{\text{eq}} \right)^{-1}. \quad (3.23)$$

Next, let us plug in some values to get an order of magnitude estimate for the RHN mass. Taking $m_3 = \sqrt{\Delta m_{\text{atm}}^2} = 0.05$ eV, $v = 246$ GeV and $\delta = 1$, we obtain $M_1 > 10^9$ GeV. This is the famous David-Ibarra (DI) bound on the lightest RHN mass [173].

Hence, we have derived a lower bound on the mass of the lightest RHN required for successful vanilla leptogenesis and neutrino masses. Note that this scale is once again very high and challenging to test experimentally. This motivates the development of low-scale seesaw models, some of which were discussed in the previous chapter; a subset of them can also accommodate low-scale leptogenesis scenarios. However, it is important to emphasize that the above lower bound on M_1 applies only to strongly hierarchical heavy neutrino spectra; in the case of quasi-degenerate spectrum of RHNs, the bound can be lowered to the TeV scale.

Moreover, the assumption of maximal efficiency does not hold across all leptogenesis scenarios, particularly when $SU(2)_L$ multiplets with gauge interactions are involved. Such gauge interactions tend to suppress the population of these states at early times and thus the efficiency of leptogenesis, leading to bounds like $M_{\text{triplets}} \gtrsim \mathcal{O}(10^{10} \text{ GeV})$ [168]. In Ref. [3], we studied type-III leptogenesis with an additional Higgs doublet that acquires a small VEV, which enhances the efficiency and lowers the scale to $\mathcal{O}(10^6 \text{ GeV})$.

3.3.2 Boltzmann equations

Since the lepton asymmetry depends on the number densities of leptons, antileptons, and heavy Majorana neutrinos, it is essential to carefully track their abundances and evolution by solving the relevant Boltzmann Equations (BEs). The detailed formalism of BEs is presented in Appendix C. Here, we write the generic form of the BE for a particle a and mass m_a involved in the interaction $a + B \leftrightarrow C$, assuming Maxwell-Boltzmann distributions for all species involved:

$$\frac{dY_a}{dx} = \frac{1}{sHx} \sum_{B,C} \left[\frac{Y_C}{Y_C^{\text{eq}}} \gamma(C \rightarrow a + B) - \frac{Y_a}{Y_a^{\text{eq}}} \cdot \frac{Y_B}{Y_B^{\text{eq}}} \gamma(a + B \rightarrow C) \right], \quad (3.24)$$

Table 3.2: Different processes contributing to the evolution of the number density of N_1 and lepton asymmetry ΔL .

Class	Process	Modifies Y_N	Modifies Y_L	Washout
Decay	$N \rightarrow LH, N \rightarrow \overline{LH}$	Yes	Yes	No
Inverse Decay	$LH \rightarrow N, \overline{LH} \rightarrow N$	Yes	Yes	Yes
$\Delta L = 1$ Scattering	$NL \leftrightarrow Q\overline{u}, NQ \leftrightarrow Lu, N\overline{u} \leftrightarrow L\overline{Q}$	Yes	Yes	Yes
$\Delta L = 2$ Scattering	$LH \leftrightarrow \overline{LH}, LL \leftrightarrow \overline{HH}$	No	Yes	Yes

where $x = m_a/T$ is a parameter that tracks the temperature evolution of the universe, $Y = n/s$ is the yield, s is the entropy density, and H is the Hubble parameter in the radiation-dominated era, and γ is the thermally averaged reaction rate density. All the definitions can be found in Appendix C.

In the case of $N \rightarrow LH$ decays, we have⁶

$$\gamma_D = \gamma(N \rightarrow LH) + \gamma(N \rightarrow \overline{LH}) = n_N^{\text{eq}} \frac{K_1(z)}{K_2(z)} \Gamma_D, \quad (3.25)$$

where K_1 and K_2 are the modified Bessel functions of the first and second kind, respectively. The ratio K_1/K_2 can be interpreted as a time dilation factor. The total decay width Γ_D is given by

$$\Gamma_{\text{tot}} = \Gamma_D \equiv \Gamma(N \rightarrow L + H) + \Gamma(N \rightarrow \overline{L} + \overline{H}) = \frac{(y^\dagger y)_{11}}{8\pi} M_1. \quad (3.26)$$

Now, let us see the processes that one needs to consider for the BEs. Primarily, we have the decays and the inverse decays which source the asymmetry and contribute to washout, respectively. Additionally, we also need to take into account the scattering processes, which are categorized as $\Delta L = 1$ scatterings mediated via Higgs exchange ($\gamma_{H,s}$ and $\gamma_{H,t}$), and $\Delta L = 2$ scatterings mediated via N_1 exchange ($\gamma_{N,s}$ and $\gamma_{N,t}$).⁷ Note that all the processes except $\Delta L = 2$ scatterings modify N_1 abundance. We summarize these processes in Table 3.2 and show the scatterings are in Fig. 3.3.

Decays and inverse decays

First, we focus on the simple case of decays and inverse decays, which occur at order $\mathcal{O}(y^2)$. These processes are usually dominant compared to scatterings, which are either phase-space suppressed or scale as $\mathcal{O}(y^4)$.

Due to the CP asymmetry generated in the decays (see eq. (3.8)), the reaction rates for the inverse decays differ from those of the decays. As CPT invariance states that any physical process is indistinguishable from its time-reversed CP conjugate process, we

⁶ Here, we drop the generation subscript on N , so $N \equiv N_1$ is implied.

⁷ We also have $\Delta L = 1$ scatterings involving gauge bosons such as $NL \rightarrow \overline{H} + V_\mu$ and $N\overline{H} \rightarrow LV_\mu$ with $V_\mu = W_\mu^i, B_\mu$, with reaction rates given by $\gamma_{V,s}$ and $\gamma_{V,t}$. These processes are usually sub-dominant compared to the ones mentioned in Table 3.2.

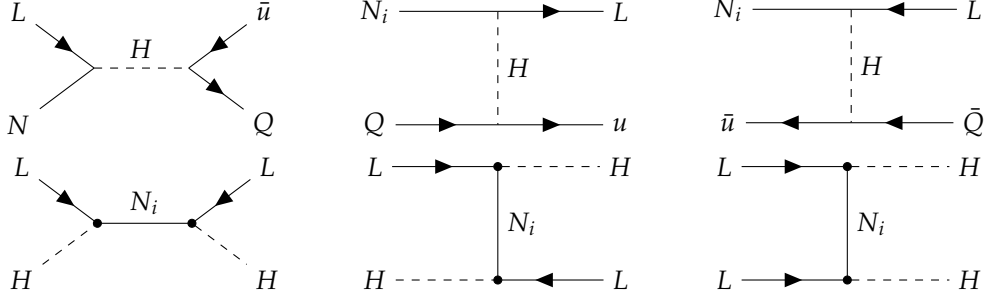


Figure 3.3: The Feynman diagrams for $\Delta L = 1$ and $\Delta L = 2$ scattering processes shown in the first and second row, respectively, where u denotes up-type quarks. Only the processes involving top quarks are dominant due to their large Yukawa coupling.

can write

$$\begin{aligned}\gamma(N \rightarrow LH) &= \gamma(\overline{LH} \rightarrow N) = \frac{1}{2}(1 + \varepsilon)\gamma_D, \\ \gamma(N \rightarrow \overline{LH}) &= \gamma(LH \rightarrow N) = \frac{1}{2}(1 - \varepsilon)\gamma_D.\end{aligned}\quad (3.27)$$

For $\varepsilon \rightarrow 0$, we see that all the four rates above become equal. Using eq. (3.24), we can write the BE for N abundance as

$$\begin{aligned}\frac{dY_N}{dz} &= -\frac{1}{sHz} \left[\frac{Y_N}{Y_N^{\text{eq}}} \left(\gamma(N \rightarrow LH) + \gamma(N \rightarrow \overline{LH}) \right) \right. \\ &\quad \left. - \frac{Y_L}{Y_L^{\text{eq}}} \frac{Y_H}{Y_H^{\text{eq}}} \gamma(LH \rightarrow N) - \frac{Y_{\bar{L}}}{Y_{\bar{L}}^{\text{eq}}} \frac{Y_{\bar{H}}}{Y_{\bar{H}}^{\text{eq}}} \gamma(\overline{LH} \rightarrow N) \right],\end{aligned}\quad (3.28)$$

where $z \equiv M_1/T$. We can simplify the above equation for $T \geq M_1 \gg m_{L,H}$ by using $Y_X \simeq Y_X^{\text{eq}}$ ($X = L, \bar{L}, H$), and the rates in eq. (3.27). We further use $(Y_L + Y_{\bar{L}})/Y_L^{\text{eq}} \simeq 2$. Hence, the equation that tracks the abundance of N , as well as its departure from thermal equilibrium becomes

$$\frac{dY_N}{dz} = -\frac{1}{sHz} \gamma_D \left[\frac{Y_N}{Y_N^{\text{eq}}} - 1 \right]. \quad (3.29)$$

Next, we can write the BEs for the lepton and antilepton yields $Y_{L,\bar{L}}$, considering only the decays and inverse decays

$$\begin{aligned}\frac{dY_L}{dz} &= \frac{1}{sHz} \left[\frac{Y_N}{Y_N^{\text{eq}}} \gamma(N \rightarrow LH) - \frac{Y_L}{Y_L^{\text{eq}}} \frac{Y_H}{Y_H^{\text{eq}}} \gamma(LH \rightarrow N) \right], \\ \frac{dY_{\bar{L}}}{dz} &= \frac{1}{sHz} \left[\frac{Y_N}{Y_N^{\text{eq}}} \gamma(N \rightarrow \overline{LH}) - \frac{Y_{\bar{L}}}{Y_{\bar{L}}^{\text{eq}}} \frac{Y_{\bar{H}}}{Y_{\bar{H}}^{\text{eq}}} \gamma(\overline{LH} \rightarrow N) \right].\end{aligned}\quad (3.30)$$

Subtracting the latter from the former, gives us the BE for ΔL

$$\frac{dY_{\Delta L}}{dz} = \frac{1}{sHz} \gamma_D \left[\varepsilon \left(\frac{Y_N}{Y_N^{\text{eq}}} + 1 \right) - \frac{Y_{\Delta L}}{2Y_L^{\text{eq}}} \right]. \quad (3.31)$$

Notice the erroneous '+' sign in the round brackets multiplied by ε which sources the asymmetry. It suggests that a net lepton asymmetry can be generated even in thermal equilibrium, a clear violation of Sakharov's third condition (which means we have done something wrong!). Here lies a subtlety: we saw that the CP asymmetry is proportional to y^4 , therefore we must include the $\Delta L = 2$ scatterings mediated by N in the s -channel $\gamma_{N,s}$, that are also $\mathcal{O}(y^4)$. However, their on-shell contribution has already been taken into account by successive inverse decay and decays, therefore, we only need the off-shell contribution. We can estimate the on-shell contribution as

$$\gamma^{\text{on-shell}}(LH \rightarrow \overline{LH}) = \gamma(LH \rightarrow N) \text{Br}(N \rightarrow \overline{LH}), \quad (3.32)$$

where we can use $\text{Br}(N \rightarrow \overline{LH}) = (1 - \varepsilon)/2$. These need to be subtracted from the s -channel N exchange ($\gamma_{N,s}$) to get the off-shell contribution [174]:

$$\begin{aligned} \gamma'(LH \rightarrow \overline{LH}) &= \gamma_{N,s} - \frac{(1 - \varepsilon)^2}{4} \gamma_D \\ \gamma'(\overline{LH} \rightarrow LH) &= \gamma_{N,s} - \frac{(1 + \varepsilon)^2}{4} \gamma_D. \end{aligned} \quad (3.33)$$

Including these off-shell contributions and ignoring $\mathcal{O}(\varepsilon^2)$ terms, we can arrive at the correct BE for lepton asymmetry

$$\frac{dY_{\Delta L}}{dz} = \frac{1}{sHz} \gamma_D \left[\varepsilon \left(\frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) - \frac{Y_{\Delta L}}{2Y_L^{\text{eq}}} \right] - 2\gamma' \frac{Y_{\Delta L}}{Y_L^{\text{eq}}}. \quad (3.34)$$

While the BEs derived above suffice for our purposes, one can similarly include the contributions from $\Delta L = 1$ scatterings to arrive at the complete set of BEs for thermal leptogenesis with hierarchical RHNs

$$\begin{aligned} \frac{dY_N}{dz} &= -\frac{1}{sHz} \left[\frac{Y_N}{Y_N^{\text{eq}}} - 1 \right] (\gamma_D + \gamma_{H,s} + 4\gamma_{H,t}), \\ \frac{dY_{\Delta L}}{dz} &= \frac{1}{sHz} \gamma_D \left[\varepsilon \left(\frac{Y_N}{Y_N^{\text{eq}}} - 1 \right) - \frac{Y_{\Delta L}}{2Y_L^{\text{eq}}} \right] - \frac{Y_{\Delta L}}{Y_L^{\text{eq}}} (2\gamma' + 2\gamma_{H,t} + \gamma_{H,s} \frac{Y_N}{Y_N^{\text{eq}}}). \end{aligned} \quad (3.35)$$

The expressions for the respective scattering rates can be found in Ref. [174]. If you want to include gauge-boson-mediated scatterings, you can make the following sub-

stitutions

$$\begin{aligned}\gamma_{H,s} &\rightarrow \gamma_{S_s} \equiv \gamma_{H,s} + \gamma_{V,s} , \\ \gamma_{H,t} &\rightarrow \gamma_{S_t} \equiv \gamma_{H,t} + \gamma_{V,t} .\end{aligned}\tag{3.36}$$

Washouts and efficiency

We can solve the coupled Boltzmann equations of eqs. (3.29) and (3.34) numerically in the case of vanilla leptogenesis to determine the efficiency. Alternatively, we can dig into the out-of-equilibrium dynamics to get some analytical estimates. Let us define the decay parameter as

$$K \equiv \frac{\Gamma_{\text{total}}}{H(z=1)} = \frac{\tilde{m}}{m^*} ,\tag{3.37}$$

where $\tilde{m}$ and m^* are the effective and equilibrium neutrino mass parameters, respectively, defined as [175]:

$$\tilde{m} \equiv \frac{8\pi v^2}{2M_1^2} \Gamma_D = \frac{(y^\dagger y)_{11} v^2}{2M_1} , \quad m^* \equiv \frac{8\pi v^2}{2M_1^2} H(z=1) \simeq 1.08 \times 10^{-3} \text{eV} ,\tag{3.38}$$

where Γ_D is defined in eq. (3.26). Depending on the value of K , we have two regimes:

1. Weak washout regime ($\tilde{m} \ll m_*$):

N decays out of equilibrium and the $\Delta L = 1, 2$ washout processes are inefficient. The dominant contribution to the lepton asymmetry comes only from decays and inverse decays, which are efficient at $T \sim M_1$. While a large lepton asymmetry can be generated in this regime, it is sensitive to the initial conditions, i.e., whether N 's have thermal or vanishing initial abundance. The efficiency factor is estimated to be

$$\delta \simeq \begin{cases} 1 & \text{for } Y_N(T \gg M_1) \simeq Y_N^{\text{eq}}(T \gg M_1) \\ \tilde{m}^2/m^{*2} & \text{for } Y_N(T \gg M_1) \simeq 0 \end{cases} .\tag{3.39}$$

2. Strong washout regime ($\tilde{m} \gg m_*$): The final lepton asymmetry is independent of the initial conditions and the total lepton asymmetry production is highly suppressed at $T \sim M_1$. Any asymmetry produced earlier is washed out completely and one gets $Y_L = 0$, since no asymmetry can be generated in equilibrium. A net asymmetry survives once the inverse decay processes go out of equilibrium. In this case, the efficiency factor is estimated as $\delta \simeq m^*/\tilde{m}$. The numerical solutions of eq. (3.28) and eq. (3.34) for a choice of benchmark values are shown in Fig. 3.4.

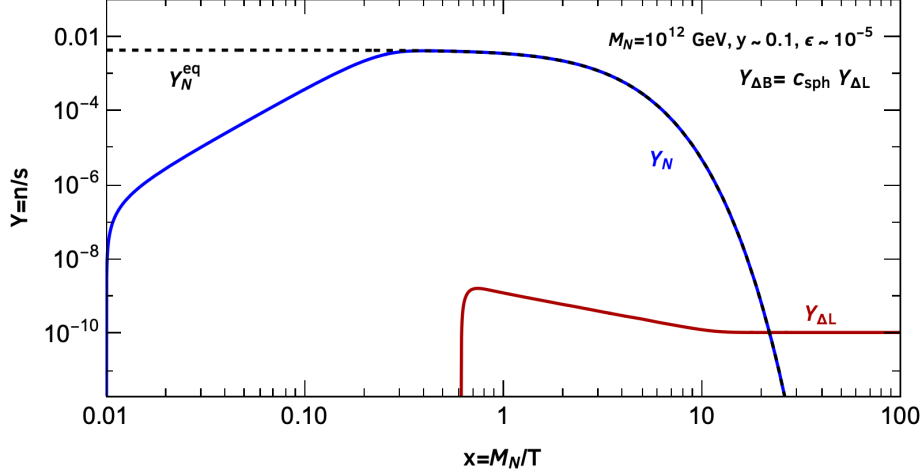


Figure 3.4: The numerical solutions of the BEs for the evolution of N abundance (blue) and lepton asymmetry (red) in the strong washout regime. The dashed black curve represents the N equilibrium yield. We fix $M_1 = 10^{12}$ GeV, $|y| = 0.1$ and $\varepsilon = 10^{-5}$.

3.3.3 Sphaleron processes

We now turn our attention to the processes responsible for converting the lepton asymmetry into a baryon asymmetry, which also play a crucial role in electroweak baryogenesis scenarios. In the renormalizable SM Lagrangian, baryon number B and lepton number L (for each flavour) are conserved. However, non-perturbative gauge field configurations known as sphalerons induce violations of B and L . Although these effects are typically negligible at low energies, they become significant at the very high temperatures of the early universe ($T \gg T_{\text{EWPT}}$), leading to rapid violation of $B + L$. Here, we discuss the anomalous $B + L$ violation that is crucial during the leptogenesis epoch and derive the sphaleron conversion that processes a net lepton asymmetry into a baryon asymmetry.

Anomalous $B+L$ violation

The global $U(1)_B$ and $U(1)_L$ symmetries of SM lead to the conservation of baryon and lepton number, respectively. We have

$$B = \int d^3x J_0^B(x), \quad L = \int d^3x J_0^L(x), \quad (3.40)$$

where the associated currents are given by

$$\begin{aligned} J_\mu^B &= \sum_i (\bar{Q}_{L,i} \gamma_\mu Q_{L,i} + \bar{u}_{R,i} \gamma_\mu u_{R,i} + \bar{d}_{R,i} \gamma_\mu d_{R,i}), \\ J_\mu^L &= \sum_i (\bar{L}_i \gamma_\mu L_i + \bar{e}_{R,i} \gamma_\mu e_{R,i}). \end{aligned} \quad (3.41)$$

Here, Q_L , L are the left-handed quark and lepton doublets, u_R , d_R , e_R are the right-handed up-type, down-type and charged lepton singlets, and summation over color

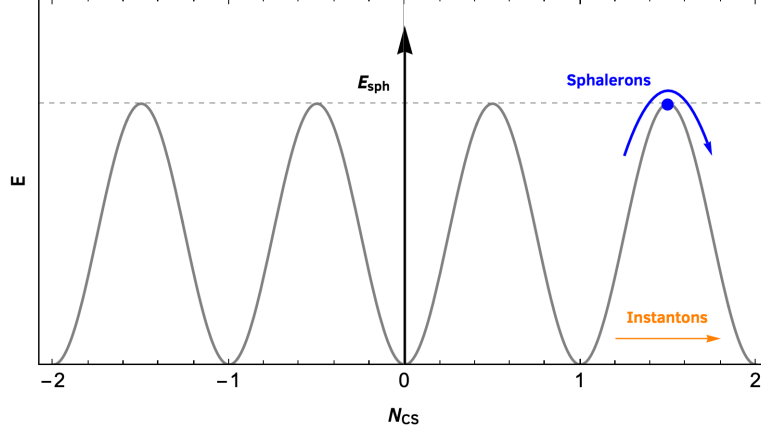


Figure 3.5: The vacuum energy as a function of Chern-Simons number N_{CS} , which are positive and negative integers. The blue and the orange curves denote the sphaleron and instanton configurations that change N_{CS} .

indices is implied.

Now, we know from the famous Noether's theorem [176], that the classical currents associated with these symmetries are conserved, i.e., $\partial_\mu J_\mu^B = \partial_\mu J_\mu^L = 0$. However, the chiral nature of the electroweak theory leads to the Adler-Jackiw-Bell triangle anomaly [177, 178] and the divergences associated with these currents do not vanish. This makes B and L anomalous at the quantum level. Upon analysing the anomaly equations for the left-handed and right-handed chiral components of the vector currents (for the SM gauge group), we get

$$\partial^\mu J_\mu^B = \partial^\mu J_\mu^L = \frac{N_f}{32\pi^2} \left(-g^2 W_{\mu\nu}^i \tilde{W}_{\mu\nu}^i + g'^2 B_{\mu\nu} \tilde{B}^{\mu\nu} \right), \quad (3.42)$$

where the gauge fields and couplings have the same definition as in Chapter 1. N_f denotes the number of flavors. We note two things from the equation above: $\partial^\mu (J_\mu^B - J_\mu^L) = 0$, so $B - L$ is exactly conserved and $B + L$ is violated

$$\begin{aligned} \Delta(B + L) &= \int d^4x \partial^\mu (J_\mu^B + J_\mu^L) = \int d^4x \frac{2N_f}{32\pi^2} \left(-g^2 W_{\mu\nu}^i \tilde{W}_{\mu\nu}^i + g'^2 B_{\mu\nu} \tilde{B}^{\mu\nu} \right) \\ &= 2N_f \Delta N_{CS}. \end{aligned} \quad (3.43)$$

where $\Delta N_{CS} = \pm 1, \pm 2, \dots$ is the Chern-Simons number and is related to the topology of the involved gauge fields [179]. The non-zero value of $\Delta(B + L)$ is a consequence of the non-Abelian structure of the gauge theory, whose vacuum structure consists of infinitely many degenerate distinct ground states separated by a potential barrier E_{sph} , as shown in Fig. 3.5. The number N_{CS} classifies the transitions among these ground states corresponding to the jumps between different topological sectors of the gauge theory [149, 180, 181]. Therefore, a transition between the two different vacua will change N_{CS} , which in turn changes B and L . The sphalerons mediate these transitions over the potential barrier [182]. We also show another non-perturbative solution

called *instantons* which determine the probability of tunneling between the neighbouring vacua in Fig. 3.5.

From eq. (3.43), we see that for 3 SM flavours, $\Delta(B + L) = \pm 6$ for $\Delta N_{\text{CS}} = \pm 1$ or $\Delta B = \Delta L = \pm 3\Delta N_{\text{CS}}$. Hence, the transition from one vacuum to another changes $\Delta B, \Delta L$ in multiples of 3. The configurations give rise to a 12-fermion interaction involving all three families, such as

$$\bar{u} + \bar{d} + \bar{c} \rightarrow d + 2s + 2b + t + \nu_e + \nu_\mu + \nu_\tau. \quad (3.44)$$

The associated instanton action is large and the transition rate

$$\Gamma^{T=0} \sim e^{S_{\text{instanton}}} = e^{-4\pi/\alpha_W} = \mathcal{O}(10^{-165}), \quad (3.45)$$

with $\alpha_W = g^2/4\pi$, is highly suppressed at zero temperatures. Therefore, the mediated $B + L$ violation is negligible [183].

However, for finite temperatures, it has been shown in Refs. [152, 184] that the transitions between vacua can be made by thermal fluctuations over the barrier instead of tunnelling, i.e., via sphalerons. When the temperature is larger than the height of the barrier, the Boltzmann suppression vanishes and the $B + L$ violating interactions can occur at a significant rate and are in thermal equilibrium.

The rate of the sphaleron transitions at a finite temperature is determined by the height of the potential barrier E_{sph} , which is proportional to the Higgs VEV at that finite temperature T . For $T < T_{EW} \simeq 100$ GeV, this rate is Boltzmann suppressed, $\Gamma \sim \exp(-E_{\text{sph}}/T)$ [184], and these transitions are suppressed. On the other hand, for temperatures $T \geq T_{EW}$, the transition rate goes as $\mathcal{O}(\alpha_W^5 T^4)$ [185]. Therefore, for temperatures higher than T_{EW} , these baryon number violating sphaleron processes are unsuppressed and very potent. By comparing these rates with the Hubble expansion rate, we find that the sphalerons are in thermal equilibrium for $100 \text{ GeV} \leq T \leq 10^{12} \text{ GeV}$. In this window, they can reprocess the asymmetries between baryons and leptons.

Let us recall Sakharov's third condition: no asymmetry can be generated in thermal equilibrium. Although sphalerons violate $B + L$, they will erase any net asymmetry in equilibrium at temperatures $T \geq T_{EW}$, unless there is a non-zero $B - L$ asymmetry from another source. This is because sphalerons preserve $B - L$ and therefore cannot erase it. Once sphaleron processes freeze out, they partially convert this $B - L$ asymmetry into a baryon asymmetry. We derive this conversion factor below.

Conversion factor

Consider the formula for the number density of a particle

$$n = g \int \frac{d^3p}{(2\pi)^3} f(E, t) = g \int \frac{d^3p}{(2\pi)^3} \frac{1}{\exp[(E - \mu)/T] \pm 1}, \quad (3.46)$$

where the $+$ ($-$) corresponds to fermions (bosons), g are the internal degrees of freedom, and μ is the chemical potential. In the very early universe, we have $\mu_x/T \ll 1$ and $m_x/T \ll 1$ for a particle species x . For the corresponding antiparticle $\bar{x}$, the sign of μ is flipped. Using eq. (3.46), we can determine the asymmetry in the particle number densities, $\Delta n_x \equiv n_x - n_{\bar{x}}$ as

$$\Delta n_x \simeq \begin{cases} \frac{1}{6} g_x T^3 (\mu_x/T) & \text{for fermions,} \\ \frac{1}{6} g_x T^3 (2\mu_x/T) & \text{for bosons.} \end{cases} \quad (3.47)$$

Starting with N_f generations of fermions and N_H Higgs doublets, we have the left-handed quark doublets Q_i , right-handed up-type (down-type) u_i (d_i) quarks, left-handed lepton doublet L_i and right-handed charged leptons e_i for $i = 1, 2, \dots, N_f$, and the Higgs doublets H_j with $j = 1, 2, \dots, N_H$. At high temperatures, all these particles interact via their gauge and Yukawa couplings. Additionally, we have the non-perturbative sphaleron processes that we discussed above. In the EW unbroken phase at sufficiently high temperatures, all the gauge interactions are in thermal equilibrium. We can take the chemical potential for all the gauge bosons to be 0, while, the components within the doublet have equal chemical potential. Further, we can relate the chemical potentials to each other from processes that are fast enough to be in equilibrium. See the relations below

1. As a consequence of the effective 12 fermion interaction induced by the sphalerons, we have

$$\sum_i (3\mu_{Q_i} + \mu_{L_i}) = 0. \quad (3.48)$$

2. Using the hypercharge conservation of plasma at all temperature, we get

$$\sum_i (\mu_{Q_i} + 2\mu_{u_i} - \mu_{d_i} - \mu_{L_i} - \mu_{e_i} + \frac{2N_H}{N_f} \mu_H) = 0, \quad (3.49)$$

where we have assumed all Higgs doublets to have the same chemical potential, i.e., $\mu_{H_j} = \mu_H$. Note the factor of 2 in front of N_H because of its bosonic nature.

3. The Yukawa interactions between the left-handed and right-handed fermions give the following relations

$$\mu_{Q_i} + \mu_H - \mu_{u_i} = 0, \quad \mu_{Q_i} - \mu_H - \mu_{d_i} = 0, \quad \mu_{L_i} - \mu_H - \mu_{e_i} = 0 \quad (3.50)$$

Finally, using eq. (3.47), we can write the baryon and lepton number densities in terms of the chemical potentials as

$$n_B = \frac{1}{6} T^2 \sum_c \sum_i^{N_f} (2\mu_{Q_i} + \mu_{u_i} + \mu_{d_i}), \quad n_L = \frac{1}{6} T^2 \sum_i^{N_f} (2\mu_{L_i} + \mu_{e_i}), \quad (3.51)$$

where c denotes colour. We further assume that the chemical potentials are flavour independent: $\mu_{Q_i} = \mu_Q$, $\mu_{u_i} = \mu_u$, $\mu_{d_i} = \mu_d$, $\mu_{l_i} = \mu_l$ and $\mu_{e_i} = \mu_e$.

We can now solve the linear equations above and express the chemical potential in terms of, say μ_e , to obtain [186]

$$n_B = \frac{1}{6}T^2 \frac{8N_f + 4N_H}{2N_f + 3N_H} N_f \mu_e, \quad n_L = -\frac{1}{6}T^2 \frac{14N_f + 9N_H}{2N_f + 3N_H} N_f \mu_e. \quad (3.52)$$

We can now finally relate B and $B - L$ as

$$n_B = \frac{8N_f + 4N_H}{22N_f + 13N_H} (n_B - n_L) \equiv C_{\text{sph}} (n_B - n_L). \quad (3.53)$$

In the SM, we have $N_f = 3$ and $N_H = 1$, so the conversion factor C_{sph} comes out to be $28/79$ (if you relate B to L instead, you get $C_{\text{sph}} = 28/51$). Hence, we have obtained all the ingredients for baryogenesis via leptogenesis. Substituting the value of C_{sph} in eq. (3.7), we get

$$Y_{\Delta B} = 8.58 \times 10^{-11} \left(\frac{\varepsilon}{6.2 \times 10^{-6}} \right) \left(\frac{\delta}{10^{-2}} \right). \quad (3.54)$$

3.3.4 Beyond Vanilla leptogenesis

Since its proposal by Fukugita and Yanagida [161], the vanilla leptogenesis scenario—with its hierarchical spectrum of RHNs and flavour-blind approach—has helped us understand the qualitative features of thermal leptogenesis and has been widely used in model building. For example, we apply this mechanism to generate asymmetries in both the visible and dark sectors in Ref. [2], and we can similarly incorporate it in the model of Ref. [1] as both these models are based on the extension of type-I seesaw. It is also worth mentioning the Davidson-Ibarra bound, which requires successful leptogenesis to have $M_1 > 10^9$ GeV [173].

Over the past two-three decades, several modifications to the standard scenario have been developed. These include, but are not limited to: adding quantum effects, thermal corrections, and spectator processes [174, 187, 188]; leptogenesis in type-II/III seesaw models [168, 189–191]; flavour effects [192–196]; N_2 dominated leptogenesis [197–199]; resonant leptogenesis [170, 172]; leptogenesis via oscillations [200–202]; leptogenesis in the scotogenic model [203–206], and so on.⁸

In many of these scenarios, it is possible to lower the leptogenesis scale to $\mathcal{O}(\text{GeV})$, allowing for testable predictions. Below, we briefly discuss flavour effects in leptogenesis and the case of Dirac leptogenesis.

⁸ In Ref. [202], it was shown that resonant leptogenesis and leptogenesis via oscillations may be actually united, corresponding to the freeze-out and freeze-in of the heavy neutrinos, respectively.

Flavour effects

In the discussion above, we did not consider the distinct Yukawa interactions of the charged leptons, which enable us to distinguish between flavours. By summing over all flavours, we arrive at a single equation for the lepton asymmetry in the one-flavour, or flavour-blind, approach. This approximation is justified for $M_1 \geq 10^{12}$ GeV. However, below $T \sim 10^{12}$ GeV, the charged lepton Yukawa interactions become rapid enough to be in thermal equilibrium, breaking the coherence of the lepton doublets produced in N decays. This requires treating each lepton flavour asymmetry separately, such that $\Delta L = \sum_{\alpha} \Delta L_{\alpha}$ with $\alpha = e, \mu, \tau$.

In the temperature range $10^9 \text{ GeV} \leq T \leq 10^{12} \text{ GeV}$, the τ -Yukawa interactions are in equilibrium and we need to treat them separately. However, the e, μ flavours are indistinguishable in this regime, so we are in the two-flavoured regime corresponding to the τ component and the non- τ component. Both the CP asymmetry as well as the BEs with washouts are decomposed in two parts. In Ref. [3], we present a simplified analysis for flavour effects in this regime for triplet fermions.

Below $T \simeq 10^9 \text{ GeV}$, the μ -Yukawa interactions also equilibrate, making all three flavours distinguishable. Here, the washout rates differ for each flavour due to their distinct interactions with the thermal bath. Consequently, flavour effects become crucial for low-scale leptogenesis scenarios and must be properly accounted for, as they can significantly modify the final lepton asymmetry compared to the standard flavour-blind treatment. A rigorous treatment requires the use of the density matrix formalism instead of the classical Boltzmann equations [196, 207, 208].

Dirac leptogenesis

An interesting variation of the standard leptogenesis scenario is Dirac leptogenesis [209], where neutrinos acquire purely Dirac masses, and thus total lepton number is conserved, unlike the Majorana case. In these scenarios, RHNs ν_R are added to the SM. Due to the presence of CP-violating phases, decays of a heavy particle can generate equal and opposite lepton asymmetries in the left- and right-handed neutrino sectors. However, because neutrino masses are extremely small, the interactions involving ν_R are typically too weak to equilibrate before electroweak sphalerons freeze out.

Since the $SU(2)_L$ sphalerons only couple to left-handed fields, they process the lepton asymmetry stored in the left-handed sector into a baryon asymmetry, even when the total initial $(B - L)$ is zero. This mechanism depends on the interplay between the slow left-right equilibration and sphaleron dynamics. Dirac leptogenesis is particularly appealing in models with left-right symmetry: $(SU(2)_R \times SU(2)_L)$, and is realised in various BSM extensions; see Refs. [210–213] for detailed discussions.

DARK MATTER

In the previous chapters, we discussed the matter we can see and interact with—so-called visible matter. The term is somewhat oxymoronic, since these particles are far too small to be seen directly. By visible matter, we mean the matter that interacts with light or the electromagnetic force. Several astrophysical and cosmological observations point to the existence of a new type of matter that cannot be seen (non-luminous), termed Dark Matter (DM).

The term was first coined by the Swiss astrophysicist Fritz Zwicky in 1933. While observing the Coma Cluster and applying the virial theorem, he deduced that some unseen matter must be responsible for the rapid motions of galaxies within the cluster. He called this new form of non-interacting matter *Dunkle Materie* (dark matter). Initially, this proposal was met with scepticism and largely dismissed. It was only in the 1960s and 70s, with improved calculations and instrumentation—thanks in large part to the efforts of Vera Rubin and Kent Ford—that dark matter gained widespread acceptance. Today, we know that dark matter accounts for roughly 26% of the total energy budget of the universe, about five times more than visible matter.

Since all evidence for dark matter arises from its gravitational effects, alternative explanations such as Modified Newtonian Dynamics (MOND) and other modifications to gravity have been proposed. However, there are compelling reasons to believe in the particle nature of dark matter instead, which we elaborate on in the next section. Regarding its particle nature, dark matter should be non-baryonic and cold—hence the ‘CDM’ in the standard Λ CDM cosmological model. Note that for cosmologists, the SM matter content falls under the category of baryonic matter, not just baryons. We will use ‘visible matter’ and ‘baryonic matter’ interchangeably when referring to SM particles. The SM, as we know it, does not contain any particle that can serve as a viable dark matter candidate. This will become clear once we discuss the necessary properties of a viable dark matter candidate in the next section. Therefore, the origin and nature of dark matter not only point to physics beyond the SM but also hold the key to

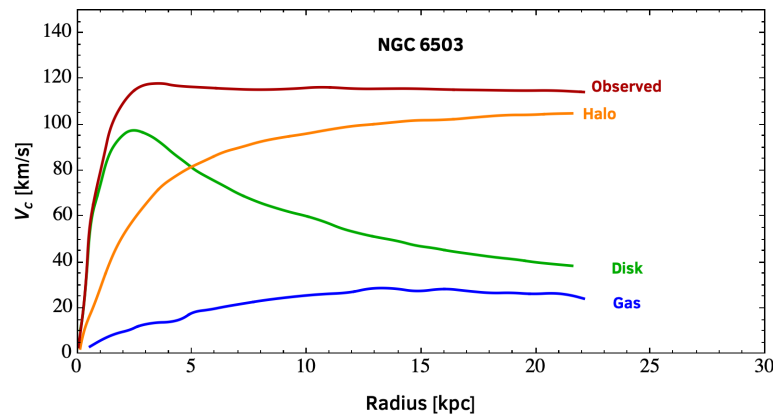


Figure 4.1: Observed rotation curve of the galaxy NGC 6503 (solid red curve) as a function of distance from the galactic core [214]. The contribution from the DM halo, disk, and gas are represented by orange, green, and blue curves, respectively.

understanding what the vast majority of our universe is made of.

In this chapter, we first review the evidence and properties of particle dark matter. We then discuss common production mechanisms and candidate particles, followed by their experimental and observational probes.

4.1 Observations and evidences

Let us briefly review the evidences that we gathered for the existence of DM across several scales.

Galaxy rotation curves

Going back to high-school physics and recalling Gauss’s law, we know that the rotational velocity of objects gravitationally bound to a galaxy is given by $v = \sqrt{GM(r)/r}$, where $M(r)$ is the mass enclosed in the system and r is the distance from the galactic centre. For a spiral galaxy, we can approximate the visible matter as being spherically distributed in the galactic bulge. Therefore, in the outer regions of the galaxy, we expect the rotational speed to decrease because there is little visible matter.

However, the pioneering work of Rubin, Ford, and Freeman showed that the rotation curves fall off very slowly and instead remain approximately flat [215–218] (see Fig. 4.1). This observed flatness cannot be explained by modifying the weights or spatial distribution of the known galactic components [219, 220]. Instead, it points to the presence of an additional, non-luminous matter component with a mass distribution scaling as $M(r) \propto r$. In Fig. 4.1, the contribution from the DM halo is shown in orange. The observed flat curve is well reproduced by summing the contributions from the DM halo and the baryonic disk.¹

¹ The discrepancy in observed rotation curves can also be addressed using MOND [221, 222], an alternative to the particle DM hypothesis. However, MOND fails to explain other key observations, such as those from Bullet Cluster and the CMB.

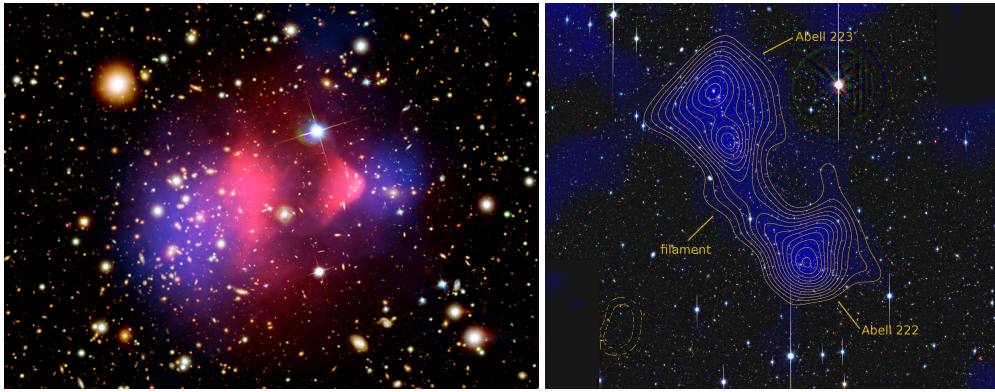


Figure 4.2: Left: Image of the Bullet Cluster, formed by the collision of two large galaxy clusters. The matter distribution inferred from gravitational lensing is shown in blue, while the pink region corresponds to the X-ray emission from the hot gas, which constitutes the luminous matter. Right: Dark matter filaments linking the galaxy cluster Abell 222 and Abell 223.

Galaxy clusters

Masses of astrophysical objects are usually estimated through the effects of gravitational lensing.² Lensing observations across numerous astrophysical systems have consistently confirmed the dominant presence of dark matter in galaxies and galaxy clusters [223–226]. The collision and subsequent merging of two galaxy clusters—as observed in the Bullet Cluster (1E 0657-56), gives the most compelling evidence for the existence of dark DM through gravitational lensing. Weak lensing data reveal a separation between the total gravitational mass and the luminous matter, as shown on the left side of Fig. 4.2. This displacement between the lensing-derived mass distribution and the X-ray luminous gas is naturally explained by the presence of collisionless DM, which does not interact electromagnetically and passes through the collision unaffected. In contrast, MOND cannot account for the gravitational effects observed in regions devoid of visible matter.

Structure formation and CMB

Weak lensing data show that luminous matter has a filamentary structure on extremely large scales, typically $\sim O(100 \text{ Mpc})$. Simulations with cold DM reproduce this observed pattern of filaments with remarkable accuracy. They also support the existence of DM filaments: large-scale, invisible, web-like structures, that connect galaxy clusters and superclusters, forming what is known as the cosmic web. These DM filaments were first detected linking the galaxy clusters Abell 222 and Abell 223, as shown on the right side of Fig. 4.2 [227].

Moreover, DM is essential in the evolution and growth of cosmic structures. In the early universe, density fluctuations in matter led to the clumping of DM before baryonic matter, due to its higher abundance and lack of interactions with radiation. These clumps

² According to general relativity, massive bodies such as galaxy clusters warp spacetime, bending and distorting the light from more distant background objects. As a result, these massive structures act as gravitational lenses.

formed gravitational potential wells, into which baryonic matter later fell. Therefore, DM leaves observable imprints on the CMB, much like baryonic matter (as discussed in the previous chapter). Cosmological measurements from probes such as WMAP and Planck have determined the DM relic density to be [77]

$$\Omega_{DM}h^2 = 0.1200 \pm 0.0012, \quad (4.1)$$

which is approximately five times the abundance of baryonic matter (see eq. (3.5)). This $\mathcal{O}(1)$ ratio between visible and DM densities is known as the *cosmic coincidence*. Explaining this coincidence is a major question in particle cosmology.

4.2 Properties of a DM particle

Having established the evidence for a large DM abundance in our universe, we now turn to its possible properties. We refer the reader to Refs. [228] for a recent comprehensive review on particle DM. While we know that DM interacts gravitationally, it may also have weak or feeble interactions with the SM particles. Here, we outline some general properties that a viable DM candidate should possess:

Neutral

Since DM does not emit or absorb light, it must be electrically neutral. However, it is possible for DM to possess an effective fractional electric charge, or *millicharge* [229, 230], since current constraints are placed on the charge-to-mass ratio [231].³ It is possible to hide the effects of its charge if the DM is ultra-heavy, for example, see Charged Massive Particles (CHAMPs) [234]. Based on null results from exotic atom searches and analyses of the CMB power spectrum, the upper bound on DM charge is given by [235]

$$\epsilon \leq 2.24 \times 10^{-4} \left(\frac{M}{\text{TeV}} \right)^{1/2}. \quad (4.2)$$

Moreover, the DM candidate should also transform as a singlet under $SU(3)_C$, implying that it is colourless and does not interact via strong interactions with the SM baryons. There are stringent constraints on the strong interactions between DM and nucleons from the Earth's heat flow [236].

Non-baryonic

Combined results from the CMB and BBN indicate that baryonic matter constitutes only 5% of the universe's total energy budget. Given the evidence for a significantly larger DM abundance, as discussed above, it is generally concluded that DM must be

³ Millicharged DM particles can interact with electrons and protons during the recombination era, thereby leaving imprints on the CMB [232, 233].

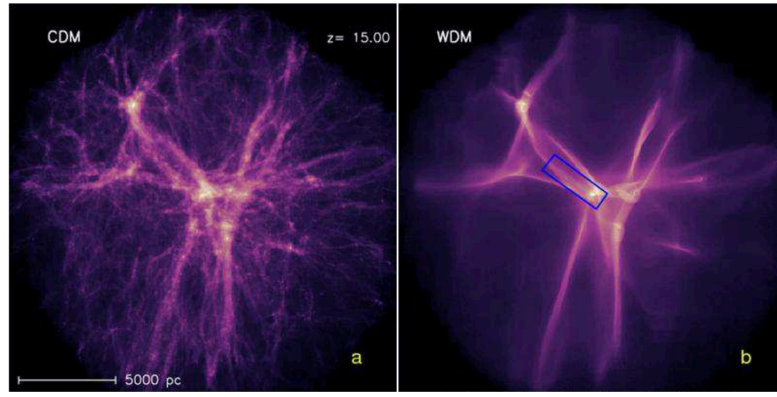


Figure 4.3: A comparison of small-scale structure formation with cold DM (left) and warm DM (right). Taken from Ref. [244].

non-baryonic. In other words, DM must consist of a new type of particle not present in the SM.

In the past, several baryonic DM candidates have been proposed. One class includes **Massive Compact Halo Objects (MACHOs)**, such as brown dwarfs, white dwarfs, and cold gas clouds. However, lensing observations have placed strong constraints on MACHOs as viable DM candidates, indicating that they can account for at most 10-20% of the DM halo mass [237–239]. Another baryonic DM candidate that has gained attention in recent years is the **Primordial Black Hole (PBH)**. Although composed of ordinary matter, PBHs are hypothesized to have formed in the early universe under extreme conditions. However, a wide range of PBH masses have been excluded due to observational constraints [240]. Currently, PBH masses in the range $10^{17} - 10^{23}$ g can be a viable DM candidate [241].

It should be point out that the SM neutrinos being extremely abundant, electrically neutral, and feebly interacting were historically considered as viable DM candidates. However, SM neutrinos are now ruled out as DM for two main reasons: i) They are hot, i.e., relativistic during the epoch of structure formation (see the discussion below), which is inconsistent with the observed bottom-up hierarchy of structure formation; ii) They are too light to account for the observed DM relic abundance.

Cold (Warm)

The observed matter power spectrum implies that DM must be cold (non-relativistic) at the time of decoupling and during the epoch of structure formation. If DM was hot at decoupling as in the case of neutrinos, its large free-streaming length (the average distance a particle travels before becoming gravitationally bound) would wash out small-scale structures and suppress the hierarchical growth of cosmic structures. Hence, the current formulation accepts the CDM model, first proposed by Peebles in 1982 [242] and developed further in Ref. [243].

While hot DM is ruled out by the observed presence of small-scale structures, an in-

intermediate scenario remains viable: Warm Dark Matter (WDM) [245,246]. These are particles that were semi-relativistic at the time of decoupling, typically with masses in the keV range. WDM suppresses structure formation only on very small scales, while still allowing larger structures to form, in agreement with current observations. A comparison of how warm and cold DM affect structure formation is shown in Fig. 4.3.

The WDM models can help address several small-scale issues within the standard Λ CDM framework, such as the *Missing Satellites Problem*, *Core-Cusp Problem* and *Too Big to Fail Problem* on small scales, see Refs. [247,248] for details. We refer the reader to Ref. [249] for a review on WDM.

Collisionless

Several observations suggest that DM does not interact significantly with itself or with baryonic matter. For instance, the Bullet Cluster—discussed earlier, provides strong evidence of this non-interacting behaviour. Such observations are used to place an upper bound on the DM self-interaction cross-section, typically $\sigma/M_{\text{DM}} \lesssim 1 \text{ cm}^2 \text{ g}^{-1}$ at cluster scales [250].

Moreover, the structure of galaxies tells us that while baryonic matter forms a thin disk, DM forms an extended, approximately ellipsoidal halo. This difference points towards the non-dissipative or collisionless nature of DM. In addition, significant interactions between DM and baryons would leave observable imprints on the CMB power spectrum (especially in the acoustic peaks) and would also alter the efficiency of structure formation. However, numerical simulations based on cold, collisionless DM reproduce both the observed CMB anisotropies and the large-scale structure of the universe accurately.

While cold, collisionless DM explains large-scale observations well, it faces challenges on smaller scales, as mentioned above. These tensions can be alleviated by introducing mild DM self-interactions. Models incorporating such interactions fall under the category of **Self-Interacting Dark Matter (SIDM)** [251,252]. In these models, DM particles interact via elastic scattering, with cross-sections in the range: $0.1 < \sigma/M_{\text{DM}} < 10 \text{ cm}^2 \text{ g}^{-1}$ [253]. These interactions are strong enough to address the small-scale issues while remaining consistent with constraints from large-scale structure and CMB observations. For example, SIDM models predict more spherical halos compared to CDM.

Long-lived

Given the observed relic abundance of DM in the universe today, DM must be long-lived on cosmological timescales, if not absolutely stable. For heavier DM candidates, ensuring their stability is a key requirement in BSM frameworks. A common approach is to impose discrete symmetries—such as a $\mathbb{Z}_2$ symmetry—that render the lightest

particle with odd parity stable by forbidding its decay. This mechanism is employed, for example, in the scotogenic model discussed in Section 2.5.2.

Note that the observations do not strictly forbid DM from decaying. We require that DM lifetime exceeds at least the age of the universe, approximately 13.8 billion years.⁴ In fact, DM decay can provide distinctive, *smoking-gun* signatures for new physics. For example, if the decay products include active neutrinos, the resulting flux can be searched for at neutrino telescopes such as IceCube or KM3NeT. Such a signal was considered in the model proposed in Ref. [2] by the author.

4.3 Thermal Freeze-out and WIMPs

Having discussed the general properties of a viable DM candidate, the next natural question concerns the origin and dynamics of such massive species. To address this, let us consider a DM particle, denoted by χ . We will also revisit the concepts of thermal decoupling and out-of-equilibrium dynamics introduced in the previous chapter.

In the early universe, SM particles coexisted in an extremely hot and dense plasma, the so-called thermal bath. At high temperatures, particle interactions occur frequently enough to maintain thermal equilibrium, characterized by $\Gamma > H$, where Γ is the interaction rate and H is the Hubble expansion rate. If a massive species were to remain in thermal equilibrium until the present day, its number density would be exponentially suppressed as $n_\chi \propto e^{-m_\chi/T}$, making it cosmologically irrelevant. Therefore, to obtain a significant relic abundance as observed today, the species must depart from equilibrium—a process known as *freeze-out* [257].

Freeze-out occurs when $\Gamma \simeq H$ at a temperature $T_{\text{fo}} < m_\chi$ (for the non-relativistic case), called the freeze-out temperature. At this point, the interactions are no longer efficient enough to maintain equilibrium. Thus, the particle χ decouples from the thermal plasma and eventually its comoving number density becomes effectively constant. This mechanism, known as **thermal freeze-out**, provides a compelling explanation for the origin of stable relic species in the early universe.

Consider the case of a stable DM, where the dominant interactions are annihilations and the corresponding reverse (creation) processes

$$\chi\bar{\chi} \leftrightarrow X\bar{X}, \quad (4.3)$$

where X denotes any SM (or beyond) particles into which the DM can annihilate. These interactions determine the comoving number density of DM.⁵ The relic abundance of

⁴ In particular, the DM lifetime should be much larger than the age of the universe if its decay products involve SM particles, for example, $\tau > 10^{23}$ s for a GeV-scale DM decaying into neutrinos [254–256].

⁵ We assume χ and $\bar{\chi}$ are equally abundant, with no particle-antiparticle asymmetry.

such a species can be computed by solving the Boltzmann equation (see Appendix C)

$$\frac{dY_\chi}{dx} = -\frac{s}{Hx} \langle \sigma v \rangle (Y_\chi^2 - Y_\chi^{\text{eq}2}), \quad (4.4)$$

where $Y_\chi = n_\chi/s$ is the comoving yield, $x \equiv m_\chi/T$, and $Y_\chi^{\text{eq}} = 0.145(g_*/g_{*,s})x^{3/2}e^{-x}$ is the equilibrium yield. The thermally averaged pair-annihilation cross-section is given by [258]

$$\langle \sigma v \rangle(T) = \frac{1}{8\pi m_\chi^4 T K_2(m_\chi/T)} \int_{4m_\chi^2}^{\infty} ds \sigma(s)(s - 4m_\chi^2) \sqrt{s} K_1\left(\frac{\sqrt{s}}{T}\right), \quad (4.5)$$

where s is the total centre of mass energy squared (not to be confused with entropy).

The distinctive features of freeze-out can be inferred from Fig. 4.4, where we numerically solve the BE above for a benchmark value. As long as $\Gamma \equiv n \langle \sigma v \rangle > H$ —where n is the number density of incoming particles—interactions are fast enough to maintain thermal equilibrium, and we have $Y_\chi \approx Y_\chi^{\text{eq}}$. Creation and annihilation processes occur at equal rates in this regime.⁶

When the temperature falls below the DM mass ($T < m_\chi$), the creation processes become exponentially suppressed, as the X particles no longer have sufficient energy to produce χ pairs. Eventually, at a temperature T_{fo} (or $x = x_f$), when $\Gamma \simeq H$, the annihilation process also freezes out. Beyond this point, Y_χ deviates from its equilibrium value and becomes approximately constant.

In the post-freeze-out regime ($\Gamma < H$ or $x > x_f$), the interaction rate is too slow to affect the number density significantly. This can be understood in terms of time scales: the reaction time $\tau \sim 1/\Gamma$ becomes larger than the age of the universe at that time $t_H \sim 1/H$. Thus, $\tau > t_H$ implies that χ and $\bar{\chi}$ annihilate too infrequently to impact their abundance. From this point onward, the number density continues to decrease only due to the expansion of the universe, not due to interactions.

4.3.1 Relic abundance

The DM abundance can be simply approximated as $Y_\chi \simeq Y_\chi^{\text{eq}}$ for $x < x_f$, and $Y_\chi \simeq Y_\chi^{\text{eq}}(x_f)$ for $x \geq x_f$. In the case of hot relics, where the species i is relativistic at freeze-out, the asymptotic yield $Y_i^\infty = Y_i^{\text{eq}}$ becomes nearly independent of x . This applies to neutrinos, whose equilibrium yield is given by

$$Y_\nu^{\text{eq}} \simeq 0.278 \frac{g_{\text{eff}}}{g_{*,s}(T_{\text{dec}})}, \quad (4.6)$$

where g_{eff} is the total number of relativistic degrees of freedom. Neutrinos decouple at $T_{\text{dec}} \sim 2$ MeV, for which $g_* = g_{*,s} = 10.75$. For a 2-component, single-generation neu-

⁶ Note that $n \propto T^3$ in the relativistic regime, and $n \propto (mT)^{3/2}e^{-m/T}$ in the non-relativistic regime. As the universe cools, n and hence Γ decreases.

trino, we have $g_{\text{eff}} = 3/2$. Plugging in these values, we can estimate the relic abundance of hot relics

$$\Omega_\nu h^2 = \frac{\sum_i m_i Y_\nu^\infty s_0 h^2}{\rho_c} \simeq \frac{\sum_i m_i}{91.5 \text{ eV}}, \quad (4.7)$$

where s_0 is the present-day entropy density and ρ_c is the critical energy density of the universe. As discussed earlier, hot relics like neutrinos cannot constitute the dominant component of DM. For example, using the Planck bound $\sum_i m_i < 0.12 \text{ eV}$, the above equation implies that SM neutrinos contribute only about 0.0013% to the total DM relic abundance.

For cold relics, freeze-out occurs at $x_f \gtrsim 3$, and the computation of the asymptotic yield typically requires numerically solving the Boltzmann equation, see eq. (4.4). However, approximate analytical estimates can be obtained by assuming the annihilation cross-section takes the form $\langle \sigma v \rangle = \sigma_0 x^{-n}$, where $n = 0$ for s -wave annihilation and $n = 1$ for p -wave processes. The resulting relic abundance is approximately given by [158]

$$\Omega_\chi h^2 \approx 10^{-10} \frac{x_f^{n+1} (n+1)}{(g_{*,s}/g_*^{1/2}) \sigma_0}, \quad (4.8)$$

where $g_*, g_{*,s}$ are evaluated at $T = m_\chi/x_f$. Note the inverse dependence of the relic abundance on the annihilation cross-section. The larger the cross-section, the longer the particle tracks its equilibrium abundance, which consequently leads to a smaller frozen-out yield.

Finally, we comment on the relic abundance of warm relics. In general, the abundance depends on the specific production mechanism, which need not be freeze-out. For example, WDM particles can be produced through oscillations or decays. In such scenarios, the relic abundance must be computed numerically. If WDM is produced via freeze-out, the freeze-out temperature T_{fo} is higher than that of cold DM. Typically, WDM candidates in the keV mass range can account for the observed DM abundance.

4.3.2 The WIMP miracle

The most popular example of cold relics are Weakly Interacting Massive Particles (WIMPs). As the name suggests, these are particles with weak-scale masses, $\mathcal{O}(1 - 1000 \text{ GeV})$, that interact with the SM particles via electroweak-strength interactions, with cross-sections roughly of the order $\sigma \sim G_F^2 m_\chi^2$.⁷ These candidates were proposed in the 1970s, following the growing recognition of the DM problem as a serious issue in cosmology. We refer the reader to the historical account: How DM Came to Matter [259].

Depending on their mass and interaction strength, freeze-out for WIMPs typically occurs in the range $x_f \sim 20\text{--}25$. This value can be considered a *magic number* for non-

⁷ There exists a wide variety of WIMP models. In each case, the coupling G_F should be replaced by the relevant effective coupling, depending on the interaction specifics of the model.

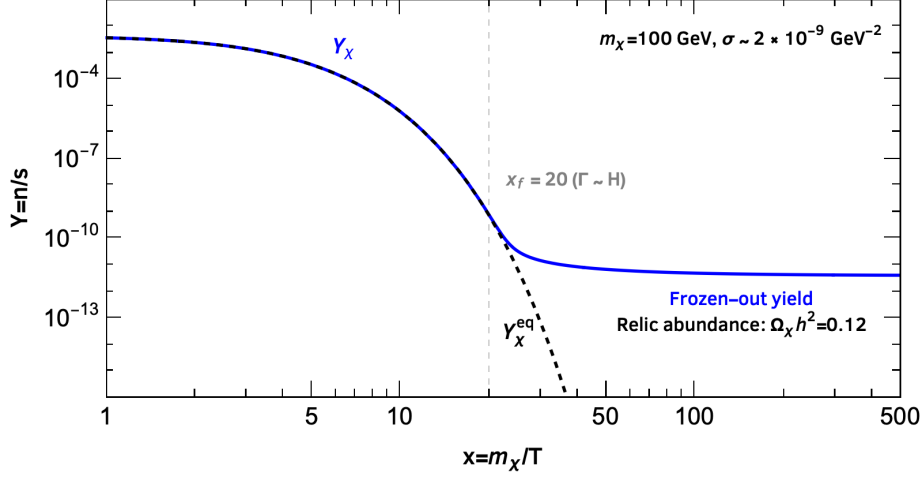


Figure 4.4: Freeze-out of a massive species (WIMP) with $m_\chi = 100 \text{ GeV}$ and $\sigma \sim 2 \times 10^{-9} \text{ GeV}^{-2}$. The black dashed curve represents the equilibrium yield, and the vertical dashed gray line denotes $x_f = 20$, the point at which freeze-out occurs.

relativistic thermal relics, marking the temperature at which their abundance decouples from the thermal bath.

Using $x_f = 20$, in eq. (4.8), we find that the observed relic abundance $\Omega_\chi h^2 = 0.12$ can be obtained for a cross-section in the weak range: $\sigma \sim 10^{-9} \text{ GeV}^{-2}$ (alternatively, $\langle \sigma v \rangle = 3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}$). This is known as the ‘WIMP miracle’. The numerical solution of the BE is shown in Fig. 4.4 for $m_\chi = 100 \text{ GeV}$ and $\sigma \sim 2 \times 10^{-9} \text{ GeV}^{-2}$. The asymptotic yield $Y^\infty = Y(x \gg x_f)$ can be related to the relic density in this case as

$$\Omega_\chi h^2 = 2.755 \times 10^8 m_\chi Y_\chi^\infty, \quad (4.9)$$

where the factor of 2.755×10^8 comes from the substituting the value of s_0, ρ_c in eq. (4.7). Substituting the value of $Y_\chi^\infty \sim 4 \times 10^{-12}$ and $m_\chi = 100 \text{ GeV}$, we get the correct relic abundance.

Given our understanding of the electroweak scale, this remarkable coincidence became a major motivation for considering WIMPs as viable DM candidates and for extensive experimental efforts to search for them (discussed below). Moreover, supersymmetric extensions of the SM (popular in those days) naturally predict particles that behave like WIMPs (for example, neutralino, gravitinos, right-handed sneutrino), further adding to their appeal in the past decades.

There are also theoretical bounds on how light or heavy these cold relics can be, approximately given by

$$\mathcal{O}(10 \text{ GeV}) \lesssim m_\chi \lesssim \mathcal{O}(100 \text{ TeV}). \quad (4.10)$$

The upper bound comes from the fact that the effective couplings cannot be arbitrarily large (the unitarity limit [260]), whereas, the lower bound can be derived from

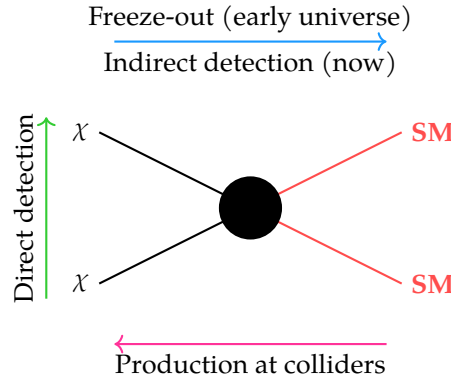


Figure 4.5: Illustrative representation of search strategies for WIMPs. The annihilations that lead to their freeze-out in the early universe now lead to indirect detection signals (blue arrow). The reverse process (creation) would lead to production at colliders (pink arrow). The green arrow denotes the χ -SM scattering that leads to direct detection signals.

not overclosing the universe for a weak scale cross-section (also known as the Lee-Weinberg limit [261]).

Before concluding our discussion on WIMPs, it is worth noting three important caveats where the simple relic abundance estimates discussed above break down [262]:

1. Resonances: When $m_A \simeq 2m_\chi$, where A is the particle mediating the annihilation, the cross-section can be significantly enhanced near the resonance.
2. Threshold effects: At high temperatures, the initial DM pair can have enough energy to produce heavier particles, which will result in $\langle\sigma v\rangle_{t_0} > \langle\sigma v\rangle_0$, the current abundance.⁸
3. Co-annihilation: When other particles (e.g., χ_2) are nearly degenerate with the DM ($m_{\chi_2} - m_{\chi_1} \lesssim T_{f_0}$), processes like $\chi_1\bar{\chi}_2 \leftrightarrow X\bar{X}$ become important and modify the relic abundance.

4.3.3 Experimental searches

We now focus on the experimental searches and strategies to detect WIMPs. The three main ways to look for WIMPs are: Direct Detection (DD), Indirect Detection (ID), and collider searches. A schematic illustration of these three approaches, based on the directionality of the annihilation processes $\chi\bar{\chi} \leftrightarrow X\bar{X}$, is shown in Fig. 4.5. Hence, the key assumption behind our current experimental searches for DM is that it interacts with the SM in some way. We discuss these briefly below.

Direct detection

The direct detection experiments are based on the scattering of the WIMP with the SM particle. As we are surrounded by the DM halo, the galactic WIMPs crossing the Earth can scatter off the atomic nuclei and deposit a fraction of its energy. This should lead

⁸ An interesting case of thresholds is that of forbidden DM, where DM particles annihilate into states heavier than themselves [263].

Type	Material	Experiments
Noble liquid	Liquid Xenon Liquid Argon	XENONnT [266], LZ [267], PandaX [268] DarkSide [269], DEAP-3600 [270]
Cryogenic crystals	Germanium Silicon	SuperCDMS [271], EDELWEISS [272] SuperCDMS-Si [271], DAMIC [273]
Crystal scintillators	NaI (Sodium Iodide) CaWO ₄ (Calcium Tungstate)	DAMA/LIBRA [274], COSINE-100 [275] CRESST [276]
Superheated liquids	C ₃ F ₈ , CF ₃ I	PICO [277]

Table 4.1: Direct detection experiments classified on the basis of target material.

to a small but detectable recoil of the atomic nuclei of the detector material. The idea of direct detection goes back to the 1980s, when Goodman and Witten proposed using the neutrino detectors to detect DM candidates [264].

As these signals arise from weak interactions, and are further subject to the DM distribution in the halo and velocity, they are expected to be very small, typically the nuclear recoil energy (E_R) is in the 0.1 – 50 keV range. Therefore, such signals must be distinguished from those induced by cosmic rays or radioactive decay products, and the experiment must be shielded. The DD experiments are usually placed in highly sensitive deep underground environment and use ultrapure materials. Depending on the detector material, the recoil energy is manifested in different ways, for example, via heat (phonon excitations of solid materials) or light (scintillating photons, direct ionization). The DD experiments are thus categorized on the basis of target material, they are summarized in Table 4.1. For a review of DM DD experiments, see Ref. [265]. The detection rate for a DM χ scattering off a nucleus (A, Z) is given as

$$\frac{dN}{dE_{NR}} = \epsilon \frac{\rho_\chi}{m_A m_\chi} \int_{v_{\min}} v f(v) \frac{d\sigma}{dE_{NR}} dv, \quad (4.11)$$

where N is the total number of detected DM particles, m_A is the atomic mass, and $\rho_\chi \approx 0.4 \text{ GeV}/\text{cm}^{-3}$ is the local DM density [278], σ is the DM-nucleon scattering cross-section, and $\epsilon = tM_T$ is the experimental exposure for observation time t and mass of the target, M_T . The velocity distribution is assumed to be isotropic and Gaussian.

Finally, the DM-nucleus cross-section is distinguished in two types: Spin-independent (SI) interactions that are independent of the DM spin and nucleus' angular momentum, and arise from scalar and vector couplings to quarks; Spin-dependent (SD) inter-

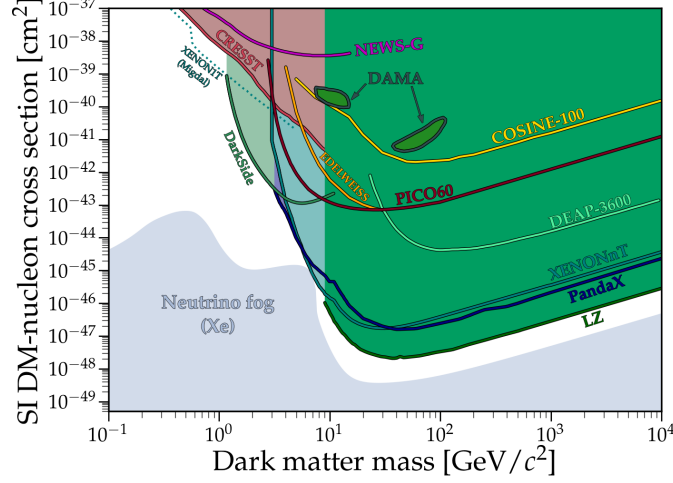


Figure 4.6: Current bounds on the spin independent DM-nucleon cross-sections from various DD experiments. The gray shaded region at the bottom denotes the neutrino fog: the irreducible neutrino background, where it becomes difficult to distinguish DM signals from that of neutrino induced scatterings. Credits: Ciaran O’ Hare.

actions that come from axial-vector couplings to quarks. They are given by

$$\begin{aligned}\sigma_{\text{SI}}(E) &= \sigma_p \frac{\mu_{\chi A}^2}{\mu_{n\chi}^2} \left(\frac{f_p Z + f_n (A - Z)}{f_p} \right)^2 F^2(E), \\ \sigma_{\text{SD}}(E) &= \frac{\sigma_p}{2J + 1} \frac{\mu_{\chi A}^2}{\mu_{n\chi}^2} \left(\frac{a_p^2 S_{pp}(E) + a_p a_n S_{pn}(E) + a_n^2 S_{nn}(E)}{a_p} \right)^2,\end{aligned}\quad (4.12)$$

where σ_p is the DM-proton cross-section; $f_{n,p}$, $a_{n,p}$ denote the couplings to neutrons and protons, respectively; J is the spin of the nucleus, $\mu_{\chi A,n}$ is the reduce mass of the DM-nucleus/nucleon system, and F, S are the nuclear form factors [279, 280]. The DD experiments thus place constraints on these scattering cross-sections and the corresponding DM mass. Over the years, the sensitivity of these experiments have increased manifolds. The null results from the searches at LZ, PandaX and XENONnT set severe constraints on the WIMP parameter space as shown in Fig. 4.6, ruling out the cross-sections consistent with the WIMP miracle as well as WIMP masses larger than 10 GeV. One can also see how close the bounds have come to the neutrino floor, the ultimate sensitivity limit where neutrino-induced scatterings can mimic the DM signals. While the neutrino fog is not a hard limit, the detection becomes more challenging in this region. Current efforts are directed at understanding the Coherent Elastic Neutrino-Nucleus Scattering (CEvNS) better [281], as well as at advanced techniques that can help us disentangle the DM and neutrino signals.

Indirect detection

The indirect detection aims at probing the annihilation (or decay) products of DM particles, as illustrated in Fig. 4.5. As WIMPs are supposed to interact with the SM particles, the annihilation products that one can look for could be photons, neutrinos,

electrons, positrons, protons, antiprotons or anti-deuterons. Since the annihilation rate scales with the DM number density ($\Gamma \propto n_\chi^2$), these annihilations can be significant in regions with high DM density. Therefore, key targets to look for such signals include dwarf spheroidal galaxies, galaxy clusters, the galactic centre or in dense stars or macroscopic bodies that can gravitationally capture the DM particles.

Concerning ID signals, there are telescopes searching for excess γ -rays such as Fermi-LAT [282–284], H.E.S.S [285, 286] and CTA [287]; neutrino telescopes and detectors such as IceCube [288], KM3NeT [289] and Super-K [67], as well as cosmic-ray detection based experiments such as AMS-02 that look for charged particle fluxes (discussed in the previous chapter) [138]. However, this detection task too is challenging due to the presence of significant astrophysical background/noise (for example, from pulsars or supernovae remnants). A smoking gun signal could be a γ -ray line or a neutrino line coming directly (undeflected) from a local source such as the galactic centre, dwarf galaxies or the Sun.

Like the DD case, we have not yet obtained any positive result from ID probes that can confirm a DM discovery. Note that there have been some interesting observations, such as the surplus of γ -rays coming from the Milky Way galactic centre detected by the Fermi telescope, which have been claimed as evidence for DM [290, 291]. However, the origin of this signal is widely debated within the theory community, without any consensus.

Collider searches

Finally, we comment on the production of DM particles from SM particles in high-energy collisions at collider-based experiments. Given the weak interactions and lifetime of the DM candidate, they are expected to interact very weakly with the detectors and are more likely to escape without a trace. This is characterized by missing transverse energy ($\cancel{E}_T$) in the final state.

As collider probes rely on detecting visible particles and their momenta, the production of an undetected particle leads to a momentum imbalance of the visible particles. The measurement of this recoil of visible particles against invisible ones is termed a mono-X signature. The ATLAS and CMS experiments search for DM particles based on these signatures, depending on the model and the mediator.

Collider probes are complementary to DD and ID experiments and are also more sensitive to the light DM mass range than DD probes. However, these searches have a limitation, as they cannot provide information about the stability of the DM particle or whether these particles can produce the correct relic abundance. Despite this, collider searches can provide interesting signals that can then be contrasted with DD and ID probes to lock in on the DM candidate.

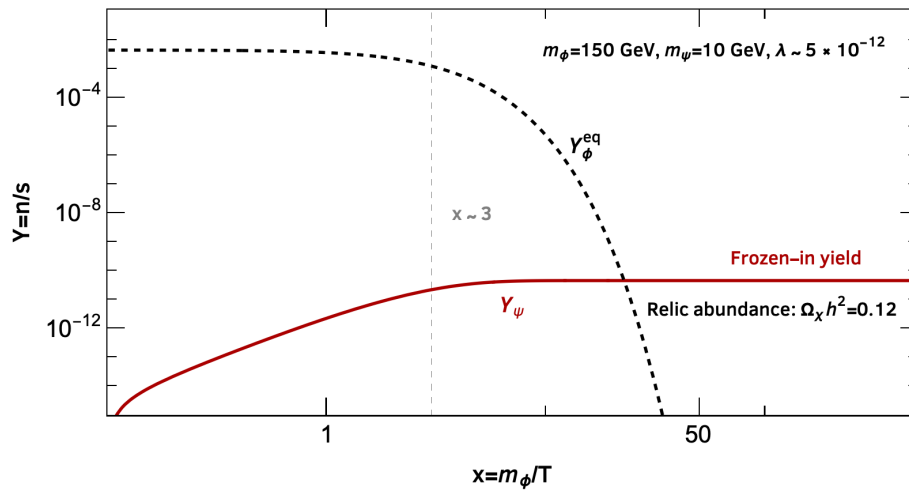


Figure 4.7: Freeze-in of a massive species (FIMP) produced from decays $\phi \rightarrow \psi\bar{\psi}$, with $m_\phi = 150$ GeV and $m_\psi = 10$ GeV. The black dashed curve denotes the equilibrium yield of ϕ . The vertical gray line denotes $x \sim 3$, the temperature at which production from decays is most efficient. The observed DM relic abundance is produced for a feeble coupling $\lambda \sim 5 \times 10^{-12}$.

4.4 Beyond WIMPs

While the freeze-out of WIMPs has long been a leading explanation for the observed DM relic abundance, the experimental searches have yielded null results (for a review, see Ref. [292]). This has motivated the exploration of alternatives—both beyond the traditional WIMP paradigm and beyond the standard freeze-out mechanism. Several such scenarios have been proposed, including those presented in the author’s works [1, 2, 4], which we now review.

4.4.1 Thermal Freeze-in and FIMPs

Instead of starting with a thermal abundance like in the case of WIMPs, consider a scenario in which the DM candidate, say ψ , has a vanishing initial abundance. This can be justified if ψ interacts with the SM only through extremely weak–*feeble* interactions, such that it never reaches thermal equilibrium with the thermal bath. Despite these suppressed interactions, a significant abundance of ψ can still be produced.

For example, ψ particles can be generated via the decay of a heavier particle ϕ that is part of the thermal bath: $\phi \rightarrow \psi\bar{\psi}$. A small population of ψ builds up from these decays until the abundance of ϕ becomes Boltzmann suppressed. (A similar argument holds for production through scatterings.) After this point, the production effectively ceases, and the abundance of ψ freezes in. This is known as the **thermal freeze-in** mechanism, and the DM candidate ψ is referred to as a **Feebly Interacting Massive Particle (FIMP)** [293].

In contrast to freeze-out, where the final abundance is inversely proportional to the interaction strength, in the freeze-in scenario, the relic abundance is directly proportional to the (feeble) coupling. The calculation of the relic abundance is significantly

simplified in the case of FIMPs, as processes involving them in the initial state can be safely neglected due to their negligible initial abundance and feeble interactions.

Consider a model where the FIMP candidate χ couples very weakly to a heavier scalar ϕ (with $m_\phi > 2m_\chi$) through an interaction term of the form $\lambda\phi\chi\bar{\chi}$, leading to its production via decays.⁹ In this scenario, the BE simplifies considerably and takes the following form:

$$\frac{dY_\psi}{dx} = \frac{1}{sHx} \gamma_\phi \left(\frac{Y_\phi}{Y_\phi^{\text{eq}}} \right), \quad (4.13)$$

where $x \equiv m_\phi/T$, and $\gamma_\phi \propto \lambda^2$ is the decay rate density (see eq. (3.25)) related to the decay width $\Gamma_\phi = \frac{\lambda^2}{8\pi} m_\phi$.

Note that in the simpler case of ϕ being in equilibrium, $Y_\phi = Y_\phi^{\text{eq}}$. Hence, the factor in the round brackets becomes one. The relic abundance can be analytically estimated in this case as [293]

$$\Omega_\psi h^2 \approx 10^{27} \frac{g_\phi}{g_{*,s}/g_*^{1/2}} \frac{m_\chi \Gamma_\phi}{m_\phi^2}. \quad (4.14)$$

As Γ_ϕ scales as λ^2 , a larger relic abundance is produced from a larger coupling strength. We show the numerical solution of the BE above in Fig. 4.7 for a benchmark value. It can be seen that for weak-scale masses, the correct relic abundance is typically produced for extremely small couplings, $\lambda \sim 10^{-12}$.

Due to the extremely feeble interactions of FIMPs, detecting them is very challenging. As a result, the freeze-in scenario remains relatively unconstrained experimentally and offers considerable flexibility for DM model building. The freeze-in mechanism can accommodate a wide variety of DM candidates (as discussed in the ‘DM Zoo’ section) and can also be combined with the freeze-out mechanism to yield various possibilities—where the FIMP may or may not be the actual DM candidate. Four such scenarios are illustrated in Fig. 4.8 [293], where LOSP stands for Lightest Observable Sector Particle, i.e., a particle that is present in the thermal bath and interacts feebly with FIMP. In Ref. [2], we utilized the freeze-in mechanism of scenario 2 to produce a FIMP DM from the freeze-out of another particle in an extended dark sector (see below). A review of FIMP models can be found in Ref. [294].

The example discussed earlier involving decays of weak-scale particles corresponds to what is known as **IR freeze-in**. In this case, the DM relic density is independent of any high-scale cutoff or UV completion and instead depends only on the masses and couplings of the particles involved. Conversely, if the interactions responsible for freeze-in arise from non-renormalizable operators, the dominant production occurs at very high

⁹ We consider a FIMP DM candidate with the same interaction in our model for theogenesis of DM, baryon asymmetry, and neutrino masses [9].

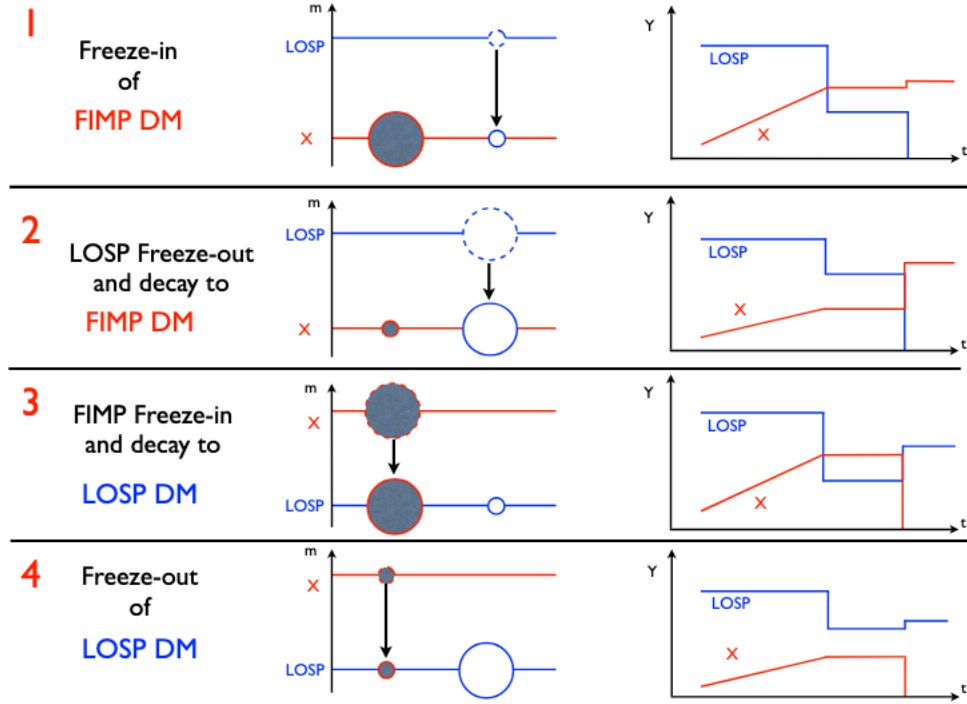


Figure 4.8: Scenarios utilizing the freeze-in and freeze-out mechanism for the production of FIMP (red) and LOSP (blue) DM, respectively. On the left side, the large (small) circle sizes represent the dominant (sub-dominant) production mechanism. The dashed (solid) circle boundary signifies that the particle is unstable (stable). The abundance evolution of LOSP and FIMP as a function of time is shown on the right side. Adapted from Ref. [293].

temperatures—either near the cutoff scale Λ or close to the reheating temperature T_{rh} . In such cases, the relic abundance becomes sensitive to both Λ and T_{rh} , a scenario referred to as **UV freeze-in** [295–297]. Our works in Refs. [2, 9] utilized the IR freeze-in mechanism.

4.4.2 Rich Dark Sectors

So far, our discussion has centred around scenarios where the total DM abundance arises from a single DM particle or component. However, this contrasts with the visible sector, which consists of multiple stable species—such as protons, electrons, neutrinos, and photons—stabilized by various symmetries. Given that the observed DM abundance is roughly five times larger than that of visible matter, it is natural to ask: Why should the dark sector consist of only one component? This question motivates the idea of a rich dark sector.

Rich dark sector models extend the SM by introducing a hidden sector composed of its own particles and interactions. These may include dark fermions, dark scalars, or even dark photons. Such sectors can give rise to a variety of phenomena, including strong self-interactions (which can help address small-scale structure issues in Λ CDM), dark phase transitions, or dark radiation, among others.¹⁰ Our works in Refs. [1, 2, 4] are all

¹⁰This hidden sector could also be an exact copy of the SM sector, termed as *Mirror DM*. However, this

examples of models featuring a rich dark sector.

Importantly, particles in the dark sector are not charged under the SM gauge groups. Instead, they interact via their own dark gauge groups or symmetries, making them secluded from the SM. However, they may still couple weakly to SM states through portal interactions mediated by connector particles (e.g., scalar, vector, or neutrino portals) [299]. This feeble connection provides a natural explanation for why WIMPs or other DM candidates have thus far evaded direct detection.

Portal interactions

The visible and dark sectors can be connected via weak or feeble interaction terms known as portals. These interactions act as bridges between the two sectors, allowing us to study DM phenomenology and guide our detection strategies. Portals are typically classified based on the type of mediator connecting the two sectors. The most common ones are:

- **Higgs portal:** Dark scalars connected to the SM Higgs. For example, a dark singlet real scalar ϕ can mix with the SM Higgs via a term like $\lambda |H|^2 \phi^2$ [300–302].¹¹ The scalar mixing can lead to the SM Higgs decay into dark sector particles, and the interaction strength can be constrained from the Higgs invisible decay width [305,306].
- **Vector portal:** Kinetic mixing between dark gauge bosons and SM gauge bosons. Consider a dark $U(1)'$ gauge boson Z' , it can mix with the SM Z bosons via a term like $\epsilon F'_{\mu\nu} F^{\mu\nu}$ [229,307]. This mixing (controlled by the value of ϵ) would then lead to Z' decays to the SM fermions.
- **Neutrino portal:** Sterile neutrinos (N) can serve as a link between the two sectors, for example, either they themselves are DM [308,309] (discussed later), or they can couple to a dark fermion ψ and a dark scalar ϕ via a Yukawa term $y_\psi N \psi \phi$ along with the usual $y_\nu N L H$ term. The neutrino portal to DM has been widely studied, see Refs. [310–314].

Despite these possibilities, the continued absence of experimental evidence raises the question: what if a rich dark sector exists but is highly or even completely decoupled from the SM? In Ref. [315], it was shown that it is possible to have thermal DM from highly decoupled dark sector, which evolves with its own temperature independently of the visible sector (we use this framework in Ref. [1]). In the extreme case of a completely decoupled dark sector, DM interacts with the SM only through gravity and lacks any portal interaction.

scenario faces tight constraints from astrophysics and cosmology, see Ref. [298] for a review.

¹¹The scalar singlet DM is a simple and minimal model, that gained popularity as a non-SUSY model. However, the parameter space of the original model is now mostly ruled out from experimental searches, recent status is presented in Ref. [303]. There are still a few workarounds to make it work [304].

Multi-component DM

In rich dark sectors, it is possible to have multiple stable particles contributing to the total DM relic abundance. This framework is known as multi-component DM [316,317]. Such scenarios typically involve more interactions, leading to a richer phenomenology and offer new ways to evade the constraints on WIMP models (see also Dynamical DM [318]). For instance, consider a dark-sector with n WIMP components, χ_i with $i = 1, 2, \dots, n$. In addition to the annihilation of eq. (4.3), we could also have: *i*) Conversions: $\chi_i \bar{\chi}_i \leftrightarrow \chi_j \bar{\chi}_j$ for $m_i > m_j$. *ii*) Semi-annihilations: $\chi_i \chi_j \leftrightarrow \chi_k X$ for $n > 2$ [319–322]. A general BE for a DM component χ_i including annihilation and conversions is given by

$$\frac{dY_i}{dx} = -\frac{s}{Hx} \left[\langle \sigma v \rangle_{i,\text{ann}} \left(Y_i^2 - Y_i^{\text{eq}2} \right) - \langle \sigma v \rangle_{ij,\text{con}} \left(Y_i^2 - \frac{Y_j^2}{Y_j^{\text{eq}2}} Y_i^{\text{eq}2} \right) \right], \quad (4.15)$$

where $x \equiv m_i/T$. We studied the implications of these conversion processes in our work in Ref. [4], and found that the contribution of each stable DM component to the total relic abundance follows a power-law dependence on its mass.

4.4.3 Asymmetric Dark Matter

Recall the $\mathcal{O}(1)$ relation between the observed DM and baryonic energy densities: $\Omega_{\text{DM}} \simeq 5 \Omega_B$. As we saw in the previous chapter, the observed baryon abundance arises from the BAU. On the other hand, we have just discussed that in the common case of WIMPs, the DM abundance is set by thermal freeze-out and depends on the annihilation cross-section. This leads to two immediate questions: Why do the two seemingly unrelated sectors yield similar abundances? And why are the mechanisms responsible for setting these abundances—an asymmetry in one case and freeze-out in the other—so different?

Taking inspiration from the visible sector, one may ask: what if the DM abundance is also set by an initial asymmetry in the dark sector? This is the central idea behind Asymmetric Dark Matter (ADM) frameworks [323]. In such models, the relic abundances are related as

$$\frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 \sim \frac{\eta_D m_{\text{DM}}}{\eta_B m_b}, \quad (4.16)$$

where η_D is the initial dark sector asymmetry, $\eta_B \sim \mathcal{O}(10^{-10})$ is the baryon asymmetry, and $m_b \sim 1$ GeV is the proton mass. In contrast to the freeze-out mechanism, the DM relic abundance in ADM scenarios is determined by the dark sector asymmetry and the DM mass, with no dependence on the annihilation cross-section. Assuming $\eta_D \simeq \eta_B$, the relation becomes predictive, suggesting $m_{\text{DM}} \sim 5$ GeV. This naturally explains the observed cosmic coincidence as arising from an $\mathcal{O}(1)$ ratio between the DM and proton masses—a key feature of the *ADM hypothesis*.

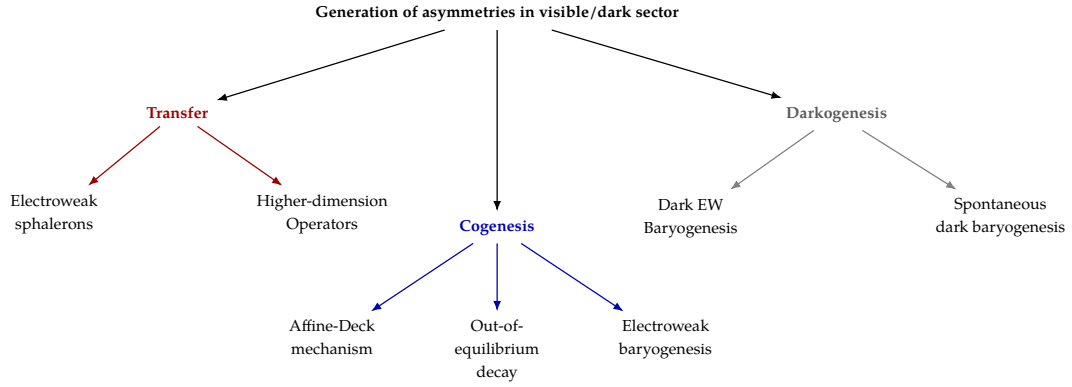


Figure 4.9: Mechanisms for generating an asymmetry in the visible and dark sectors. These have been described in Chapter 3.

This scenario typically favours DM masses lighter than those in the standard WIMP paradigm. While other values of η_D could allow for heavier or lighter DM candidates, it would not entirely be in the spirit of the ADM hypothesis. For comprehensive reviews on ADM, see Refs. [324,325].

Generic features

To begin with, an initial asymmetry must be generated in the dark sector. As discussed in the previous chapter, this asymmetry should be dynamically produced through a mechanism that satisfies the Sakharov conditions. This requires the existence of a conserved quantum number in the dark sector, denoted by D .¹²

To relate the visible and dark sectors in the ADM framework, a mechanism must establish a connection between their respective asymmetries. Broadly, this connection can arise in one of three distinct ways:

1. **Transfer:** The primordial asymmetry in the visible sector is communicated to the dark sector via some higher-dimension operators (for details, see Sec. 3 of the author's work in Ref. [4]).
2. **Cogenesis:** The asymmetries in both the sectors can be generated simultaneously, for example, via decays of a heavy particle connecting both the sectors. We use this approach for generating an asymmetry in Ref. [2].
3. **Darkogenesis:** The initial asymmetry is first generated in the dark sector and then communicated to the visible sector, for example via an FOPT in the dark sector or dark sphalerons.

Some examples of mechanisms to realize the generation of asymmetry in the two sectors are illustrated in Fig. 4.9.

Furthermore, the processes responsible for communicating the asymmetry between

¹²An alternative approach involves a generalized baryon number, wherein dark sector particles carry a baryon number equal in magnitude but opposite in sign to that of SM baryons. In such scenarios, the total baryon number of the universe is conserved.

the two sectors must eventually decouple at some temperature. This decoupling is essential to preserve the generated asymmetries in both sectors; otherwise, ongoing interactions could wash them out. In addition, the dark sector must contain interactions that efficiently deplete the symmetric component—just as in the visible sector, where the symmetric component annihilates into radiation. This ensures that the relic abundance is dominantly asymmetric.

To achieve sufficient annihilation, ADM models often introduce a light mediator or a dark force (typically present in rich dark sectors). For instance, consider a minimal scenario featuring a $U(1)_D$ gauge symmetry in the dark sector, containing fermions ψ and χ , along with a dark gauge boson Z' . In this setup, viable annihilation channels for the DM candidate ψ include: $\psi\bar{\psi} \rightarrow Z' \rightarrow \chi\bar{\chi}$ or $\psi\bar{\psi} \rightarrow Z'Z'$. If the Z' is massive, it can decay into SM fermions via kinetic mixing. If massless, it contributes to the radiation content as dark radiation.¹³

Asymmetric freeze-out

In Ref. [326], it was shown that DM can be *partially asymmetric*, meaning that a small symmetric component can survive after freeze-out. This framework—referred to as the *asymmetric WIMP* scenario—interpolates between the standard WIMP freeze-out mechanism and the completely asymmetric DM case. In this intermediate regime, the relic abundance depends exponentially on the annihilation cross-section, while retaining a linear dependence on both the initial dark-sector asymmetry η_D and the DM mass.

Consider a dark sector characterized by an initial asymmetry η_D , and a DM candidate ψ , such that the net asymmetry is $Y_\psi^+ - Y_\psi^- = \eta_D$. Due to this asymmetry, the two components ψ^+ and ψ^- will freeze out with different abundances, satisfying $Y_\psi^+ \gg Y_\psi^-$.

It is convenient to define the asymmetry ratio $r = Y_\psi^-/Y_\psi^+$, so that $r = 0$ corresponds to the case of completely asymmetric DM and $r = 1$ corresponds to the standard symmetric WIMP case. Therefore, we can express the yields as

$$Y_\psi^+ = \frac{\eta_D}{1-r}, \quad Y_\psi^- = \frac{r\eta_D}{1-r}. \quad (4.17)$$

We can write the corresponding BE following eq. (4.4), but now explicitly denoting $Y_\psi^\pm$ as

$$\frac{dY_\psi^\pm}{dx} = -\frac{s}{Hx} \langle \sigma v \rangle \left(Y_\psi^+ Y_\psi^- - Y_\psi^{+,eq} Y_\psi^{-,eq} \right). \quad (4.18)$$

This can be rewritten in terms of η and r to obtain

$$\frac{dr}{dx} = -\frac{s\eta_D}{Hx} \langle \sigma v \rangle \left[r - r^{eq} \zeta^2(r) \right], \quad (4.19)$$

¹³This model is similar to the one explored in our work in Ref. [1], although in that case, no dark-sector asymmetry was present.

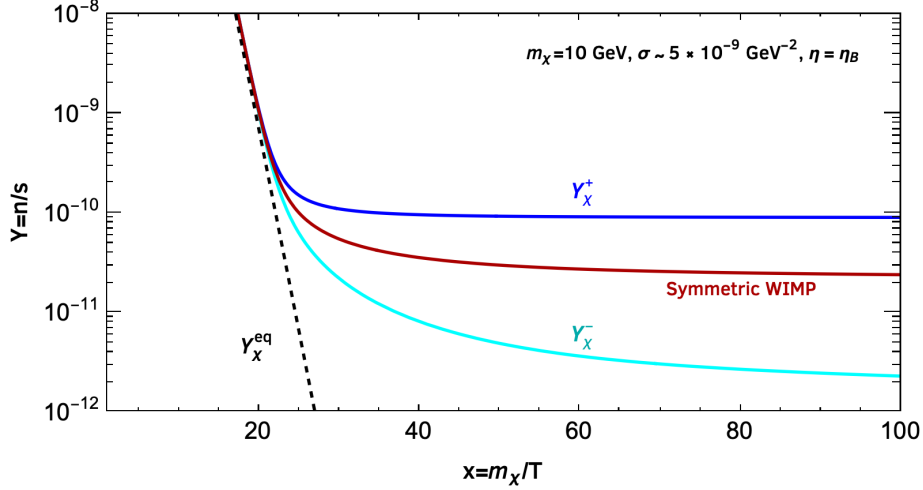


Figure 4.10: Asymmetric freeze-out of a massive species with $m_\chi = 10 \text{ GeV}$, $\sigma \sim 5 \times 10^{-9} \text{ GeV}^{-2}$ and $\eta = \eta_B$. The blue and cyan curves denote the abundance of Y_χ^+ and Y_χ^- , respectively. The red dashed curve denotes the freeze-out of a symmetric WIMP with the same mass and annihilation cross-section.

where $\zeta(r) \equiv 1 - r/1 - \bar{r}$, and $r^{\text{eq}} \equiv e^{-2\xi(x)}$ with $\xi = \sinh^{-1}(\eta/2Y_\psi^{\text{eq}})$. The BE above is solved numerically (see Fig. 4.10), however, a good analytical estimate is given by [326]

$$r(x) = r^{\text{eq}}(x_f) \exp[-\lambda \eta_D \Phi(x, m_\psi)], \quad (4.20)$$

where $\lambda = (\pi/45)^{1/2} M_{\text{Pl}} m_\psi \sigma_0$, and

$$\Phi(x, m) \equiv \int_{x_f}^x dx' x'^{m-2} g_*^{1/2}. \quad (4.21)$$

Here, we have once again used $\langle \sigma v \rangle = \sigma_0 x^{-n}$. The relic abundance is finally determined from the asymptotic solution r_∞

$$\rho_\psi = s m_\psi \eta_D (1 + \frac{2r_\infty}{1 - r_\infty}). \quad (4.22)$$

We recover the abundance for completely asymmetric DM by setting $r_\infty = 0 \Rightarrow \rho_\psi = s m_\psi \eta_D$. Further, the analytical solution above can be recast in a manner to compare with a typical WIMP cross-section ($\sigma_{0,\text{WIMP}} \sim \mathcal{O}(10^{-9}) \text{ GeV}^{-2}$) as

$$r_\infty \approx \exp \left[-2 \left(\frac{\sigma_0}{\sigma_{0,\text{WIMP}}} \right) \left(\frac{1 - r_\infty}{1 + r_\infty} \right) \right] \xrightarrow{r_\infty \ll 1} \exp [-2\sigma_0/\sigma_{0,\text{WIMP}}], \quad (4.23)$$

Therefore, in order to achieve $r < 1$, the exponential dependence leads to $\sigma_0 > \sigma_{0,\text{WIMP}}$. We extend this approach to multi-component dark sectors, to include several stable asymmetric WIMP candidates in Ref. [4], where in addition to the annihilations, we also include the conversions of the heavier DM components into lighter ones.

Having discussed the asymmetric freeze-out, we asked ourselves about the possibility of **asymmetric freeze-in** [327–329]. This was the main goal of our work in Ref. [2]

where we use the cogenesis path to generate an initial asymmetry in the visible and dark sectors (via two-sector thermal leptogenesis [330]). We then used the asymmetric freeze-out mechanism to have an asymmetric population of a fermion ψ in the dark sector. The ψ further decays to the FIMP like DM χ in the model. The DM χ turns out to be asymmetric if the decays of ψ peak once ψ^+ and ψ^- have frozen out with different abundances.

Before concluding this section, we briefly comment on the phenomenology of ADM. The DD signals in these models are expected to be similar to those of WIMPs and can be probed by the DD experiments discussed above. However, due to the asymmetric nature of the DM and the vanishing abundance of antiparticles, the ID signals are significantly suppressed. In the case of partially asymmetric DM, however, residual annihilations can still occur in the present epoch and may be constrained by CMB observations.

Interesting implications come from astrophysics and cosmology. The ADM particles can accumulate in compact objects such as neutron stars and white dwarfs, potentially reaching the critical mass for gravitational collapse. This phenomenon provides constraints on DM–nucleon interactions and serves as a complementary probe to direct detection experiments. Furthermore, the presence of dark radiation in ADM scenarios can impact BBN and CMB observations, for instance by contributing to an increase in the effective number of relativistic species, N_{eff} . For a detailed discussion on ADM phenomenology, see Refs. [324, 325].

4.4.4 The Dark Matter Zoo

The absence of any compelling evidence for particle DM has led to the formulation of a plethora of theoretical DM models and candidates over the years.¹⁴ Analogous to a real-life zoo, these DM models (and candidates) are thus clubbed under the term of ‘DM Zoo’. The foremost classification can be done on the basis of DM mass, as shown in Fig. 4.11. We have DM candidates ranging from 10^{-21} eV, all the way up to the Planck scale, about 50 orders of magnitude! Below, we discuss the categories on the basis of DM mass.

Heavy and super-heavy DM

The DM masses in the GeV–TeV range (corresponding to the standard WIMP mass scales) are categorized as intermediate or heavy DM. Most of our current experimental efforts are targeted at probing this range. The super-heavy range ($m_{\text{DM}} > \mathcal{O}(100 \text{ TeV})$) corresponds to super-heavy relics and macroscopic objects such as PBHs that have been discussed before. There are two super-heavy exotic DM candidates motivated by early universe physics: non-thermal DM candidates—WIMPZillas [336]—and Q-balls,

¹⁴There are several creative models: Cannibal DM [331], Selfish DM [332], Super-cool DM [333], Homeopathic DM [334], Pandemic DM [335], etc.

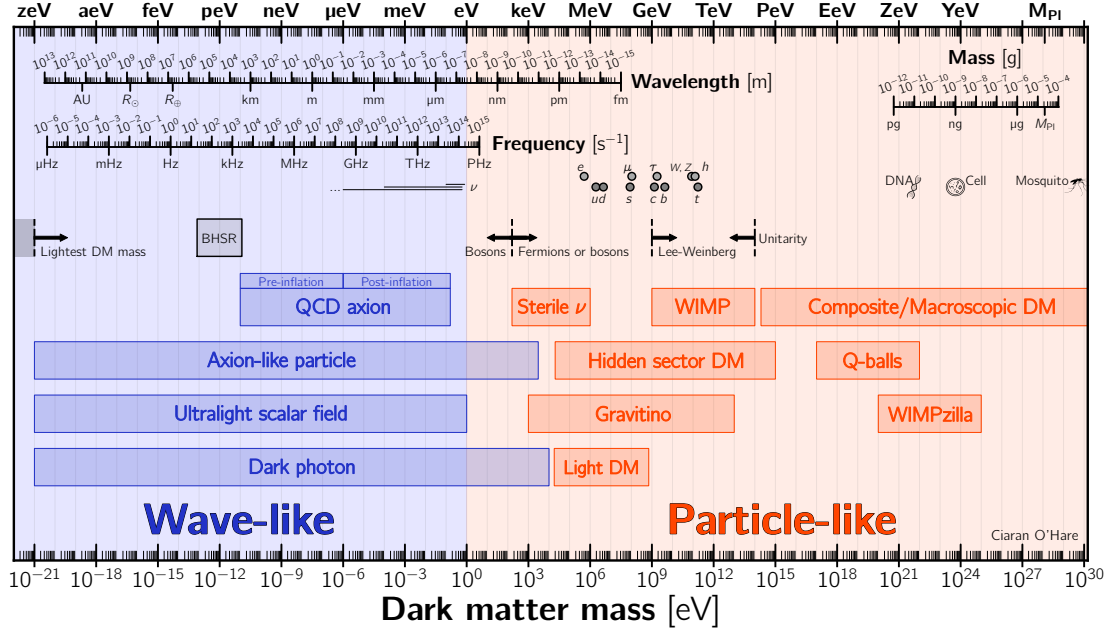


Figure 4.11: Different DM candidates categorized on the basis of mass. The blue region corresponds to wave-like DM (bosonic DM), while the red region corresponds to particle DM (fermionic or bosonic DM). The limits on the lightest fermion ($m > \text{keV}$) and boson mass ($m > 10^{-22} \text{ eV}$) are shown by the small dashed black line. Similarly, the bounds on WIMP masses discussed before are also depicted. Also shown are the corresponding sizes of SM particles and small entities for comparison with macroscopic DM like PBHs. Image courtesy: Ciaran O’Hare.

macroscopic scalar field configurations [337,338], both shown in Fig. 4.11. In general, super-heavy DM candidates are produced non-thermally, for example via gravitational production, and are mostly non-interacting or have extremely weak interactions.

Light DM and Sterile neutrinos

DM masses below the GeV scale, going down to the eV–keV scale, correspond to light DM candidates. There are several constraints on light DM scenarios from BBN, CMB, and structure formation (as the mass range indicates WDM candidates). These DM candidates can be probed through electron recoils, beam dump experiments, rare meson decays, or more recently via table-top condensed matter-based experiments [339].

The most common candidates in this range are **sterile neutrinos**, which can be produced via oscillations from SM neutrinos. This is known as the *Dodelson–Widrow (DW) mechanism* [308]. This mechanism involves a very simple extension of the SM by a sterile neutrino ν_4 , characterized by its mass ($\sim \mathcal{O}(\text{keV})$) and mixing (θ) with the active neutrinos, so that we have

$$\nu_4 = \cos \theta \nu_s + \sin \theta \nu_a, \quad (4.24)$$

where ν_s denotes the new interaction eigenstate and ν_a represents a combination of $\nu_{e,\mu,\tau}$.

Initially, the sterile neutrinos are not in equilibrium with the SM thermal bath, while the active ones are in equilibrium and abundant. The latter propagate freely; however, they develop a ν_s component due to the mixing. A small component of them thus oscillates into sterile neutrinos and remains thereafter, thus producing their relic abundance.

The DW mechanism can be easily accommodated in the Neutrino Minimal Standard Model (ν MSM), where the SM is extended by three RHNs as in the type-I seesaw [340, 341]. In these models, the lightest RHN is the sterile neutrino DM candidate, while the heavier RHNs lead to two non-zero neutrino masses. Moreover, the BAU can also be explained by leptogenesis, making it an elegant unified framework for addressing all three issues together.¹⁵

We accommodate a sterile neutrino DM residing in a dark sector in our model in Ref. [1]. The sterile neutrino relic abundance is determined by freeze-out in the dark sector at a temperature smaller than the visible one, and is therefore different from the mechanism discussed above.

Ultralight DM and Axions

These are DM candidates with super small masses (below the eV scale). For these mass scales, only bosonic DM candidates are allowed—ultralight fermions are mainly excluded due to the Pauli exclusion principle [346]. Furthermore, at these masses, the wave-like effects become important, as their corresponding deBroglie wavelength exceeds the inter-particle separation. Hence, they are termed as **wave-like** DM (see Ref. [347] for a review).

The QCD axions (mentioned in Chapter 1) are one of the most popular examples of ultralight DM. Axions have extremely feeble interactions with other SM particles, making their detection challenging, though several novel search strategies are currently being developed. More generally, even if the QCD axion is excluded as DM, there exists a class of models with **Axion-like Particles (ALPs)** that can arise in some extensions of the SM and behave as DM; see Refs. [348, 349].

¹⁵This minimal scenario is now mostly in tension (in the regions where ν_s makes up the total relic abundance) with X-ray and Lyman- α observations [342–345].

SUMMARY AND OUTLOOK

The introductory chapters serve as the background and motivation for the works carried out in this thesis. In this chapter, we summarise the motivation, objectives, and results of the research works. The research works on dark sectors and leptogenesis (Refs. [1–4]) by the author are presented in the following part of ‘Scientific Research’. On the other hand, the research works of Refs. [6–9] are not included in Part II, though their central ideas and results are summarised below.

5.1 Neutrino Sector

In Chapter 2, we discussed the neutrino mass mechanisms and the standard seesaw models. We also discussed the low-scale variants of type-I seesaw, along with a few radiative neutrino mass models. The testability of LNV in these high-scale standard seesaws remains an issue.

The scale can be parametrically suppressed if there are new scalars in the theory, which in turn could lead to new Weinberg-like operators and seesaws [6,7]. Moreover, these seesaws can also explain the baryon asymmetry via leptogenesis, as discussed in Chapter 3. We discussed the vanilla leptogenesis cases as well as going beyond the canonical approach. The work in Ref. [3] is one such approach that aims at lowering the scale of type-III leptogenesis.

New Weinberg-like operators and seesaw models

In Refs. [6,7], we investigate new Weinberg-like operators involving one or two additional $SU(2)$ scalar multiplets, extending up to the quintuplet representation. These operators can generate Majorana neutrino masses post EWSSB. Tree-level UV completions of these effective operators will lead to novel seesaw models. These scenarios are shown in Table 5.1.

As these new scalars acquire VEVs, they are expected to remain close to the electroweak

Model	Scalar Multiplets	Mediators	Wilson Coefficients
A₁	$\Phi_1 = 4_{-1/2}^S$	$\Sigma = 5_0^F$	$C_5^{(2)} = y_1 M_{\Sigma}^{-1} y_1^T$
A₂	$\Phi_1 = 4_{-3/2}^S$	$\mathcal{F} = 3_{-1}^F$	$C_5^{(1)} = y_1 M_{\mathcal{F}}^{-1} y_H^T + y_H M_{\mathcal{F}}^{-1} y_1^T$
B₁	$\Phi_1 = 4_{1/2}^S, \Phi_2 = 4_{-3/2}^S$	$\mathcal{F} = 5_{-1}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B₂	$\Phi_1 = 3_0^S, \Phi_2 = 5_{-1}^S$	$\mathcal{F} = 4_{-1/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B₃	$\Phi_1 = 5_{-2}^S, \Phi_2 = 5_1^S$	$\mathcal{F} = 4_{3/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B₄	$\Phi_1 = 5_{-1}^S, \Phi_2 = 5_0^S$	$\mathcal{F} = 4_{1/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$

Table 5.1: List of genuine models obtained from new Weinber-like operators with 1 new scalar multiplet Φ_1 (A_i) and with two new scalar multiplets $\Phi_{1,2}$ (A_i) [7]. The UV completions corresponding to Majorana fermions or vector-like fermions are shown in the mediator column. The Wilson coefficients in terms of the mediator mass and Yukawa couplings $y_{1,2,H}$ with the scalars $\Phi_{1,2}, H$ are shown in the last column. The multiplets are represented as $N_Y^{S,F}$ where N denotes the $SU(2)$ representation, S, F denote scalar and fermion nature of the multiplet, and Y its hypercharge.

scale. Moreover, the ρ parameter constrains these new VEVs to be $\sim O(1)$ GeV. We analyse the scalar potentials for naturally induced small VEVs and discuss their implications for pseudo-Goldstone bosons.

Additionally, we establish general constraints on the new scalar multiplets from loop corrections to EWPTs and through direct searches for doubly charged scalars at colliders (discussed in Section 2.5.3), see Fig. 5.1. The production cross-sections at $\sqrt{s} = 13$ TeV were determined using MadGraph [350].

Furthermore, we had explored the possibility of explaining the CDF measurement of the W -boson mass [351].¹

We computed the contribution to neutrino masses at both tree-level and one-loop level, and discussed low-energy variants either based on inverse seesaws or assuming hierarchies among the two different Yukawas. Depending on these hierarchies, it is possible to have a rich phenomenology from multiplets with the larger Yukawas. We also examined the generation of dimension-6 operators and explored their implications such as non-unitarity of the leptonic mixing matrix, non-universal Z -boson interactions, and LFV.

Additionally, in the presence of specific symmetries, we construct scotogenic or generalized scotogenic-like models that can not only produce viable DM candidates but also generate neutrino masses. To satisfy the DM relic abundance, we provide mass ranges for the new scalar and fermion multiplets. As these are seesaw models, we can also incorporate leptogenesis, however, EW baryogenesis may be realized directly due to the presence of new scalars.

¹ More recent W -mass measurements are consistent with earlier results, hence, the CDF anomaly could be erroneous. In any case, whether the anomaly disappears or exists, addressing it was not the main goal of our work.

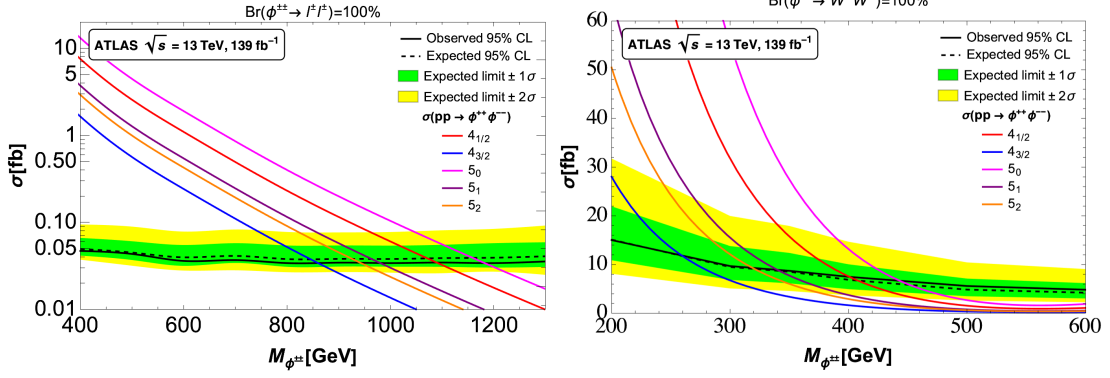


Figure 5.1: Upper limits on the production cross-section of doubly-charged scalars times their branching ratio into same-sign dileptons (left) and to a pair of W bosons (right). The region above the observed limit is excluded. The green and yellow bands denoted the expected exclusion curves within $1 - 2\sigma$ [136, 352].

Triplet leptogenesis at lower scales

We saw in Chapter 3 that the canonical type-I leptogenesis mechanism requires the fermion singlet masses $M_N > \mathcal{O}(10^9)$ GeV. However, for fermion triplets (type-III seesaw), this threshold increases to $M_\Sigma > \mathcal{O}(10^{10})$ GeV [168]. The heightened requirement for the triplets is due to the presence of gauge interactions, which compete with the Yukawa interactions, leading to a decrease in their population.

While several works demonstrate the plausibility of reducing the scale of type-I leptogenesis to $\mathcal{O}(\text{TeV})$ via an extended scalar sector, for example, in two-Higgs doublet models or scotogenic models, this possibility had remained unexplored for type-III leptogenesis. In Ref. [3], we investigate if it would be possible to lower down the scale of fermion triplet leptogenesis when it couples to a second Higgs doublet Φ_2 , which assumes a much smaller VEV compared to the SM one.

We show that it is indeed possible to bring down the scale to $\mathcal{O}(10^6)$ GeV, when flavour effects are included. This is attributed to the enhanced Yukawa interactions with the second doublet. We also discuss the compatibility with naturalness and constraints from strong washout. The results are summarized in Fig. 5.2. The region above the dashed line corresponds to the naturalness constraint, where the correction to the Φ_2 mass from M_Σ is larger than $\Lambda_{\text{relaxed}} = 10$ TeV. The blue shaded region is excluded from inefficient leptogenesis, whereas the orange shaded area represents the region where one can have potentially strong $\Delta L = 2$ washout.

This study though done in the context of type-III seesaw, can be extended for exploring leptogenesis in the new seesaw models discussed above because of common features: *i*) new multiplets have gauge interactions; *ii*) the new VEVs are comparatively smaller leading larger Yukawas.

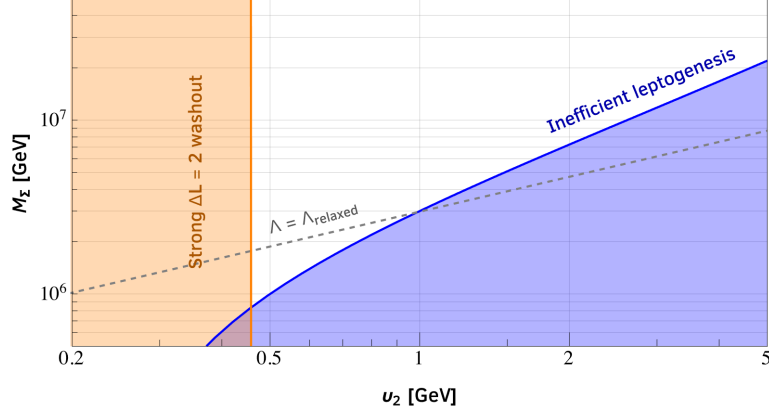


Figure 5.2: The viable parameter space for successful leptogenesis in a type-III seesaw model where the triplet fermions Σ couple to a second Higgs doublet Φ_2 with VEV v_2 . See Ref. [3] for more details.

5.2 Dark Sector

In Chapter 4, we discussed that DM constitutes roughly 25% of the total energy budget of the universe. Given this proximity of baryonic and DM densities, there's a plausible assumption that the dark sector could be quite rich, comprising multiple stable particles. Moreover, the DM relic abundance might be established by an initial asymmetry akin to the baryonic one. This idea of having multi-component and asymmetric DM served as the motivation behind the works in Refs. [2,4].

Multiple symmetric and asymmetric components

In Ref. [4], we consider an extended dark sector composed of several components, encompassing symmetric, partially asymmetric, and completely asymmetric DM candidates. In the presence of multiple stable components, one requires a combination of stabilising symmetries, which can guarantee at least n stable DM components. With just one discrete symmetry $\mathbb{Z}_N$, multi-component DM requires $N \geq 4$. It is more convenient to use combinations of both discrete and continuous symmetries to realize a multi-component framework. We show that for a product of n discrete symmetries $\mathbb{Z}_N$, at least n stable DM components are allowed, but can accommodate more stable components depending on the kinematics. We then discuss the correlation between the asymmetries in visible and dark sectors, depending on the conserved charges in each sector and the higher-dimensional operator that transfers the asymmetries.

Further, as discussed, this setup with more interactions than the conventional WIMP approach, leads to interesting phenomenological signatures. We investigate the role of conversions among particles in the dark sector, showcasing how these conversions tend to alter the abundance of the heavier and lighter components compared to scenarios with only annihilations. We also explored the role of conversions in the presence of initial asymmetries in the dark sector for the first time. Our findings show that conversions tend to make heavier DM components highly asymmetric, conversely, lighter

Model	σ	Symmetric		Partially Symmetric
		Ann.	Ann. + Conv.	Ann. / Ann. + Conv.
Mass-independent	$\neq f(m_i)$	$\sim m^0$	$\sim m^{-1}$	$\sim \eta m^{[1,0]}$
Heavy mediators	$\sim m_i^2/\Lambda^4$	$\sim m^{-2}$	$\sim m^{-[3,2]}$	$\sim \eta m^{[1,-2]}$
Light mediators	$\sim 1/m_i^2$	$\sim m^2$	$\sim m^{[0,2]}$	$\sim \eta m^{[1,2]}$

Table 5.2: Scaling of the relic abundance with the DM mass for all but the lightest component in a multi-component DM, where η denotes the dark sector asymmetry. The left/right exponents in the square brackets correspond to strong/weak conversions (and annihilations) in the symmetric (partially-asymmetric) case.

components exhibit predominantly symmetric behaviour.

We also conduct an effective parameterization considering heavy/light mediators to determine the power-law scaling for DM yield based on its mass. This framework should be applicable to a wide range of models, provided the masses/cross-sections lie at the weak scale, see Table 5.2.

It can be seen that in several cases $\Omega \sim m^0$. In these cases, the DM components in presence of conversions can have similar contributions to the total relic abundance. The effect of conversions in both symmetric and asymmetric components shows the abundances can be reduced, and the dark sector could remain hidden from the current experimental reach.

Cogenesis and asymmetric freeze-in

In Ref. [2] introduces a cogenesis model (see Section 4.4.3) to simultaneously generate an asymmetry in the visible and dark sectors. This model expands upon the two-sector thermal leptogenesis mechanism by augmenting the dark sector to encompass two distinct DM components: a scalar S and a fermion ψ . In various scenarios within the model, either one or both of the DM components exhibit an asymmetry.

The motivation was to realize a mechanism of asymmetric freeze-in. To this end, we proposed a feebly interacting particle serving as a potential asymmetric DM candidate produced in the decays of a TeV scale fermion χ that is an asymmetric WIMP. However, whether the decay products also exhibit an asymmetry depends on the strength of the freeze-in coupling. This is illustrated in Fig. 5.3. If the dominant production from decays takes place before χ has frozen out (for larger values of the feeble coupling), then the resulting DM population (ψ) is symmetric. On the other hand, if the feeble coupling is comparatively smaller, then the dominant production takes place from late decays, once χ has asymmetrically frozen out. In the latter case, the decay products are also asymmetric.

Hence, the late decays of the parent particle χ play a pivotal role in determining the asymptotic nature and yield of the DM component. These late decays can also significantly populate the symmetric component at later times, effectively restoring annihilation

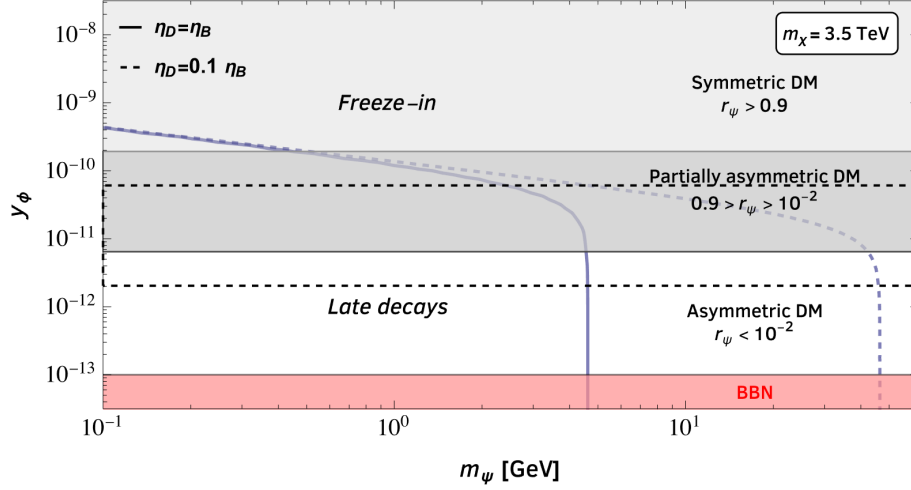


Figure 5.3: The parameter space to realize asymmetric freeze-in. The FIMP DM ψ is produced in the decays of χ via feeble Yukawa interactions y_ϕ . In the gray (dark gray) shaded region for larger values of y_ϕ , we find that the DM ψ produced via freeze-in is symmetric (partially asymmetric) as denoted by the asymmetry ratio r . In the white region for smaller values of y_ϕ , the ψ is produced dominantly in late decays of χ and turns out to be asymmetric. The pink region represents the constraints from BBN on the lifetime of χ [2].

lations and enhancing signals in indirect detection. Notably, in this scenario, the DM abundance is governed by the initial asymmetry, despite being asymptotically symmetric in nature. Hence, we have $O(1 - 10)$ GeV DM masses, consistent with the ADM hypothesis.

Furthermore, our model presents an intriguing phenomenological signal in the form of a monochromatic flux of neutrinos originating from the late decays of the heavier of the two DM components, potentially detectable at neutrino telescopes. While the main model is based on an extension of type-I seesaw and difficult to probe, we also provide an inverse seesaw model that can work at low scales and successfully realize the mechanism.

Light dark sectors and Warm DM

In Chapter 2, we discussed the tension between the neutrino mass bounds from cosmology and oscillation experiments. Further, these cosmology bounds take at face value would imply that it would be difficult to detect neutrino masses in terrestrial experiments such as KATRIN.

Since cosmological mass bound are sensitive only to the energy density of neutrino and not directly to their masses, the neutrino mass bounds can be relaxed if there is a mechanism to suppress their number density. In Ref. [1], we present a mechanism to relax the cosmological mass bound by introducing a light dark sector with several massless fermions (dark radiation) and a sterile neutrino DM candidate in the $O(10 - 100)$ keV mass range.

The dark sector comes into thermal equilibrium with the active neutrinos after neu-

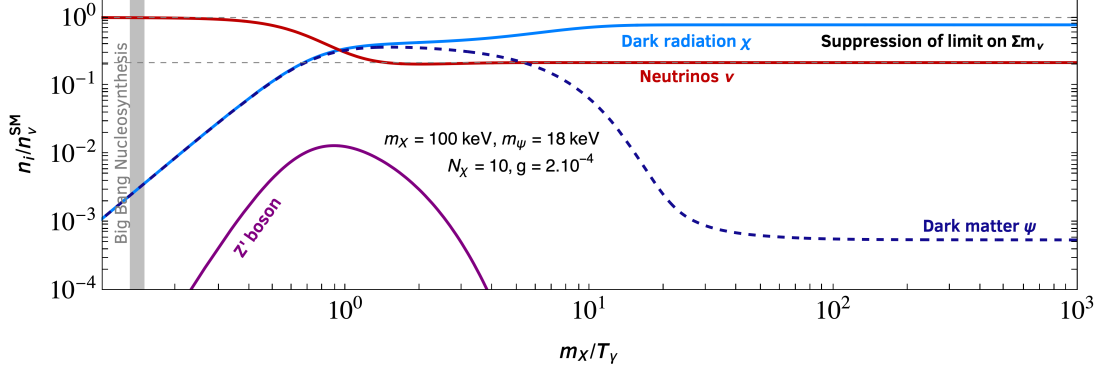


Figure 5.4: Schematic representation of the mechanism to reduce neutrino number density and relax the neutrino mass bound in Ref. [1]. The lower gray dashed line corresponds to $n_\nu/n_\nu^{SM} \simeq 0.2$.

trino decoupling and before recombination in the temperature range $O(100 \text{ keV}) \gtrsim T \gtrsim O(10 \text{ eV})$ via a $U(1)$ gauge interaction in the dark sector. The light dark sector states including the DM ψ , dark radiation χ , and the dark gauge boson Z' are all populated at the expense of active neutrinos via these new gauge interactions. Note that as the neutrinos have already decoupled, they are not produced any more from the SM particles. Hence, the mechanism suppresses the abundance of active neutrinos and therefore reconciles sizeable neutrino masses with cosmology.

The schematic representation of the mechanism obtained via numerically solving the relevant BEs is shown in Fig. 5.4. After neutrinos decouple, they come into equilibrium with the dark sector particles: Z' , DM ψ and N_χ generations of massless fermions χ . The interactions $\nu\nu \rightarrow Z'$ reduce the neutrino population, while interactions such as $Z' \rightarrow \chi\chi$ and $Z' \rightarrow \psi\psi$ (depending on the hierarchy) populate the fermions in the dark sector. The N_χ generations of χ contribute as dark radiation and compensate for the decrease in N_{eff} due to reduction in neutrino number density. The DM relic abundance is determined once the annihilations of the sterile neutrino DM ψ such as $\psi\psi \rightarrow \chi\chi$ freeze out in the dark sector (see the distinctive freeze-out curve).

We embed the mechanism of Ref. [353] in a UV complete model based on the minimally extended type-I seesaw [354]. This model is similar to our model in Ref. [2], however, it includes a richer dark sector and features no dark asymmetry. We explore the phenomenology of the WDM candidate in detail, including constraints from stability and structure formation.

The cosmological neutrino mass bound can be relaxed by a factor of ~ 5 (8) for 10 (20) generations of massless dark fermions. One of the main signatures of this model is slightly increased N_{eff} at CMB than at BBN, as well as suppression of matter power spectrum at smaller scales. Furthermore, the model is compatible with successful vanilla leptogenesis, though this was not the main goal of the work.

5.3 Unified frameworks

We saw that the type-I seesaw can be extended to include a DM candidate. These unified frameworks are based on the motivation of finding a common origin of the three key BSM issues that this thesis focuses on. This common origin was the motivation behind the works carried out in Refs. [8,9].²

A continuous symmetry argument

In Ref. [9], we propose a model featuring a heavy scalar mediator with non-zero lepton number (L) that has lepton number conserving and lepton number violating interactions to the visible and dark sectors, respectively. The scalar decays to the heavy right-handed neutrinos (charged under $U(1)_L$) through substantial interactions (leading to type-I seesaw and leptogenesis) and to dark fermions (uncharged under $U(1)_L$) via feeble interactions, thus giving rise to a FIMP-like DM. The heaviness and smallness of the parameters characterizing the new physics are naturally ascribed to the subsequent conservation and violation of this lepton number symmetry, respectively.

A discrete symmetry argument

In Ref. [8], we present another model based on a discrete $\mathbb{Z}_4$ symmetry, where a heavy scalar once again connects the neutrino sector with a feebly-interacting dark sector of the universe. The $\mathbb{Z}_4$ symmetry breaks down to a remnant $\mathbb{Z}_2$ when the scalar acquires a VEV, thereby generating tiny neutrino masses and explaining the observed baryon asymmetry through leptogenesis. This model introduces new contributions to CP asymmetry from scatterings. The scale of $\mathbb{Z}_4$ breaking occurs at significantly higher energy scales, potentially leading to the formation of domain walls (DW). These defects could dominate the universe's energy density, resulting in a cosmological catastrophe unless they are rendered unstable. Consequently, if the discrete symmetry is only approximate, the collision and annihilation of these DWs can generate a stochastic gravitational wave (GW) background. Such a signal could be detected by current and future GW detectors, offering a window to explore not only high-scale seesaw models but also to probe the feebly interacting dark sector.

5.4 Final remarks

The origin of neutrino masses, the explanation for the observed baryon asymmetry of the universe, and the nature of DM are among the clearest indications of BSM physics. While minimality has long served as a guiding principle in building new models, we have seen that many minimal and elegant solutions are either severely constrained or extremely difficult to probe. This necessitates efforts to construct new models that

² Note that the models in Refs. [1,2] are all also unified frameworks, though they were primarily focused on the DM phenomenology, and therefore have been discussed above.

can evade existing constraints and yield predictions testable via astrophysical, cosmological, or collider observables. Model-building and phenomenology have thus been central to the research works presented in this thesis to go beyond the SM.

In the light (or perhaps darkness!) of null results from various experiments in the search for new physics, it becomes reasonable to infer that new physics may be feebly coupled to the SM or may reside at high energy scales. A quick glance at the titles of the included articles immediately reveals a focus on BSM extensions that go beyond minimal frameworks. This approach stands in contrast to the principle of Occam's razor (presented at the beginning of this thesis). Nevertheless, most of the works presented here tackle these issues within unified frameworks. Though these frameworks are not strictly minimal, they are interesting in that they establish connections between neutrino physics and the dark sector, while retaining the elegance of the minimal extension upon which they were built upon.

In models with rich dark sectors, we see how new physics scenarios can remain hidden or evade detection due to their feeble coupling to the SM particles. Nevertheless, they may leave observable imprints on astrophysical and cosmological data. Turning to the early universe as a 'cosmic laboratory' thus becomes an important strategy for testing their viability. In contrast, the models discussed in the neutrino sector were motivated by their low mass scales and testable predictions, which can be probed by current or near-future experiments. In this way, the thesis adopts two complementary approaches to BSM phenomenology: i) Construct testable models that can be constrained or ruled out by experiments; ii) Develop models that explain why new physics has so far remained elusive, hinting at a deeper underlying theory.

Looking back at the SM, it took about 115 years from the first discovery of the electron in 1897 to the last discovery of the Higgs boson in 2012. Over these years, our understanding of the universe has expanded significantly, thanks to the extensive experimental efforts of the past few decades. Yet, much remains unknown. We are now entering the precision era of neutrino physics, with numerous experiments either underway or in development. In the coming decade, we hope to obtain answers to the open questions in the neutrino sector. While these insights will also help us explore possible connections between neutrinos and the dark sector, DM searches themselves are becoming increasingly sensitive. With upcoming space-based missions and novel strategies to detect light and ultralight DM, we may finally be able to shed some light on the dark!

Despite the rough patch the field is currently facing, well-motivated, interdisciplinary BSM frameworks will lead the way towards a new SM. While we may not know how long it will take for this 'New SM' to emerge, ongoing model-building efforts, constraints, and null results will continue to bring us closer to unravelling the secrets of our universe.

Part II: Scientific Research

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Large neutrino mass in cosmology and keV sterile neutrino dark matter from a dark sector

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ABSTRACT: We consider an extended seesaw model which generates active neutrino masses via the usual type-I seesaw and leads to a large number of massless fermions as well as a sterile neutrino dark matter (DM) candidate in the $\mathcal{O}(10\text{--}100)$ keV mass range. The dark sector comes into thermal equilibrium with Standard Model neutrinos after neutrino decoupling and before recombination via a $U(1)$ gauge interaction in the dark sector. This suppresses the abundance of active neutrinos and therefore reconciles sizeable neutrino masses with cosmology. The DM abundance is determined by freeze-out in the dark sector, which allows avoiding bounds from X-ray searches. Our scenario predicts a slight increase in the effective number of neutrino species N_{eff} at recombination, potentially detectable by future CMB missions.

KEYWORDS: dark matter theory, neutrino masses from cosmology, cosmology of theories beyond the SM, particle physics - cosmology connection

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1 Introduction

Non-zero neutrino masses leave imprints on cosmological observables, which in turn can be used to constrain their properties in several cosmological surveys [1]. In combination with *Planck* cosmic microwave observations [2], the latest results from the DESI collaboration [3] place a stringent bound on the sum of neutrino masses in the standard Λ CDM model:

$$\sum m_\nu \equiv \sum_{i=1}^3 m_i < 0.072 \text{ eV (95\% CL)}, \quad (1.1)$$

which can be compared with the earlier bound from the *Planck* collaboration $\sum m_\nu < 0.12 \text{ eV}$ (95% CL) [2]. As these bounds from cosmology are getting stronger and stronger, they are nearing or already in tension with the laboratory constraints on the same. In particular, neutrino oscillation data put a lower limit on the sum of neutrino masses: $\sum m_\nu > 0.058$ (0.098) eV (95% CL) for normal (inverted) neutrino mass ordering, when the lightest neutrino is massless ($m_{\text{lightest}} \rightarrow 0$) [4], implying that the DESI bound is in tension with inverted ordering at the 2σ level. Recently, it has even been argued that cosmology would

prefer a negative effective neutrino mass [5–7] (see, however, also [8–10]), suggesting that these bounds might become even more stringent in the future, leading to significant discord between the cosmology and neutrino oscillations [11].

At the same time, the current best model-independent limit on the neutrino mass measurement from the KATRIN experiment gives $m_\beta < 0.45$ eV [12], which would correspond to $\sum m_\nu \lesssim 1.35$ eV. Similarly, the search for neutrinoless double beta decay ($0\nu\beta\beta$) to prove the Majorana nature of neutrinos at the KamLand-Zen experiment also gives a constraint on the effective Majorana mass, $m_{\beta\beta} < 0.028\text{--}0.122$ eV (90% CL) [13], which gives $m_{\text{lightest}} < 0.084\text{--}0.353$ eV at 90% CL.

Hence, if our universe is well described by the standard cosmological scenario, we might not be able to observe finite absolute neutrino mass in the laboratory. If we are to expect positive results from the ongoing experiments or if cosmological constraints become significantly inconsistent with neutrino oscillations, we need a mechanism to reconcile the bounds from cosmology and laboratory. In the literature, several non-standard scenarios to relax the cosmological mass bound have been proposed. For example, they invoke neutrino decays [14–20], time-varying neutrino masses [21–26], presence of dark radiation [27–30], and neutrinos with distribution different from Fermi-Dirac distribution [31–33]. Phenomenological implications of some of these so-called *large-neutrino mass cosmologies* have been discussed in ref. [34].

Another cosmological puzzle that can be related to the neutrino sector is the problem of dark matter (DM). Indeed, if sterile neutrinos exist in nature, they can be considered viable dark matter candidates, see refs. [35–37] for reviews. While keV sterile neutrino DM produced via the Dodelson-Widrow mechanism is excluded, there have been attempts in the literature to modify the standard picture to accommodate sterile neutrino DM by introducing active neutrino self-interactions [38–40] and sterile neutrino self-interactions [41–44], see also [45, 46]. Moreover, in refs. [47, 48], it was shown that it is possible to have thermal DM below the MeV scale, if the DM comes into equilibrium with the standard model (SM) below the temperature of neutrino-photon decoupling.

In this work, we present a way to relax the cosmological neutrino mass bound with the assistance of a dark sector comprising dark radiation and sterile neutrino dark matter (DM) by employing self-interactions of active and sterile states. We depart from the model from ref. [27], which is based on a *Minimal Extended Seesaw* scenario [49] and realises a mechanism formulated by Farzan and Hannestad [28] to allow for large neutrino masses. We show that this model can be extended in a straight forward way to accommodate a keV scale DM candidate, whose cosmic abundance is obtained via a freeze-out mechanism in the dark sector, reconciling it with constraints from astrophysical and cosmological observations.

The outline of the work is as follows: in section 2, we review the mechanism to relax the cosmological neutrino mass bound with a light dark sector. We then present the model to realize the mechanism in section 3. The dark matter phenomenology is discussed in section 4 and the constraints on the parameter space of the model, suppression of the neutrino mass as well as predictions for the effective relativistic degrees of freedom (i.e., the effective number of neutrino species N_{eff}) are provided in section 5. Finally, we conclude in section 6 and provide some appendices with additional details.

2 Relaxing the cosmological neutrino mass bound due to a light dark sector

We employ the mechanism to relax the cosmological bound on the neutrino mass introduced in [28] and further elaborated in [27]. Here, we review it briefly and discuss how it is modified in the presence of a DM candidate.

The bound on the sum of neutrino masses in eq. (1.1) is derived indirectly from the cosmological observations, which are sensitive to the energy density in neutrinos. The CMB and large-scale structure observations, which are insensitive to the exact distribution function of neutrinos [32], place a bound only on the energy density of non-relativistic neutrinos [2]

$$\Omega_\nu h^2 \equiv \frac{\sum m_\nu n_\nu^0 h^2}{\rho_{\text{critical}}} < 1.3 \times 10^{-3} \text{ (95\% CL)}. \quad (2.1)$$

As can be seen from the equation above, the cosmological observations thus place a bound on the product of neutrino masses and their number density

$$\sum m_\nu \times \left(\frac{n_\nu^0}{56 \text{ cm}^{-3}} \right) < 0.12 \text{ eV (95\% CL)}, \quad (2.2)$$

with n_ν^0 denoting the background neutrino number density per helicity state. For illustration purposes, here we use the bound from *Planck* [2]. Hence, if we introduce a mechanism that can reduce the number density of neutrinos, larger neutrino masses can be accommodated.

At earlier times, when neutrinos are ultra-relativistic, their energy density is characterized by the number of effective ultra-relativistic neutrino species, N_{eff}

$$N_{\text{eff}} \equiv \frac{8}{7} \left(\frac{11}{4} \right)^{4/3} \left(\frac{\rho_{\text{rad}} - \rho_\gamma}{\rho_\gamma} \right), \quad (2.3)$$

where ρ_{rad} is the total energy density in relativistic species, and we have $N_{\text{eff}} \propto \langle p_\nu \rangle n_\nu$. Since the current measurement of $N_{\text{eff}} = 2.99 \pm 0.17$ [2] is in good agreement with the standard model (SM) calculation $N_{\text{eff}}^{\text{SM}} = 3.044(1)$ [50, 51], the decrease in n_ν must be compensated by the addition of light/massless beyond standard model species, and both this reduction of neutrino number density and addition of BSM states should happen before recombination. Moreover, in order to not spoil the successful predictions of Big Bang Nucleosynthesis (BBN), the mechanism should be activated once proton-neutron conversions have frozen-out around $T_\gamma \sim 0.7 \text{ MeV}$. Therefore, the mechanism would involve BSM states that thermalize with the neutrinos in the period between BBN and recombination, $10 \text{ eV} \lesssim T_\gamma \lesssim 100 \text{ keV}$, so that N_{eff} remains unchanged at earlier times.

As the neutrinos decouple at $T_\gamma \sim 2 \text{ MeV}$, they cannot be produced anymore from the particles in the SM thermal plasma. This implies that the production of new states would happen at their expense, thus reducing their number density and consequently relaxing the cosmological mass bound, such that cosmology becomes sensitive to the quantity

$$\left[\sum m_\nu \right]_{\text{eff}} = \sum m_\nu \frac{n_\nu}{n_\nu^{\text{SM}}}, \quad (2.4)$$

where n_ν^{SM} is the neutrino number density as predicted in the standard ΛCDM model. Hence, for $n_\nu < n_\nu^{\text{SM}}$, a large value of $\sum m_\nu$ can be accommodated. The suppression factor depends

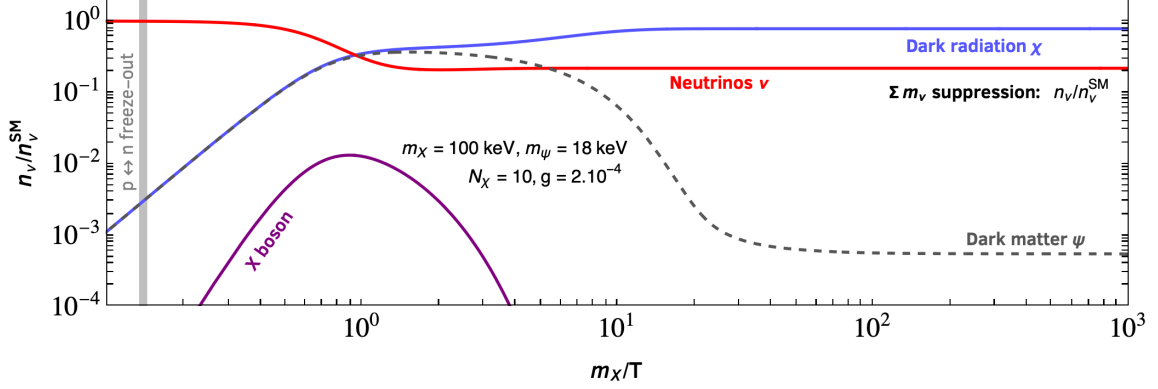


Figure 1. Schematic representation of the mechanism to reduce the active neutrino number density relative to the one in the Standard Model (shown in red) by equilibrating a dark sector consisting of $N_\chi = 10$ massless Dirac fermion species and the DM candidate ψ with mass $m_\psi = 18$ keV due to a mediator X with mass $m_X = 100$ keV and coupling $g = 2 \times 10^{-4}$. The X boson is produced initially from the inverse decays $\nu\bar{\nu} \rightarrow X$ and then decays dominantly to dark sector species, i.e. $X \rightarrow \chi\chi$ and $X \rightarrow \psi\psi$. The DM relic abundance $\Omega_\psi h^2 = 0.12 \pm 0.0012$ [52] is obtained once the interactions $\psi\psi \leftrightarrow \chi\chi$ freeze out.

on the massive and massless degrees of freedom in the dark sector, and we will discuss its values in our model in section 5.2 below.

Ref. [27] realized this mechanism by considering a new boson X with mass $100 \text{ eV} \lesssim m_X \lesssim 1 \text{ MeV}$, that thermalizes with neutrinos and numerous massless BSM species (χ) via processes such as $\nu\bar{\nu} \rightarrow X$ and $X \rightarrow \chi\chi$, with the requirements $\langle \Gamma(\nu\bar{\nu} \rightarrow X) \rangle > H$ and $\langle \Gamma(X \rightarrow \chi_i + \chi_j \text{ or } \nu) \rangle > H$ in the temperature window $10 \text{ eV} \lesssim T \lesssim 1 \text{ MeV}$, where H is the Hubble expansion rate. This reduces the neutrino population to produce the new boson, which then decays dominantly into the massless states χ , thereby populating the dark radiation component at the expense of the massive SM neutrinos.

In this work, we enlarge the dark sector to also contain a light ($\mathcal{O}(\text{keV})$) fermionic DM candidate ψ with mass $m_\psi \gtrsim 10 \text{ keV}$. In the case in which $m_X > 2m_\psi$, the X boson produced at the expense of neutrinos will not only populate the dark radiation component but also produce the DM candidate via $X \rightarrow \psi\psi$, which thermalizes with the massless χ states and takes part in conversion processes such as $\psi\psi \leftrightarrow \chi\chi$, analogous to multi-component DM frameworks [53–55]. Alternatively, in the case where $m_X < m_\psi$, the DM can still be produced by $2 \leftrightarrow 2$ interactions such as $\psi\psi \leftrightarrow XX$ in addition to $\psi\psi \leftrightarrow \chi\chi$, which would establish an equilibrium in the dark sector.

At temperatures below m_ψ , the $2 \leftrightarrow 2$ conversions between ψ , χ and X will freeze-out, resulting in a slight increase in the χ population as well as determining the DM relic abundance. A schematic illustration of this mechanism is shown in figure 1 for a chosen set of benchmark values, that allows us to relax the bound on the sum of neutrino masses by a factor $n_\nu/n_\nu^{\text{SM}} \approx 0.2$. In the next section, we present a UV complete model to realize this mechanism. The parameter space of the scenario is rather constrained by various astrophysical and cosmological bounds, which we discuss in section 5.1. In general, dark sector-neutrino interactions such as $\nu\bar{\nu} \leftrightarrow X$ are potentially also subject to some constraints from the

	Field	Species	U(1) _X
Scalar	Φ	1	+1
Fermions	χ	N_χ	-1
	ψ	1	-1
	N	3	0
	N'	1	0

Table 1. List of particles added to the SM. We give the number of copies for each type of particle (“Species”) and their charge assignments.

laboratory experiments, e.g., involving decays of SM particles such as mesons ($\pi^\pm, K^\pm$), τ leptons and $Z_0, W^\pm$ gauge bosons, which we discuss at the end of section 3.

3 Model

For a minimal realization of the mechanism described above, we extend the SM by adding in total N_{heavy} heavy right-handed neutrinos (denoted by N, N'), involved in the generation of neutrino masses via the seesaw mechanism. A suitable choice for our purposes will be 3 copies of N and one copy of N' , i.e., $N_{\text{heavy}} = 4$. Furthermore, we introduce a dark sector comprising a scalar singlet Φ and $N_{\text{light}} = N_\chi + 1$ generations of light vector-like fermions (denoted by χ, ψ). The dark sector particles ψ, χ and Φ are charged under a $U(1)_X$ symmetry, which can be either global or local. In ref. [27], it was shown that the gauged version is more appealing than the global one, therefore, in the following we will consider the former case. In table 1, we list the new particles and their charge assignments.

Further, a $\mathbb{Z}_2$ symmetry is imposed, under which all particles except χ_R and ψ_R are even, in order to avoid the vector-like mass terms and their interactions with Φ and N . (The χ_R and ψ_R fields are needed only for anomaly cancellation.) Thus, the Lagrangian consistent with the charge assignments can be written as

$$\begin{aligned}
-\mathcal{L}_{\text{int}} = & Y_\nu \bar{N} l_L \tilde{H}^\dagger + Y_\chi \bar{N} \chi_L \Phi + Y_\psi \bar{N} \psi_L \Phi + Y'_\nu \bar{N}' l_L \tilde{H}^\dagger + Y'_\chi \bar{N}' \chi_L \Phi + Y'_\psi \bar{N}' \psi_L \Phi \\
& + \frac{1}{2} M \bar{N} N^c + \frac{1}{2} M' \bar{N}' N'^c + \text{H.c.},
\end{aligned} \tag{3.1}$$

where l_L and H are the SM lepton and Higgs doublets respectively, $\tilde{H} = i\sigma_2 H^*$, and M, M' are the right-handed neutrino masses of N, N' , respectively. The scalar potential can be written as

$$V(H, \phi) = \mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 + \mu_\phi^2 |\Phi|^2 + \lambda_\phi |\Phi|^4 + \lambda_{H\phi} |\Phi|^2 H^\dagger H, \tag{3.2}$$

where we take $\lambda_{H\phi} \approx 0$, to avoid the mixing between the scalars. After electroweak symmetry breaking, the SM Higgs takes a VEV which can be parameterized as $\langle H \rangle = 1/\sqrt{2} [0, v_{\text{EW}}]^T$ with $v_{\text{EW}} = 246 \text{ GeV}$. Further, $\langle \Phi \rangle = v_\phi/\sqrt{2}$ breaks the $U(1)_X$ symmetry, with $v_\phi^2 = -\mu_\phi^2/\lambda_\phi$, also generating the mass for the associated gauge boson Z' , $m_{Z'} = gv_\phi$, where g is the $U(1)_X$ gauge coupling.

3.1 Neutral fermion mixing

The symmetry breaking induces mixing among the neutral leptons, and gives rise to the following $(3+N_{\text{light}}+N_{\text{heavy}}) \times (3+N_{\text{light}}+N_{\text{heavy}})$ mass matrix in the basis $(\chi_L^c, \nu_L^c, \psi_L^c, N', N)$:

$$\mathcal{M}_n = \begin{pmatrix} 0 & 0 & 0 & \Lambda' & \Lambda \\ 0 & 0 & 0 & m_{D'} & m_D \\ 0 & 0 & 0 & \kappa' & \kappa \\ \Lambda'^T & m_D^{T'} & \kappa'^T & M' & 0 \\ \Lambda^T & m_D^T & \kappa^T & 0 & M \end{pmatrix}, \quad (3.3)$$

where

$$\begin{aligned} m_D &= Y_\nu v_{\text{EW}}/\sqrt{2}, & m_{D'} &= Y'_\nu v_{\text{EW}}/\sqrt{2}, \\ \Lambda &= Y_\chi v_\phi/\sqrt{2}, & \Lambda' &= Y'_\chi v_\phi/\sqrt{2}, \\ \kappa &= Y_\psi v_\phi/\sqrt{2}, & \kappa' &= Y'_\psi v_\phi/\sqrt{2}. \end{aligned} \quad (3.4)$$

The rank of the matrix (3.3) is $2N_{\text{heavy}}$, leading to $(3+N_{\text{light}}-N_{\text{heavy}})$ massless and $2N_{\text{heavy}}$ massive states. For our purposes, we want 4 massive states in addition to the N_{heavy} heavy right-handed neutrinos: the 3 active neutrinos plus the DM candidate. Therefore, we chose $N_{\text{heavy}} = 4$, and $N_{\text{light}} = N_\chi + 1$, leaving N_χ states massless. One of the “light” dark sector fermions gets massive, which will become our DM candidate, and we single it out by denoting it with ψ to distinguish it from its massless partners χ .¹

Note that only left-handed fields appear in the Yukawa Lagrangian eq. (3.1) and receive masses according to (3.3), whereas the right-handed fields χ_R and ψ_R remain massless due to the postulated $\mathbb{Z}_2$ symmetry. Hence, we are left with

$$\tilde{N} = 2N_\chi + 1, \quad \tilde{g} = 4N_\chi + 2 \quad (3.5)$$

massless states and degrees of freedom in total, respectively, corresponding to χ_L, χ_R and ψ_R .

Upon block diagonalisation (see appendix A for details) we get the following mass eigenvalues

$$\begin{aligned} m_\chi &= 0, \\ m_\nu &= \frac{(m_D \kappa' - m_{D'} \kappa)^2 + (m_D \Lambda' - m_{D'} \Lambda)^2 + (\kappa' \Lambda - \kappa \Lambda')^2}{M'(m_D^2 + \kappa^2 + \Lambda^2) + M(m_{D'}^2 + \kappa'^2 + \Lambda'^2)}, \\ m_\psi &\approx \frac{m_D^2 + \kappa^2 + \Lambda^2}{M} + \frac{m_{D'}^2 + \kappa'^2 + \Lambda'^2}{M'}, \\ m_{N'} &\approx M', \\ m_N &\approx M. \end{aligned} \quad (3.6)$$

Here, $m_\nu = U_\nu^* \hat{m}_\nu U_\nu^\dagger$, where U_ν is the PMNS mixing matrix in the diagonal mass basis for the charged leptons and $\hat{m}_\nu = \text{diag}(m_1, m_2, m_3)$ contains the physical neutrino mass eigenvalues, and we use $a^2 = aa^T$ in order to write the equations for $m_{\nu, \psi}$ in a compact form.

¹Here we assume that all three active neutrinos are massive. If the lightest of them remains massless, we would need only $N_{\text{heavy}} = 3$ heavy right-handed neutrinos.

The mass eigenstates are then obtained as

$$(\hat{\chi}, \hat{\nu}, \hat{\psi}, \hat{N}', \hat{N})^T = \mathbf{W}^\dagger (\chi_L^c, \nu_L^c, \psi_L^c, N', N)^T, \quad (3.7)$$

where the mixing matrix $\mathbf{W}$ is given by

$$\mathbf{W} = \begin{pmatrix} 1 & \frac{\Lambda^*}{m_D^*} & 0 & \frac{\Lambda'^*}{M'^\dagger} & \frac{\Lambda^*}{M^\dagger} \\ \frac{-\Lambda^T}{m_D^T} & 1 & \frac{m_D'^*}{\kappa'^*} & \frac{m_D'^*}{M'^\dagger} & \frac{m_D^*}{M^\dagger} \\ 0 & \frac{-m_D'^T}{\kappa'^T} & 1 & \frac{\kappa'}{M'} & \frac{\kappa}{M} \\ \frac{-\Lambda'^T}{M'} & \frac{-m_D'^T}{M'} & \frac{-\kappa'^T}{M'} & 1 & 0 \\ \frac{-\Lambda^T}{M} & \frac{-m_D^T}{M} & \frac{-\kappa^T}{M} & 0 & 1 \end{pmatrix} \times \text{Diag}[1, U_\nu, 1, 1, 1]. \quad (3.8)$$

In deriving the mixing matrix above, we adopt a diagonal basis for the right-handed neutrino mass matrix, and assume the following hierarchy:

$$M \gg M' \gg m_D \gg \kappa', \Lambda \gg m_D', \Lambda', \kappa. \quad (3.9)$$

This corresponds to a regime where the dominant interactions of N' are with ψ whereas N interacts dominantly with χ and ν . Indeed, in the limit of $m_D', \Lambda', \kappa \rightarrow 0$, the mass matrix becomes block-diagonal and the $\psi - N'$ sub-block decouples. In this limit, we recover an additional $\mathbb{Z}_2$ symmetry under which ψ_L and N' are odd and all other fields even. Hence, the smallness of m_D', Λ', κ in eq. (3.9) can be motivated by this approximate $\mathbb{Z}_2$ symmetry. Indeed, for the phenomenology discussed below, including DM production, we do not need these couplings and the symmetry could even be exact, enforcing $m_D' = \Lambda' = \kappa = 0$. Note that given the assumed hierarchy, all the mixing terms are quite small, thus, contributions from oscillations amongst the light fields can be safely neglected.

Using the hierarchy above, we can simplify the expressions in eq. (3.6) for the neutrino and dark matter mass

$$\begin{aligned} m_\nu &\approx \frac{m_D^2 \kappa'^2}{M' m_D^2 + M \kappa'^2}, \\ m_\psi &\approx \frac{m_D^2}{M} + \frac{\kappa'^2}{M'}. \end{aligned} \quad (3.10)$$

We see that m_ψ receives two seesaw contributions, from both, N and N' . In order to have ψ as viable DM particle we need $m_\psi \gg m_\nu$, which can be achieved by imposing an additional condition, namely

$$M' m_D^2 \ll M \kappa'^2, \quad (3.11)$$

which does not follow from the hierarchy (3.9) but is compatible with it.² Then we obtain

$$\begin{aligned} m_\nu &\approx m_D M^{-1} m_D^T, \\ m_\psi &\approx \kappa' M'^{-1} \kappa'^T. \end{aligned} \quad (3.12)$$

²We can achieve eq. (3.11) by requiring $M \gg M'$. For instance, for large Yukawa couplings typically we will have $m_D \sim 100 \text{ GeV}$ and $\kappa' \sim 1 \text{ GeV}$, so we need $M'/M \ll 10^{-4}$.

Hence, the seesaw mechanism factorizes, such that the three N are responsible to generate active neutrino masses, whereas N' is responsible for the DM mass m_ψ . (This motivates our notation, to distinguish N and N' .)

Adopting the one-flavor approximation, the above mixing matrix can be expressed in terms of the mixing angles θ at the leading order

$$\begin{aligned}\theta_{\nu N} &= \frac{m_D}{M}, & \theta_{\nu\chi} &= \frac{\Lambda}{m_D}, & \theta_{N\chi} &= \frac{\Lambda}{M} \sim 0, \\ \theta_{\nu N'} &= \frac{m_{D'}}{M'}, & \theta_{\nu\psi} &= \frac{m_{D'}}{\kappa'}, & \theta_{N'\psi} &= \frac{\kappa'}{M'}.\end{aligned}\tag{3.13}$$

3.2 Parameters

Here, we discuss the parameters and interactions relevant for phenomenology. Using the seesaw relation in eq. (3.12), we can express $\theta_{\nu N}$ in terms of the neutrino mass as $\theta_{\nu N} = \sqrt{m_\nu/M}$, and similarly $\theta_{N'\psi} = \sqrt{m_\psi/M'}$. It can also be seen that $\theta_{\nu N'} = \theta_{\nu\psi}\theta_{N'\psi}$, hence, for our analysis, we are left with the following independent parameters

$$\{m_\nu, m_\psi, M', M, \theta_{\nu\chi}, \theta_{\nu\psi}, v_\phi, m_{Z'}, N_\chi\}.$$

We fix $m_\nu = 0.2 \text{ eV}$ for numerical estimates and are interested in the DM mass $m_\psi \sim \mathcal{O}(10\text{--}100) \text{ keV}$. For instance,

$$m_\psi \sim 10 \text{ keV} \left(\frac{\kappa'}{10^4 \text{ keV}} \right)^2 \left(\frac{10 \text{ GeV}}{M'} \right),\tag{3.14}$$

tells us about the order of magnitude of M' and $\kappa' = Y'_\psi v_\phi$. Note that the value of M is not relevant for the mechanism described above, however, it is constrained by perturbativity requirements and further constraints arise if we also want to incorporate thermal leptogenesis in the model [27].

For simplicity, we work in the one-flavor approximation and assume that the interactions between active neutrinos and all N_χ species have the same strength parameterized by $\theta_{\nu\chi}$. The mixing angle $\theta_{\nu\psi}$, instead, determines the interaction between neutrinos and DM which is subject to several constraints from DM stability and X-ray observations, as we discuss later. Hence, in order to evade these constraints, this mixing must be highly suppressed, i.e., $\theta_{\nu\chi} \gg \theta_{\nu\psi}$. In the limit of the above mentioned $\mathbb{Z}_2$ symmetry, we would even have $m'_D = 0$ and therefore $\theta_{\nu\psi} = 0$. Hence, to determine the viable parameter space of the model, the main parameters of interest are:

$$\{m_\psi, m_{Z'}, v_\phi, \theta_{\nu\chi}, N_\chi\}.\tag{3.15}$$

3.3 Interactions

Apart from the Yukawa interactions that can be directly read from the Lagrangian in eq. (3.1), we will also have gauge boson-mediated interactions. The interactions of the fermions with the $U(1)_X$ gauge boson can be written as

$$\mathcal{L} = \sum_f Q_f g Z'_\mu \bar{f} \gamma^\mu f,\tag{3.16}$$

where $f = \{\chi_L, \chi_R, \psi_L, \psi_R\}$,

$$g \equiv \frac{m_{Z'}}{v_\phi} \quad (3.17)$$

is the $U(1)_X$ gauge coupling, $Q_f = -1$ denotes the $U(1)_X$ charge for both the left and right-handed fermions. Given the mixing among the active and sterile neutrinos defined above, this interaction will give rise to the couplings $\lambda_{Z'}^{xx}$ leading to processes $Z' \leftrightarrow xx$

$$\begin{aligned} \lambda_{Z'}^{\nu\nu} &= \frac{m_{Z'}}{v_\phi} (\theta_{\nu\chi}^2 + \theta_{\nu\psi}^2) \simeq \frac{m_{Z'}}{v_\phi} \theta_{\nu\chi}^2, \\ \lambda_{Z'}^{\nu\chi} &= \frac{m_{Z'}}{v_\phi} \theta_{\nu\chi}, \\ \lambda_{Z'}^{\nu\psi} &= \frac{m_{Z'}}{v_\phi} \theta_{\nu\psi}, \\ \lambda_{Z'}^{\psi\psi} &= \lambda_{Z'}^{\chi\chi} = \frac{m_{Z'}}{v_\phi}. \end{aligned} \quad (3.18)$$

For our purposes, these gauge interactions are dominant compared to the real and pseudoscalars interactions between Φ and ν, χ which is further suppressed by m_ν . Therefore, we neglect the latter here. The couplings above will lead to various decays, inverse decays, and $2 \leftrightarrow 2$ processes, the rates of which are collected in appendix B.

We briefly comment on laboratory constraints from SM particle decays involving neutrinos, such as $Z_0, W^\pm, K^\pm, \pi^\pm$, which in principle are sensitive to neutrino- Z' or neutrino-(pseudo-)scalar interactions, as the emission of a mediator from the neutrino leg would lead to observable modifications of the decays, implying strong bounds on the corresponding coupling, see [56–61]. These bounds do not apply to our model, as all light mass states mix with the active neutrinos and χ 's contribute to the internal neutrino propagator. Therefore, the emission of the Z' is highly suppressed due to the approximate unitarity of the $\nu - \chi$ sub-block of the mixing matrix (which is unitary up to tiny corrections due to mixing with heavy states, $\theta_{\nu N}, \theta_{\nu N'}$). The same argument can be applied to interactions with the (pseudo-)scalar bosons, which in addition are further suppressed by a factor $m_\nu/m_{Z'}$ relative to Z' interactions. Therefore, such laboratory bounds do not lead to additional constraints on the parameter space of our model on top of the astrophysical and cosmological limits discussed in section 5.1 below.

4 Dark matter

The DM candidate gets its mass via the seesaw mechanism similar to the active neutrinos (see eq. (3.12)) while the χ 's remain massless. In figure 2, we show the contours of fixed DM mass for different values of M' and κ' . We also indicate the maximum bound on κ' from the perturbativity of the Yukawa coupling $\kappa'^{\max} \simeq \sqrt{4\pi} v_\phi^{\max} / \sqrt{2}$ by the dashed gray line. We take $v_\phi^{\max} \sim 10$ GeV, after considering several constraints that are discussed in the next section. Finally, we notice that considering $m_\psi \sim \mathcal{O}(10\text{--}100)$ keV corresponds to the upper bound on $M' \lesssim 10^8$ GeV. Below, we discuss the DM production, freeze-out and constraints from relic abundance, stability and structure formation.

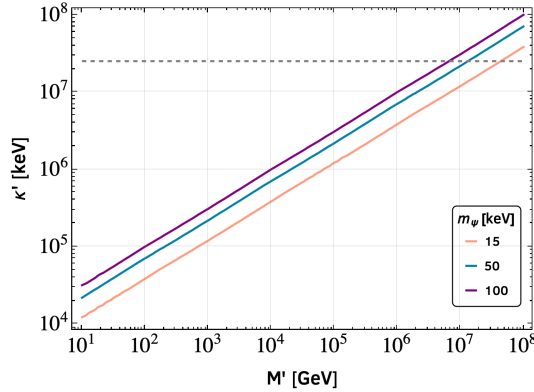


Figure 2. Contours of fixed DM mass $m_\psi = \{5, 15, 50, 100\}$ keV in the plane of N' mass M' and $\kappa' = Y'_\psi v_\phi / \sqrt{2}$. The gray dashed line indicates the bound from perturbativity of the Yukawa coupling Y'_ψ , taking $v_\phi^{\max} \sim 10$ GeV.

4.1 Production and freeze-out

At early times, DM may be produced from the decays of heavy neutrinos N, N' via the Yukawa interactions in eq. (3.1), while the production from $\nu \leftrightarrow \psi$ oscillations is negligible due to $\theta_{\nu\psi} \ll 1$. Later, once we have a significant population of Z' bosons, ψ is produced dominantly from Z' decays for $m_{Z'} > 2m_\psi$, or via processes such as $Z'Z' \leftrightarrow \psi\psi$ and $\chi\chi \leftrightarrow \psi\psi$ when $m_{Z'} < 2m_\psi$. In order to not spoil successful BBN, we need to make sure that these processes do not bring ψ or χ in equilibrium with the SM plasma before neutron-proton freeze-out.

To start, let's consider the DM production via heavy neutrino decays. Imposing the out-of-equilibrium condition $\Gamma < H$ gives $Y < 10^{-8}(M/\text{GeV})^{1/2}$. Since the Yukawa interactions of N with the SM lepton and Higgs doublet determine the ν masses, they are sizeable and expected to thermalize N with the SM plasma. In order to not bring ψ into equilibrium with the SM, we require that κ is sufficiently small-consistent with the hierarchy (3.9), such that N does not play a significant role in DM production. Furthermore, in some regions of the parameter space we have to assume a re-heating temperature $T_{\text{RH}} < M_N$, such that N are never produced, in order to avoid early equilibration of massless χ from N decays [27].

On the other hand, with the hierarchies (3.9), (3.11), for N' the relation $Y'_\nu < 10^{-8}(\frac{M'}{\text{GeV}})^{1/2}$ is fulfilled and therefore N' does not thermalize with the SM plasma. Thus also N' decays do not lead to a significant production of ψ , despite the sizeable coupling κ' required to generate the DM mass, see eq. (3.14) and figure 2.

At later times, the decays and scatterings involving Z' populate both ψ and χ , establishing an equilibrium in the dark sector, including also SM neutrinos. The equilibration happens after neutrino decoupling from the SM, therefore leading to a suppressed neutrino number density. Once $T < m_\psi$, the annihilations $\psi\psi \rightarrow \chi\chi$ and $\psi\psi \rightarrow Z'Z'$ (for $m_{Z'} < m_\psi$) will deplete the ψ population,³ eventually leading to their freeze-out and determining the DM relic abundance. This is analogous to the WIMP freeze-out, but in the dark sector characterized by a temperature T_{dark} with the DM freezing-out at a temperature of the order of a few

³Note that conversions such as $\psi\psi \rightarrow \nu\nu$ or $\nu\chi$ are also present, however, they are suppressed by the mixing angle $\theta_{\nu\chi}$ and subdominant enough to be safely neglected in determining the DM relic abundance.

	$m_{Z'} > 2m_\psi$	$2m_\psi > m_{Z'} > m_\psi$	$m_{Z'} < m_\psi$
Production	$Z' \leftrightarrow \psi\psi$	$Z'Z' \leftrightarrow \psi\psi$	$Z'Z' \leftrightarrow \psi\psi$
		$\chi\chi \leftrightarrow \psi\psi$	$\chi\chi \leftrightarrow \psi\psi$
Depletion	$\psi\psi \rightarrow \chi\chi$	$\psi\psi \rightarrow \chi\chi$	$\psi\psi \rightarrow \chi\chi$
			$\psi\psi \rightarrow Z'Z'$
Decay	$\psi \rightarrow \nu\chi\chi$	$\psi \rightarrow \nu\chi\chi$	$\psi \rightarrow \nu\chi\chi$
			$\psi \rightarrow Z'\nu$

Table 2. Processes contributing to the production, freeze-out and decay of DM candidate ψ in our model for different kinematical regimes.

keVs. This DM production mechanism is similar to the scenario discussed in refs. [47, 48], extended by the massless fermions in our model; see also refs. [62, 63] for discussions of DM freeze-out in a hidden sector.

In table 2, we summarize the dominant production and depletion channels ($\sigma \propto g^4, \Gamma \propto g^2$) for different kinematical regimes for $m_{Z'}$ and m_ψ . We also list the possible modes of DM decay ($\Gamma \propto \theta_{\nu\psi}^2$) in each regime. In order to track the evolution of number densities of different BSM species in presence of the aforementioned processes, we numerically solve a set of coupled Boltzmann equations collected in appendix B. A representative example is shown in figure 1 for the case $m_{Z'} > 2m_\psi$.

4.2 Relic abundance

The observed DM relic abundance $\Omega_\psi h^2 = 0.12 \pm 0.0012$ [52] is obtained once the interactions between ψ , χ and Z' freeze-out. The dominant DM annihilation rates are given by (see e.g., refs. [64, 65])

$$\begin{aligned}
(\sigma v)_{\psi\psi \rightarrow \chi\chi} &\approx \tilde{N} \frac{g^4}{48\pi} \frac{m_\psi^2}{(m_{Z'}^2 - 4m_\psi^2)^2} v^2, \\
(\sigma v)_{\psi\psi \rightarrow Z'Z'} &\approx \frac{g^4}{16\pi m_\psi^2} \left(1 - \frac{m_{Z'}^2}{m_\psi^2}\right)^{1/2} \left(1 + \frac{m_\psi^4}{m_{Z'}^4} v^2 + \mathcal{O}\left(\frac{m_{Z'}^2}{m_\psi^2}\right)\right),
\end{aligned} \tag{4.1}$$

and the thermal average can be obtained by substituting $\langle v^2 \rangle = 6/x_d$ with $x_d \equiv m_\psi/T_{\text{dark}}$. Here we take into account that the dark sector temperature T_{dark} will deviate from the photon temperature T_γ in our model, see discussion in section 5.2. For $m_{Z'} \lesssim m_\psi$, quickly the second term in the $\psi\psi \rightarrow Z'Z'$ cross-section dominates over the others. It can be seen that for $m_\psi < m_{Z'}$ the number of massless species $\tilde{N}$ plays an important role in determining the strength of $\psi\psi \rightarrow \chi\chi$ annihilations in addition to relaxing the cosmological mass bound.⁴ The DM relic abundance is computed as

$$\Omega_\psi h^2 \approx 2.755 \times 10^2 \left(\frac{m_\psi}{1 \text{ keV}} \right) Y_\psi^\infty, \tag{4.2}$$

⁴When not specifying chirality, we use a shorthand notation with ψ referring actually to the massive ψ_L , whereas χ collectively denotes the massless states χ_L, χ_R, ψ_R , compare eq. (3.5).

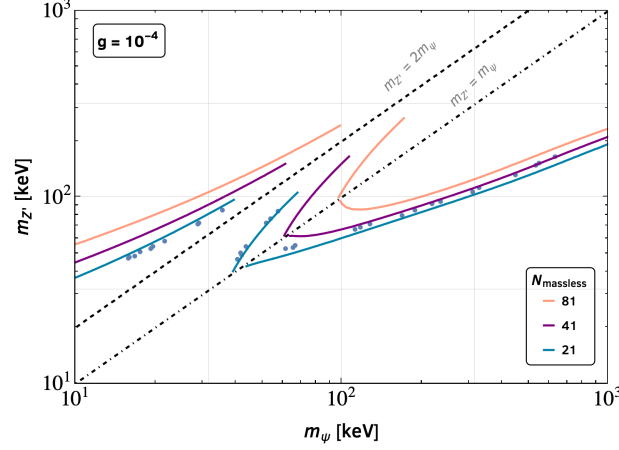


Figure 3. Contours of observed DM relic abundance for $g = 10^{-4}$ in the $m_{Z'} - m_{\psi}$ plane for different massless degrees of freedom, obtained from the analytical approximation of eq. (4.3). The blue dots represent the numerical solution for $\tilde{N} = 21$, matching closely with the analytical one. The dashed black line represents the s -channel resonance condition $m_{Z'} = 2m_{\psi}$, whereas the dash-dotted line corresponds to $m_{Z'} = m_{\psi}$, below which the t -channel annihilations $\psi\psi \rightarrow Z'Z'$ dominate.

where Y_{ψ}^{∞} is the asymptotic DM yield defined in appendix B and obtained by solving eq. (B.6) numerically.

Analogous to the visible sector freeze-out, we can use the analytical approximation for the relic abundance [66]. It is convenient to express the dark sector temperature in terms of the visible sector temperature by defining $\xi \equiv T_{\text{dark}}/T_{\gamma}$, so we have

$$\Omega_{\psi} h^2 \simeq x_f \frac{10^{-10} \text{ GeV}^{-2}}{\langle \sigma v \rangle_{\text{tot}}}, \quad (4.3)$$

with $\langle \sigma v \rangle_{\text{tot}} = \langle \sigma v \rangle_{\psi\psi \rightarrow \chi\chi} + \langle \sigma v \rangle_{\psi\psi \rightarrow Z'Z'}$. Here, $x_f = m_{\psi}/T_{\gamma,f}$ denotes the freeze-out temperature from the visible sector viewpoint, and it depends on ξ as [62]

$$x_f \approx \xi \ln \left[0.04\delta(\delta + 2)(g_{\psi}/g_{\text{eff}}^{1/2})\beta\xi^{3/2} \right] - \frac{1}{2}\xi \ln \left\{ \xi \ln \left[0.04\delta(\delta + 2)(g_{\psi}/g_{\text{eff}}^{1/2})\beta\xi^{3/2} \right] \right\}, \quad (4.4)$$

where $g_{\psi} = 2$ are the internal degrees of freedom of ψ , $g_{\text{eff}} = g_{\gamma} + g_{\text{dark}}\xi^4$ takes into account the contribution from the total relativistic degrees of freedom in the dark sector, δ is a numerical factor chosen to match the numerical solutions (we take $\delta \approx 0.7$ for p -wave annihilations), and $\beta = M_{\text{pl}} m_{\psi} \langle \sigma v \rangle_{\text{tot}}^{x=1}$ with $M_{\text{pl}} = 1.2 \times 10^{19} \text{ GeV}$. In our scenario, the dark sector temperature varies with the number of massless species $\tilde{N}$, as we will discuss in section 5.2, and therefore, x_f also depends on $\tilde{N}$ via the temperature ratio ξ . Using eq. (4.4), we find that DM freeze-out occurs at $x_f \sim 3$ – 5 for the relevant temperature window and annihilation cross-sections (in agreement with the numerical solution of the Boltzmann equations). In figure 3, we show the contours of observed relic abundance for a fixed value of the gauge coupling ($g = 10^{-4}$) and for varying $\tilde{N}$ in the $m_{Z'} - m_{\psi}$ plane using the analytical solution given above. For the case of $\tilde{N} = 21$, we also show the numerical result by solving the Boltzmann equations by blue points, matching closely the blue curve corresponding to the analytical approximation. For $m'_{Z'} > m_{\psi}$, where the annihilation to χ dominates, the dependence on $\tilde{N}$ follows explicitly from eq. (4.1).

In the region where annihilation to Z' dominates, i.e., when m_ψ is much larger than $m_{Z'}$, the $\tilde{N}$ dependence appears due to the modification of x_f via ξ . It should be noted that near the resonance region, i.e., when $m_{Z'} \simeq 2m_\psi$, the analytical solution and the thermally averaged cross-sections (eq. (4.1)) are not a good approximation, and a careful treatment of the cross-section and the freeze-out temperature is required in order to properly determine the DM relic abundance. Therefore, for dealing with resonances, we use the complete formula for thermal averaging (see appendix B) and compute the relic abundance numerically after solving the Boltzmann equations. We find that in the region very close to the resonance, the annihilation rate during freeze-out is so large that we do not obtain the correct abundance.

In figure 4, we show with solid lines of different colors contours of different DM mass which satisfy the relic abundance, highlighting the parameter space compatible with accommodating keV scale DM in the model as well as relaxing the cosmological neutrino mass bound. This is the main result of our work.

Note that in the limits far away from the resonance, the annihilation cross-sections from eq. (4.1) depend only on the combination $g/m_{Z'} = v_\phi^{-1}$ in both limits:

$$\langle\sigma v\rangle = \frac{m_\psi^2}{8\pi v_\phi^4 x_d} \times \begin{cases} \tilde{N} & (m_{Z'} \gg m_\psi) \\ 3 & (m_{Z'} \ll m_\psi) \end{cases}, \quad (4.5)$$

where we have used $\langle v^2 \rangle = 6/x_d$. Hence, in these limits, the DM relic abundance according to eq. (4.3) fixes the dark VEV for a given DM mass:

$$v_\phi \simeq 10^5 \text{ keV} \left(\frac{m_\psi}{15 \text{ keV}} \right)^{1/2} \left(\frac{3.2}{x_f} \right)^{1/2} \times \begin{cases} 2\tilde{N}^{1/4} & (m_{Z'} \gg m_\psi) \\ 2.4 & (m_{Z'} \ll m_\psi) \end{cases}, \quad (4.6)$$

in agreement with figure 4. For large mixing angles $\theta_{\nu\chi}$ (upper panel of figure 4), we are in the limit $m_{Z'} < m_\psi$ in large regions of the parameter space and therefore the DM contours are determined by the combination $m_\psi/v_\phi^2 = g^2 m_\psi/m_{Z'}^2$, and depend mildly on $\tilde{N}$ due to x_f .

4.3 Stability and X-ray constraints

In our model, the interaction of ψ_L and SM neutrinos with the massless states $\chi_{L,R}, \psi_R$ mediated by the gauge boson Z' can lead to DM decay via three body processes. The dominant ones are $\psi \rightarrow \nu\chi\chi$.⁵ Therefore, in order for the DM to be stable on cosmological timescales, we need to ensure that the lifetime of ψ relative to this channel is at least larger than the age of the universe, $t_U \approx 4.35 \times 10^{17}$ s. Using eq. (B.2) and trading the dependence on g and $m_{Z'}$ by v_ϕ we get

$$\tau_\psi \sim 45 \tau_U \frac{\left(\frac{v_\phi}{2 \text{ GeV}} \right)^4}{\left(\frac{\tilde{N}}{21} \right) \left(\frac{m_\psi}{15 \text{ keV}} \right)^5 \left(\frac{\theta_{\nu\psi}}{10^{-8}} \right)^2}. \quad (4.7)$$

⁵The decay $\psi \rightarrow 3\nu$ is also allowed, but the amplitude of this process is suppressed by the 3^{rd} power of the mixing angle $\theta_{\nu\chi}$, instead of the single power of $\theta_{\nu\chi}$ that suppresses $\psi \rightarrow \nu\chi\chi$. Therefore, to account for DM stability, we only need to consider $\psi \rightarrow \nu\chi\chi$. Moreover, due to the absence of $Z'\psi\chi$ interactions, decays such as $\psi \rightarrow 3\chi$ do not occur.

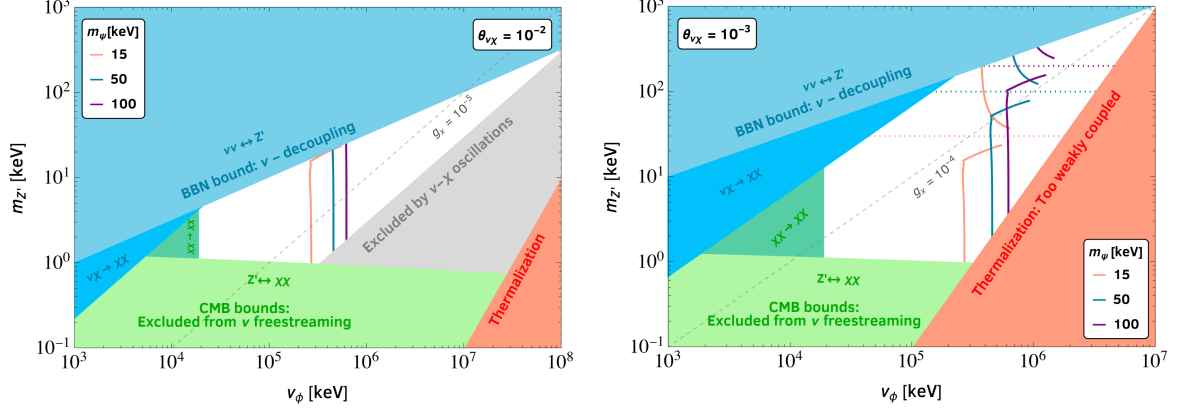


Figure 4. Parameter space of the model with the shaded areas highlighting regions of the parameter space excluded by several cosmological constraints, for a fixed value of $\nu - \chi$ mixing $\theta_{\nu\chi} = 10^{-2}$ (left) and $\theta_{\nu\chi} = 10^{-3}$ (right), and $N_\chi = 10$. Along the orange, blue, purple solid curves the observed relic DM density is obtained for $m_\psi = 15, 50, 100$ keV, respectively. The gray dashed lines indicate a fixed value of the $U(1)_X$ gauge coupling, $g = m_{Z'}/v_\phi$ and the dotted lines correspond to $m_{Z'} = 2m_\psi$ for a given DM mass. The red region is excluded from the thermalization condition, as the interactions of Z' with neutrinos are not strong enough. The blue regions are excluded from BBN by requiring that Z' is not in equilibrium with ν 's at $T > 0.7$ MeV, and the green regions show the area excluded by ν -free-streaming and CMB power spectra. The gray shaded region is excluded from production of χ via $\nu - \chi$ oscillations before BBN.

To be safe from cosmological constraints on the DM lifetime [67, 68], we impose $\tau_\psi > 20 \tau_U$, which gives

$$\theta_{\nu\psi}^2 < 2 \times 10^{-16} \left(\frac{15 \text{ keV}}{m_\psi} \right)^5 \left(\frac{21}{\tilde{N}} \right) \left(\frac{v_\phi}{2 \text{ GeV}} \right)^4. \quad (4.8)$$

We see that the DM lifetime bound sets a strong limit on the DM-neutrino mixing. We also observe preference for large values of v_ϕ in order to ensure the stability for $\mathcal{O}(10)$ keV DM.

On the other hand, if $m_\psi > m_{Z'}$, the $\psi \rightarrow Z'\nu$ decay channel opens up. Since it is a two-body decay it is comparatively unsuppressed with respect to the three-body decay discussed above (see eq. (B.2)) and the relative DM lifetime is given by

$$\tau_\psi \equiv \frac{1}{\Gamma_{\psi \rightarrow Z'\nu}} \sim 25 \tau_U \frac{\left(\frac{m_{Z'}}{10 \text{ keV}} \right)^2}{\left(\frac{g}{10^{-4}} \right)^2 \left(\frac{\theta_{\nu\psi}}{10^{-15}} \right)^2 \left(\frac{m_\psi}{40 \text{ keV}} \right)^3}, \quad (4.9)$$

where the approximation holds well for $m_\psi > 3m_{Z'}$. Similar to the previous decay, imposing $\tau_\psi > 20 \tau_U$, we get

$$\theta_{\nu\psi}^2 < 1.2 \times 10^{-30} \left(\frac{m_{Z'}}{10 \text{ keV}} \right)^2 \left(\frac{10^{-4}}{g} \right)^2 \left(\frac{40 \text{ keV}}{m_\psi} \right)^3, \quad (4.10)$$

which is much stronger than the one in eq. (4.8). This implies that the regime $m_\psi > m_{Z'}$ can work only with an effectively vanishing value of $m_{D'}$.

Sterile neutrino DM that mixes with active neutrinos can decay sub-dominantly also via radiative decay at 1-loop level as $\psi \rightarrow \nu\gamma$. For DM with mass $m_\psi \sim \mathcal{O}(10)$ keV, this

process leads to an observable monochromatic line in the X-ray frequency band. In the past two decades, a number of telescopes have scrutinized the sky in search of a signature of this kind (see e.g., section 5.4 of ref. [69]) and ruled out part of the parameter space of sterile neutrinos, setting a constraint in the $(m_\psi, \theta_{\nu\psi})$ plane.⁶ The observable measured in this case is the flux of photons detected by the telescope pointing towards DM-dominated objects. The flux of emitted photons is proportional to the decay width [74]

$$\Gamma_{\psi \rightarrow \nu \gamma} = \frac{9 \alpha G_F^2}{256 \cdot 4\pi^4} \sin^2(2\theta_{\nu\psi}) m_\psi^5. \quad (4.11)$$

Using the data put together in ref. [75], one can estimate the constraint on m_ψ and $\theta_{\nu\psi}$ to be

$$\theta_{\nu\psi}^2 \lesssim 7.65 \times 10^{-13} \left(\frac{15 \text{ keV}}{m_\psi} \right)^5. \quad (4.12)$$

Hence, X-ray limits can be satisfied by choosing $\theta_{\nu\psi}$ sufficiently small. Comparing with eq. (4.8), unless v_ϕ is much larger than 2 GeV, we see that if the lifetime bound is satisfied, the X-ray bound (4.12) is also satisfied. Similarly, if the lifetime bound of eq. (4.10) is satisfied, the X-ray bound is always satisfied.

Notice that in our framework, the production of DM happens independently of its mixing with active neutrinos. Therefore, $\theta_{\nu\psi}$ can always be chosen small enough to satisfy the strong upper bounds of eqs. (4.8), (4.12) without affecting the DM abundance. In other words, this implies that we can safely neglect any contribution to the DM abundance from its production from active neutrinos via the Dodelson-Widrow mechanism.

4.4 Bounds on DM from structure formation

In our model we find typical dark matter masses in range of 10 to 100 keV, and therefore our candidate is subject to constraints from structure formation, as any other warm DM (WDM) candidate. Potentially large free-streaming scales prevent the formation of structure at small scales, leading to a suppressed matter power spectrum at small scales. The impact of WDM particles on structure formation depends essentially on two factors: the mechanism through which they are produced that determines their momentum distribution, and their mass. Typically, for each production mechanism, one can extract a lower bound on the mass of the DM candidate from different large-scale structure observables.

The impact of WDM candidates on structure formation is conventionally parametrized by considering the half-mode halo mass M_{hm} or half-mode length-scale λ_{hm} at which the power spectrum becomes suppressed by a factor 2 compared to the reference cold dark matter case. They are related by

$$M_{\text{hm}} = \frac{4\pi}{3} \rho_{\text{DM}} \left(\frac{\lambda_{\text{hm}}}{2} \right)^3 \approx 1.9 \times 10^7 M_\odot \left(\frac{\lambda_{\text{hm}}}{0.1 \text{ Mpc}} \right)^3 \quad (4.13)$$

where ρ_{DM} is the dark matter energy density today. The limit on M_{hm} depends on the used cosmological data and analysis assumptions. For instance, ref. [76] obtains upper bounds at

⁶In 2014, two different groups claimed to have found a signal of this kind in different datasets [70–72]. However, to date, not only the interpretation of the claimed signal as being produced from decay of sterile neutrino DM, but even its effective presence in the data is still a matter of debate in the community [73].

95% CL in the range $\log_{10}(M_{\text{hm}}/M_{\odot}) = 7.0\text{--}8.6$, whereas ref. [77] finds $M_{\text{hm}} < 4.3 \times 10^7 M_{\odot}$. In the following we will adopt $\lambda_{\text{hm}} < 0.1$ Mpc as a rough estimate for consistency with structure formation.

In order to estimate the suppression scale λ_{cutoff} in our model we follow closely the discussion in [48]. Our DM candidate comes into chemical equilibrium with the thermal bath of SM neutrinos, Z' and massless χ and therefore will acquire a thermal momentum distribution. After the chemical decoupling, its co-moving abundance is constant, but it remains in thermal contact with the dark radiation via elastic $\psi\chi \leftrightarrow \bar{\psi}\chi$ interactions. The cutoff scale λ_{cutoff} is determined when the DM decouples also kinetically from the dark radiation, see e.g. [78]. Let us denote the photon temperature at kinetic DM decoupling by T_{kd} . Then there are two physically different effects which set λ_{cutoff} : (i) After kinetic decoupling, the DM will free-stream and suppress the logarithmic growth of structure before matter-radiation equality at co-moving scales smaller than

$$\lambda_{\text{FS}} \approx \frac{1}{2} \int_{t_{\text{kd}}}^{t_{\text{MRE}}} dt \frac{v_{\psi}}{a(t)} \approx \frac{1}{2} \left(\frac{4\pi^3 g_{\text{eff}}}{135} \right)^{-1/2} \sqrt{\frac{\xi}{T_{\text{kd}} m_{\psi}}} \frac{M_{\text{pl}}}{T_0} \log \frac{T_{\text{kd}}}{T_{\text{MRE}}}. \quad (4.14)$$

Here MRE refers to matter-radiation equality, v_{ψ} is the DM velocity, $a(t)$ is the cosmic scale factor, ξ is the ratio of the dark sector and photon temperatures, and $T_0 = 2.7$ K is the photon temperature today. The last approximation assumes that the DM particle is non-relativistic at kinetic decoupling [48], such that $v_{\psi} = p_{\psi}/E_{\psi} \approx \sqrt{3\xi T_{\text{kd}}/m_{\psi}} (T/T_{\text{kd}})$ for $T < T_{\text{kd}}$. (ii) As long as the DM is kinetically coupled to the dark radiation bath, modes with wavelength smaller than the horizon will undergo acoustic oscillations [79], preventing the formation of structure at these scales. This effect introduces another cutoff scale set by

$$\lambda_{\text{AO}} = \int_0^{t_{\text{kd}}} \frac{dt}{a(t)} = \frac{1}{aH} \Big|_{\text{kd}} \approx \left(\frac{4\pi^3 g_{\text{eff}}}{45} \right)^{-1/2} \frac{M_{\text{pl}}}{T_{\text{kd}} T_0}. \quad (4.15)$$

As a rough requirement for the consistency of our DM candidate with structure formation we then require

$$\lambda_{\text{cutoff}} = \max(\lambda_{\text{FS}}, \lambda_{\text{AO}}) < 0.1 \text{ Mpc}. \quad (4.16)$$

Kinetic decoupling has been studied for a range of simplified models in ref. [78]. In order to estimate T_{kd} we adopt the approximate expressions developed in the appendix of [78]. In our model, T_{kd} depends on the VEV v_{ϕ} , the DM mass m_{ψ} , and the massless degrees of freedom $\tilde{g}$, where the latter determines also the ratio of the dark sector-to-photon temperatures ξ , see eq. (5.16) below. Requiring the correct DM freeze-out abundance fixes v_{ϕ} for given m_{ψ} and $\tilde{g}$, see eq. (4.6). We show T_{kd} as a function of m_{ψ} and two representative values of $\tilde{g}$ in the left panel of figure 5. For DM masses in the range 10 to 100 keV we find decoupling temperatures $0.3 \text{ keV} \lesssim T_{\text{kd}} \lesssim 10 \text{ keV}$, depending also on the hierarchy between $m_{Z'}$ and m_{ψ} . With this result we can calculate the free-streaming and acoustic oscillation cutoff scales from eqs. (4.14) and (4.15). For the relevant parameter regions we find that $\lambda_{\text{AO}} \gtrsim \lambda_{\text{FS}}$ and therefore λ_{AO} determines the cutoff scale. From the rough criterion in eq. (4.16) we find from the right panel of figure 5 that the lower bound on the DM mass in our model is

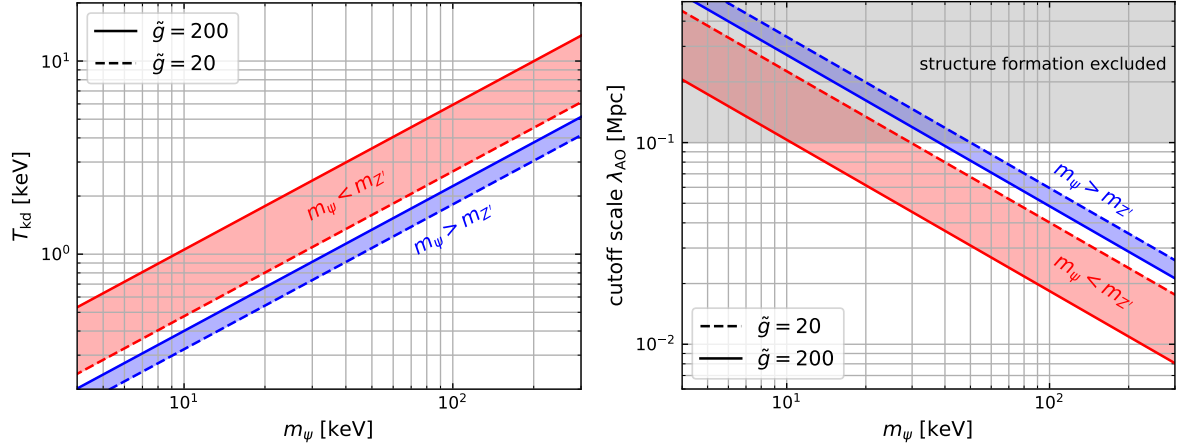


Figure 5. *Left:* photon temperature at kinetic decoupling of DM from the dark radiation bath. *Right:* cutoff scale due to acoustic oscillations λ_{AO} . In both panels we used the requirement of correct DM abundance and determined v_ϕ via eq. (4.6) either for $m_{Z'} \ll m_\psi$ (blue) or $m_{Z'} \gg m_\psi$ (red). The shaded bands indicate the range between two values of the massless degrees of freedom $\tilde{g} = 4N_\chi + 2$, where the dashed (solid) boundary of the bands correspond to $\tilde{g} = 20$ (200).

somewhere in the range from 10 keV to 50 keV, depending on the Z' - ψ mass hierarchy and the number of massless degrees of freedom. For increasing $\tilde{g}$ the lower bound decreases, as the kinetic decoupling occurs at higher temperatures. As we will see below, $\tilde{g} \approx 20$ is in tension with current limits on N_{eff} and therefore we require larger values than this. For $\tilde{g} = 200$ (corresponding to $N_\chi \approx 50$) we find that our model is consistent with structure formation for DM masses $m_\psi \gtrsim 10$ (40) keV for $m_\psi < m_Z$ ($m_\psi > m_Z$).

5 Relaxing the neutrino mass bound and predictions for N_{eff}

Let us now combine this DM production mechanism with the scenario of ref. [27] to relax the neutrino mass bound from cosmology. In section 5.1 we review the relevant constraints on the parameter space of the model, whereas in section 5.2 we discuss the thermalization of the dark sector, the relaxation of the neutrino mass bound and our predictions for N_{eff} .

5.1 Constraints and viable parameter space

The interactions between Z' and SM neutrinos are essentially unchanged from those of ref. [27], as the mixing between neutrinos and ψ is very suppressed. Therefore, the constraints on $Z' \rightarrow \nu\nu$ to determine the viable parameter space are also unchanged. In the following, for completeness, we review the most relevant constraints for the scenario.

Thermalization. From the discussion above, neutrinos should thermalize with the Z' boson in the temperature interval $0.7 \text{ MeV} > T_\gamma > 10 \text{ eV}$. The thermalization constraint is applied to the interaction $Z' \rightarrow \nu\nu$ rather than $Z' \rightarrow \nu\chi$, even though the former is suppressed by $\theta_{\nu\chi}$ as initially there are no χ present in the plasma. Since the thermally averaged decay rate peaks at $T \sim m_{Z'}/3$, we obtain a lower bound on the coupling $\lambda_{Z''}''$ by requiring the interaction rate $\langle \Gamma(Z' \rightarrow \nu\nu) \rangle > H(T \sim m_{Z'}/3)$. Using the expression for the Hubble rate

in the radiation dominated universe, we obtain

$$\lambda_{Z'}^{\nu\nu} \gtrsim 3.1 \times 10^{-12} \left(\frac{m_{Z'}}{\text{keV}} \right)^{1/2}. \quad (5.1)$$

The region where neutrinos do not thermalize with Z' in the early universe is highlighted in light red in figure 4.

BBN constraints. In order not to spoil BBN, Z' must not be in thermal equilibrium with neutrinos at $T > 0.7 \text{ MeV}$, otherwise it would deplete the population of neutrinos which participate in $p \leftrightarrow n$ conversions and would also contribute to the expansion rate of the universe along with the χ particles. Therefore, we require $\langle \Gamma(Z' \leftrightarrow \nu\nu) \rangle < H(T = 0.7 \text{ MeV})$, which leads to

$$\lambda_{Z'}^{\nu\nu} \lesssim 10^{-7} \left(\frac{\text{keV}}{m_{Z'}} \right). \quad (5.2)$$

This constraint is shown as the light blue region in figure 4. Further, if a minimal abundance of χ is present in the plasma before BBN (for example, having been produced from decays of N or N'), in order to avoid its exponential growth before BBN we need to impose

$$\langle \Gamma(\nu\chi \leftrightarrow \chi\chi) \rangle < H(T = 0.7 \text{ MeV}), \quad (5.3)$$

which gives a bound on the coupling

$$g \lesssim 10^{-3} \left(\frac{\theta_{\nu\chi}}{10^{-3}} \right)^{-1/2}, \quad (5.4)$$

excluding the dark blue shaded region in figure 4.

CMB constraints. In order for the CMB not to be distorted by $\nu\nu \rightarrow Z'$ and $Z' \rightarrow \chi\chi$ interactions, they must be rendered inefficient at $z < 10^5$, to be free of CMB constraints coming from neutrino free-streaming and power spectra [80], therefore, we require⁷

$$\langle \Gamma(\nu\nu \leftrightarrow Z') \rangle < H(T = 23 \text{ eV}). \quad (5.5)$$

Additionally, as Z' decays to a lot of massless states which are relativistic, we also require

$$\langle \Gamma(Z' \leftrightarrow \chi + \chi/\nu) \rangle < H(T = 23 \text{ eV}), \quad (5.6)$$

which excludes the region shaded in light green in figure 4. Finally, to ensure that CMB is not perturbed by the lack of free-streaming of χ , analogous to the free-streaming of ν , we impose

$$\langle \Gamma(\chi\chi \leftrightarrow \chi\chi) \rangle < H(T = 23 \text{ eV}), \quad (5.7)$$

which gives us a constraint: $v_\phi \gtrsim 2 \times 10^4 \text{ keV}$, shown by the bright green region in figure 4.

⁷We use the standard relation between temperature and redshift, $T(z) = T_0(1+z)$, where $T_0 = 2.3 \times 10^{-4} \text{ eV}$ is today's CMB temperature, and redshift $z = 10^5$ corresponds to $T(10^5) = 23 \text{ eV}$.

Constraints on active-sterile neutrino mixing. The $\nu - \chi$ mixing would also lead to the production of χ via oscillations in the early universe, well before the neutrinos decouple and BBN, therefore, the χ will contribute to the energy density of the universe and change N_{eff} . Therefore, requiring $\Delta N_{\text{eff}} < 0.3$, a bound can be placed on the mixing angle [27]

$$|\theta_{\nu\chi}| \leq 10^{-3} \left(\frac{10}{N_\chi} \right)^{1/2} \left(\frac{0.2 \text{ eV}}{m_\nu} \right)^{1/2}. \quad (5.8)$$

However, recently it was shown that considering the χ self-interactions [81], it is possible to have larger mixing angles, leading to the possibility of detecting χ from neutrino oscillation experiments. Varying the parameters of interest, one expects $\theta_{\nu\chi} \leq 10^{-4} - 10^{-1}$. The constraints from $\nu - \chi$ after considering the χ self-interactions are shown by the gray region in the top plot of figure 4, and is applicable only for $\theta_{\nu\chi} > 10^{-3}$. We find that the parameter space corresponding to the extreme limits, i.e., $\theta_{\nu\chi} = 10^{-1}, 10^{-4}$ is quite constrained and incompatible with successfully accommodating a DM candidate, hence we choose $\theta_{\nu\chi} = 10^{-3}, 10^{-2}$ in figure 4.

5.2 Suppression of neutrino mass, N_{eff} and dark sector temperature

Let us discuss now the equilibration of the dark sector in somewhat more detail. We assume that the new degrees of freedom come in equilibrium instantaneously with the active neutrinos when they have a temperature T_ν^{eq} and form a thermal system with a new temperature T_{eq} , with $m_{Z'}, m_\psi \ll T_{\text{eq}} \ll 1 \text{ MeV}$.⁸ For the parameter region of interest, the following interactions come into thermal equilibrium:

$$Z' \leftrightarrow ff, \quad Z' \leftrightarrow ff', \quad (5.9)$$

$$Z'Z' \leftrightarrow ff, \quad (5.10)$$

$$ff \leftrightarrow f'f', \quad (5.11)$$

with $f, f' = \nu, \chi, \psi$. As discussed above, the reaction $Z' \leftrightarrow \nu\nu$ is crucial to equilibrate the dark sector at first place, however, the less suppressed reactions with $f = \chi, \psi$ will be much faster once these particles are abundant. In particular, both, $1 \leftrightarrow 2$ and $2 \leftrightarrow 2$ processes involving Z' will simultaneously come into equilibrium. This implies that all chemical potentials are zero and there is no conserved particle number.⁹

⁸Note that our masses are $\mathcal{O}(10-100) \text{ keV}$ and actually the hierarchy is mild. Nevertheless, we take ψ and Z' as fully relativistic at equilibration for the estimates in this section. We have checked that they are in good agreement with the solution of the Boltzmann equations.

⁹This is different from the scenario considered in refs. [27, 28], where only the $1 \leftrightarrow 2$ processes from eq. (5.9) have been taken into account, but not the $2 \leftrightarrow 2$ processes from eq. (5.10). In that case, there is a conserved particle number (fermions f are created or destroyed only in pairs, while Z' numbers change by one unit) and chemical potentials develop, see also the discussion in appendix B of [27]. In contrast, in our model, the processes (5.10) proceed through t - and u -channel diagrams with $\sigma \sim g^4/T^2$, for which the interaction rates for coupling strength of $10^{-5} \lesssim g \lesssim 10^{-3}$ (cf. figure 4) are much larger than the expansion rate in the relevant temperature range.

Hence the system is fully characterized by its temperature T_{eq} , which is obtained from energy conservation during equilibration:

$$\rho_\nu(T_\nu^{\text{eq}}) = \sum_{f=\nu,\chi,\psi} \rho_f(T_{\text{eq}}) + \rho_{Z'}(T_{\text{eq}}), \quad (5.12)$$

with the energy density of a species i given by

$$\rho_i(T) = \left[\frac{7}{8} \right] \frac{\pi^2}{30} g_i T^4, \quad (5.13)$$

where the factor in square-brackets holds for fermions and should be dropped for bosons. Eq. (5.12) determines the ratio $T_{\text{eq}}/T_\nu^{\text{eq}}$ in terms of the involved degrees of freedom. Then the system evolves adiabatically to a temperature $T_{\text{fin}} \ll m_\psi, m_{Z'}$, when the DM ψ_L and Z' particles have become non-relativistic and therefore no longer contribute to dark radiation. Here we can use conservation of co-moving entropy:

$$a_{\text{eq}}^3 s_{\text{eq}}(T_{\text{eq}}) = a_{\text{fin}}^3 s_{\text{fin}}(T_{\text{fin}}), \quad (5.14)$$

with the cosmic scale factor a and $s = 4\rho/(3T)$ for massless particles. At the right-hand side of eq. (5.14) we take into account that only neutrinos and the massless degrees of freedom (χ_L, χ_R, ψ_R) contribute at T_{fin} . Combining eqs. (5.12) and (5.14) leads to the ratio of the dark sector temperature T_{dark} to the neutrino temperature in the SM (evaluated at $T_{\text{dark}} = T_{\text{fin}}$) as

$$\frac{T_{\text{dark}}}{T_\nu^{\text{SM}}} = \left(\frac{g_\nu + \tilde{g} + g_\psi + \frac{8}{7}g_{Z'}}{g_\nu + \tilde{g}} \right)^{1/3} \times \left(\frac{g_\nu}{g_\nu + \tilde{g} + g_\psi + \frac{8}{7}g_{Z'}} \right)^{1/4} \quad (5.15)$$

and the dark sector-to-photon temperature as

$$\xi \equiv \frac{T_{\text{dark}}}{T_\gamma} = \left(\frac{4}{11} \right)^{1/3} \frac{T_{\text{dark}}}{T_\nu^{\text{SM}}}. \quad (5.16)$$

Here $\tilde{g} = 4N_\chi + 2$ are the massless degrees of freedom and $g_\psi = 2, g_{Z'} = 3, g_\nu = 6$. Note that for the scalar in our model we expect masses of order of the dark VEV, with $10 \text{ MeV} \lesssim v_\phi \lesssim 10 \text{ GeV}$. Therefore, the scalar will not be present in the dark sector thermal bath in the relevant temperature range $T \lesssim 1 \text{ MeV}$ and does not contribute to the degrees of freedom here. For N_{eff} we then obtain

$$N_{\text{eff}} \equiv \frac{8}{7} \left(\frac{11}{4} \right)^{4/3} \frac{\rho_{\text{dark}}}{\rho_\gamma} = \frac{g_\nu + \tilde{g}}{2} \left(\frac{T_{\text{dark}}}{T_\nu^{\text{SM}}} \right)^4. \quad (5.17)$$

Using eq. (5.15), we see that this expression for N_{eff} converges to $g_\nu/2 = 3$ for both, $\tilde{g}, g_\psi, g_{Z'} \rightarrow 0$ as well as for $\tilde{g} \rightarrow \infty$. Note that this is the value for N_{eff} relevant for CMB observables, whereas for BBN N_{eff} in our scenario should be very close to the SM value. For the suppression of the neutrino number density we obtain

$$\frac{n_\nu}{n_\nu^{\text{SM}}} = \left(\frac{T_{\text{dark}}}{T_\nu^{\text{SM}}} \right)^3 = \frac{g_\nu + \tilde{g} + g_\psi + \frac{8}{7}g_{Z'}}{g_\nu + \tilde{g}} \left(\frac{g_\nu}{g_\nu + \tilde{g} + g_\psi + \frac{8}{7}g_{Z'}} \right)^{3/4}. \quad (5.18)$$

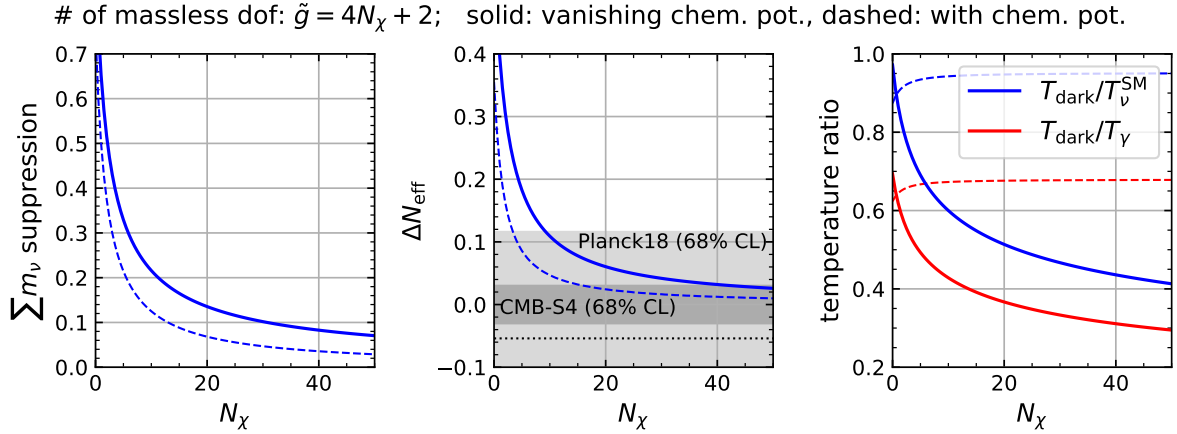


Figure 6. *Left:* the ratio n_ν/n_ν^{SM} , i.e., the effective suppression factor of $\sum m_\nu$, see eqs. (2.4), (5.18), *middle:* ΔN_{eff} , *right:* ratios of the dark-sector temperature to the neutrino temperature in the SM and to the photon temperature as a function of the number of χ generations N_χ , which is related to the number of massless degrees of freedom in our model by $\tilde{g} = 4N_\chi + 2$. Dashed curves correspond to the case with particle number conservation, shown for illustration purposes. The dotted line in the middle panel indicates the current *Planck* best fit point $\Delta N_{\text{eff}} = 2.99 - 3.044 = -0.054$ [2] with the 1σ interval ± 0.17 shown with light gray shading, and the dark shaded region indicates the sensitivity at $\pm 1\sigma$ of a future CMB-S4 mission [82].

The results of this calculation are shown in figure 6, where for comparison we show with dashed curves also the case when particle numbers are conserved, as considered in refs. [27, 28], see footnote 9. The latter are calculated following the procedure outlined in appendix B in ref. [27]. For given N_χ we observe a slightly smaller suppression of the neutrino mass bound and a larger contribution to N_{eff} than in the case with particle number conservation. For large $\tilde{g}$, the suppression factor in eq. (5.18) decreases like $\tilde{g}^{-3/4}$, compared to the factor $(1 + \tilde{g}/6)^{-1}$ for the case with number conservation [27, 28].

Our scenario predicts a sizable deviation of N_{eff} from the SM value at recombination. From the middle panel of figure 6 we see that small values of $N_\chi \lesssim 10$ are already in tension with the *Planck* result $N_{\text{eff}} = 2.99 \pm 0.17$ [2]. For the range $10 \lesssim N_\chi \lesssim 40$ (where also relevant relaxation of the neutrino mass bound occurs), deviations of N_{eff} from the SM value may be observable by future CMB missions, which should reach sensitivities of 0.07 [83] or even 0.03 [82] at 1σ . See also ref. [84] for a recent study of such dark radiation scenarios in the context of observables in the non-linear regime. Note that we predict an additional contribution to N_{eff} at CMB, whereas its value relevant for BBN should be largely unaffected in our model. Hence, different N_{eff} values inferred from BBN and CMB would be another potential signature of our model. Let us also note that the analysis in ref. [85] shows that recent DESI BAO measurements [3] leading to the tight bound on $\sum m_\nu$ from eq. (1.1) can be better accommodated by adding some amount of dark radiation (i.e., $\Delta N_{\text{eff}} > 0$), as predicted in our scenario.

Finally, we also comment on the rather different behavior of the dark sector temperature in the two cases. For conserved particle numbers, the ratio of the dark sector temperature to the neutrino temperature in the SM saturates around 0.95 for large $\tilde{g}$, whereas in the case with zero chemical potential, eq. (5.15) gives $1/\tilde{g}^{1/4}$ dependence for large $\tilde{g}$, leading to a significant cooling of the dark sector in our model.

M [GeV]	M' [GeV]	m_D [GeV]	κ' [GeV]	Λ [GeV]	v_ϕ [GeV]	m_ψ [keV]	$m_{Z'}$ [keV]	$g = m_{Z'}/v_\phi$	$\theta_{\nu\chi}$	N_χ	n_ν/n_ν^{SM}	ΔN_{eff}
10^{11}	10^2	4.47	0.043	0.004	0.5	18.5	100	2×10^{-4}	10^{-3}	10	0.216	0.109
10^{12}	10^3	14.14	0.23	0.141	0.8	53	77	9.6×10^{-5}	10^{-2}	10	0.216	0.109
10^{13}	10^2	44.7	0.1	0.044	0.6	100	32	5.3×10^{-5}	10^{-3}	20	0.135	0.060

Table 3. A set of benchmark points for the model where the mechanism to reduce the cosmological neutrino number density works and the observed DM relic abundance ($\Omega_\psi h^2 = 0.12 \pm 0.0012$) is obtained. The left columns show the relevant quantities from the fermion mass matrix and illustrate that our assumed hierarchy eq. (3.9) is satisfied. The middle columns indicate our main parameters of interest (see eq. (3.15)), and the right columns correspond to the effective suppression factor of $\sum m_\nu$ for cosmology and our predictions for ΔN_{eff} , respectively. For all three points the active neutrino mass $m_\nu = m_D^2/M$ is 0.2 eV.

6 Summary and discussion

We have presented a simple extension of the SM featuring a light dark sector with a U(1) gauge symmetry, embedded in a seesaw model for neutrino masses. The coupling of the dark sector to the SM occurs via mass-mixing of the dark fermions with neutrinos. The fermion spectrum contains heavy sterile neutrinos responsible for the seesaw mechanism, a large number of massless states χ , and a DM candidate ψ with mass in the 10 to 100 keV range. Both, active neutrino masses as well as the DM mass are generated via the type-I seesaw mechanism. We have identified regions of parameter space, where the following phenomenology is valid:

- The light dark sector particles play no role in the universe for temperatures above 1 MeV but come into thermal equilibrium with the SM neutrinos *after* neutrino decoupling from the SM but *before* recombination. This sets the typical scale for the dark gauge boson mass $m_{Z'} \sim 10$ keV.
- At late times the number density of active neutrinos is suppressed compared to the standard evolution, which allows to accommodate sizable neutrino masses with cosmology. The amount of neutrino suppression is determined by the number of massless degrees of freedom in the dark sector, which act as dark radiation during recombination.
- The DM fermion ψ first thermalizes together with the light dark sector, but then freezes out when the dark-sector temperature drops below its mass. Hence, the DM relic abundance is set by the dark sector gauge interactions and is independent of the DM mixing with SM neutrinos, which can be taken to be very small in order to satisfy constraints on DM stability, avoiding also measurable X-ray signatures.

Although the particle content of the model is complex, its free parameters are not numerous. Indeed, if we assume that the value of mixing between ν and χ is the same for all N_χ χ species, the phenomenology of the model is essentially determined by five independent parameters (see eq. (3.15)), which makes the model rather predictive. The parameter space of the model is well constrained and closed from all sides in many directions, see e.g., figure 4. In table 3 we give three benchmark points where all the above requirements are met. The left part of the table shows the relevant entries of the fermion mass matrix, which satisfy the assumed hierarchies from eqs. (3.9) and (3.11). The parameters m'_D, Λ', κ can be taken

arbitrarily small. In particular, a very small value of m'_D is required in order to have tiny mixing of the DM fermion with active neutrinos, as $\theta_{\nu\psi} = m'_D/\kappa'$. This is required in order to suppress the decays $\psi \rightarrow \nu\chi\chi$ and $\psi \rightarrow Z'\nu$ (if kinematically allowed) to have a sufficiently long-lived DM particle.

In the middle part of the table we provide the values of the VEV breaking the dark U(1) symmetry v_ϕ . Noting that the entries κ' and Λ are related to this VEV by Yukawa couplings (see eq. (3.4)), we confirm that these couplings remain within the perturbative regime. Typical gauge couplings are in the range of 10^{-5} to 10^{-4} . The mixing of the active neutrinos with the massless states $\theta_{\nu\chi} = \Lambda/m_D$ is constrained roughly as $10^{-3} \lesssim \theta_{\nu\chi} \lesssim 10^{-2}$, potentially testable in oscillation experiments only by saturating the upper bound.

The number of massless degrees of freedom in the dark sector, $\tilde{g}$ is related to the number of generations of χ fermions by $\tilde{g} = 4N_\chi + 2$ and sets the suppression factor of the neutrino density via eq. (5.18), which corresponds to the factor about which a bound on the neutrino mass from cosmology is relaxed. For instance, for the three benchmark points in table 3, the DESI bound $\sum m_\nu < 0.071$ eV from eq. (1.1) would be relaxed to 0.33, 0.33, 0.52 eV, respectively, making it safely compatible with the lower bounds from the neutrino oscillation data for both mass orderings and consistent with the value $m_\nu = 0.2$ eV assumed for the benchmark points.

While our specific DM candidate is very difficult to test observationally, the over-all model leads to indirect predictions. In particular, in our scenario we expect an increased energy density of dark radiation, expressed in terms of the effective number of neutrino generations N_{eff} . This effect emerges due to the disappearance of the massive Z' and DM particles from the relativistic bath, which leads to an effective heating of the dark sector. The amount of increase in N_{eff} decreases with $\tilde{g}$ (see eq. (5.17) and figure 6). Values of $N_\chi \lesssim 10$ are already in tension with the present CMB constraint on N_{eff} , whereas N_χ up to about 40 may lead to sizeable values in reach of planned CMB missions such as the Simons Observatory [83] or CMB-S4 [82]. With DM masses in the 10 to 100 keV range, we predict a suppression of the matter power spectrum at small scales. For instance, the three benchmark points from table 3 lead to suppression of structure formation at scales below 0.11, 0.051, and 0.054 Mpc, respectively, with the first point marginally consistent with observations, see section 4.4 for the corresponding discussion.

To conclude, the comparison of cosmological and laboratory determinations of neutrino masses may reveal exciting signs of new physics, if near future cosmological observations continue present tendencies and lead to increasingly strong bounds on the sum of neutrino masses and/or upcoming laboratory measurements of the absolute neutrino mass establish a non-zero value. We have presented an economical UV complete model offering an explanation in such a case. In addition to making large neutrino masses consistent with cosmology our model provides a simple seesaw explanation of neutrino masses, it incorporates a dark matter candidate in the 10 to 100 keV mass range in a straightforward way in the dark-fermion spectrum and its abundance is obtained naturally by freeze-out in the dark sector. Genuine predictions of the model are an increased value of the relativistic degrees of freedom at late times, potentially observable by future CMB observations and a suppressed matter power spectrum at small scales, as characteristic for warm DM candidates.

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A Diagonalisation of mass matrix

The matrix in eq. (3.3) can be reduced to a 2×2 form and block-diagonalised with a unitary transformation,

$$\mathbf{A}^T \begin{pmatrix} 0 & \Upsilon \\ \Upsilon^T & \mathbf{M} \end{pmatrix} \mathbf{A} = \begin{pmatrix} \hat{\mathbf{m}} & 0 \\ 0 & \hat{\mathbf{M}} \end{pmatrix}, \quad (\text{A.1})$$

where Υ is the 3×2 matrix containing the couplings and $\mathbf{M}$ is the diagonal matrix containing heavy neutrino masses. We choose the following Ansatz for $\mathbf{A}$ [86]

$$\mathbf{A} = \begin{pmatrix} \mathbb{1}_{3 \times 3} & \rho \\ -\rho^\dagger & \mathbb{1}_{2 \times 2} \end{pmatrix}, \quad (\text{A.2})$$

with the small mixing parameter ρ given by $\Upsilon^*(M^{-1})^\dagger$, and

$$\hat{\mathbf{m}} = -\Upsilon \mathbf{M}^{-1} \Upsilon^T, \quad \hat{\mathbf{M}} \sim \mathbf{M}. \quad (\text{A.3})$$

The resultant 3×3 matrix $\hat{\mathbf{m}}$ is a rank 2 matrix that can be further diagonalised to obtain the mass eigenvalues for the light fields

$$\mathbf{R}^T \hat{\mathbf{m}} \mathbf{R} = \text{diag}(m_\chi, m_\nu, m_\psi), \quad (\text{A.4})$$

where the rotation matrix $\mathbf{R}$ can be parameterized in terms of two mixing parameters (α, β) as follows

$$\mathbf{R} = \begin{pmatrix} 1 & \beta & 0 \\ -\beta^\dagger & 1 & \alpha \\ 0 & -\alpha^\dagger & 1 \end{pmatrix}, \quad (\text{A.5})$$

with $\beta \sim \Lambda^*/m_D^*$ and $\alpha \sim m_D'^*/\kappa'^*$. The mixing matrix $\mathbf{W}$ can then be expressed as a product of the matrices $\mathbf{A}, \mathbf{R}$, and further U_ν .

B Interaction rates and Boltzmann equations

The relevant expressions for N and Z' decay widths are

$$\Gamma_{N' \rightarrow lH/\chi\phi/\psi\phi} = \frac{M'}{8\pi} Y_{\nu/\chi/\psi}^{\prime 2}, \quad \Gamma_i \equiv \Gamma_{Z' \rightarrow ii} = N_i \frac{\lambda_{ii}^2}{24\pi} m_{Z'} \left(1 - \frac{4m_i^2}{m_{Z'}^2}\right)^{3/2}, \quad (\text{B.1})$$

where N_i is the number of generations of i Majorana fermions. The thermally averaged decay rates are given by [50]

$$\langle \Gamma_{a \rightarrow 1+2} \rangle = \Gamma_a \frac{K_1(m_a/T)}{K_2(m_a/T)}, \quad \langle \Gamma_{1+2 \rightarrow a} \rangle = \Gamma_a \frac{g_a}{2g_1} (m_a/T)^2 K_1(m_a/T), \quad (\text{B.2})$$

where $K_{1,2}$ are the modified Bessel functions of the first and second kind and $g_{a,1}$ are the internal degrees of freedom of the decaying and daughter particle.

The $2 \leftrightarrow 2$ cross-sections involving light/massless particles mediated via Z' are

$$\sigma \sim \frac{g^4}{16\pi^2} \frac{1}{m_{Z'}^4} T^2 \quad (m_{Z'} \gtrsim T), \quad \sigma \sim \frac{g^4}{16\pi^2} \frac{1}{T^2} \quad (m_{Z'} \ll T). \quad (\text{B.3})$$

The rate and thermal average for ψ to χ annihilation are given by

$$\begin{aligned} \sigma_{\psi\chi}(s) &= \tilde{N} \frac{g^4}{12\pi} \frac{(s + 2m_\psi^2)(1 - 4m_\psi^2/s)^{-1/2}}{(s - m_{Z'}^2)^2 + m_{Z'}^2 \Gamma_{Z'}^2}, \\ \langle \sigma v \rangle_{\psi\chi} &= \frac{1}{8T m_\psi^4 K_2(\frac{m_\psi}{T})^2} \int_{s_{\min}}^{\infty} ds s^{3/2} K_1\left(\frac{\sqrt{s}}{T}\right) \lambda(1, m_\psi^2/s, m_\psi^2/s) \sigma_{\psi\chi}(s), \end{aligned} \quad (\text{B.4})$$

where s is the centre of mass energy squared, $\Gamma_{Z'}$ is the total decay width of the mediator Z' and $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ is the Källén function. In the region far from resonances, we can expand σv in power of v using $s = 4m_\psi^2 + m_\psi^2 v^2$, and then take the thermal average to obtain eq. (4.1).

Finally, the DM decay widths are given by

$$\begin{aligned} \Gamma_{\psi \rightarrow \nu\chi\chi} &\simeq \tilde{N} \frac{m_\psi^5}{96\pi^3 M_{Z'}^4} (\lambda_{Z'}^{\chi\chi})^2 (\lambda_{Z'}^{\nu\psi})^2, \\ \Gamma_{\psi \rightarrow Z'\nu} &\simeq \frac{g^2 \theta_{\nu\psi}^2 m_\psi}{32\pi} \left(1 - \frac{m_{Z'}^2}{m_\psi^2}\right) \left(1 + \frac{m_\psi^2}{m_{Z'}^2} - \frac{2m_{Z'}^2}{m_\psi^2}\right). \end{aligned} \quad (\text{B.5})$$

We work with the following set of coupled Boltzmann equations for the evolution of number densities of the various BSM fields

$$\begin{aligned} \frac{dY_\nu}{dx} &= \frac{\langle \Gamma_\nu \rangle}{Hx} \left(Y_{Z'} - Y_{Z'}^{\text{eq}} \frac{Y_\nu^2}{Y_\nu^{\text{eq}2}} \right), \\ \frac{dY_{Z'}}{dx} &= \sum_{i=\nu,\chi,\psi} -\frac{\langle \Gamma_i \rangle}{Hx} \left(Y_{Z'} - Y_{Z'}^{\text{eq}} \frac{Y_i^2}{Y_i^{\text{eq}2}} \right) + \frac{s \langle \sigma v \rangle_{\psi\psi \rightarrow Z'Z'}}{Hx} \left(Y_\psi^2 - \frac{Y_\psi^{\text{eq}2}}{Y_{Z'}^{\text{eq}2}} Y_{Z'}^2 \right), \\ \frac{dY_\chi}{dx} &= \frac{\langle \Gamma_\chi \rangle}{Hx} \left(Y_{Z'} - Y_{Z'}^{\text{eq}} \frac{Y_\chi^2}{Y_\chi^{\text{eq}2}} \right) + \frac{s \langle \sigma v \rangle_{\psi\psi \rightarrow \chi\chi}}{Hx} \left(Y_\psi^2 - \frac{Y_\psi^{\text{eq}2}}{Y_\chi^{\text{eq}2}} Y_\chi^2 \right), \\ \frac{dY_\psi}{dx} &= \frac{\langle \Gamma_\psi \rangle}{Hx} \left(Y_{Z'} - Y_{Z'}^{\text{eq}} \frac{Y_\psi^2}{Y_\psi^{\text{eq}2}} \right) - \frac{s \langle \sigma v \rangle_{\psi\psi \rightarrow \chi\chi}}{Hx} \left(Y_\psi^2 - \frac{Y_\psi^{\text{eq}2}}{Y_\chi^{\text{eq}2}} Y_\chi^2 \right) \\ &\quad - \frac{s \langle \sigma v \rangle_{\psi\psi \rightarrow Z'Z'}}{Hx} \left(Y_\psi^2 - \frac{Y_\psi^{\text{eq}2}}{Y_{Z'}^{\text{eq}2}} Y_{Z'}^2 \right), \end{aligned} \quad (\text{B.6})$$

where $Y_i \equiv n_i/s$ is the yield, $s = (2\pi^2/45)g_{*s}T_\gamma^3$ is the entropy density, $H(T_\gamma) = 1.66\sqrt{g_{\text{eff}}}(T_\gamma^2/M_{\text{pl}})$ with $g_{\text{eff}} = g_\gamma + g_{\text{dark}}\xi^4$, and here we define $x \equiv m_{Z'}/T_\gamma$. The equilibrium yield of a particle with mass m_i is given by

$$Y_i^{\text{eq}} = 0.115 \frac{g_i}{g_{*s}} \xi x_i^2 K_2(x_i/\xi) \quad (\text{B.7})$$

with $x_i = m_i/T_\gamma$. For massless particles $Y_\chi^{\text{eq}} = 0.23(g/g_{*s})\xi^3$ and in the case of neutrinos corresponds to their yield after they have decoupled from the plasma. Note that the $2 \leftrightarrow 2$ processes involving ν 's are quite suppressed by the small mixing angles, and we have checked that their contribution is negligible compared to the processes considered above in eqs. (B.6). In the case $m_{Z'} < 2m_\psi$, the terms corresponding to $\langle \Gamma_\psi \rangle$ are not taken into account.

We solve these equations by taking negligible abundances for all the dark sector particles as initial conditions. Further, it should be mentioned that we assume a fixed value of the dark sector temperature when solving the Boltzmann equations. A precise treatment should also involve the temperature evolution equations. We have checked that our results for the DM yield are not very sensitive to how we chose the dark sector temperature, and therefore we are confident that this is a reasonable approximation.

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Asymmetries in extended dark sectors: aogenesis scenario

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ABSTRACT: The observed dark matter relic abundance may be explained by different mechanisms, such as thermal freeze-out/freeze-in, with one or more symmetric/asymmetric components. In this work we investigate the role played by asymmetries in determining the yield and nature of dark matter in non-minimal scenarios with more than one dark matter particle. In particular, we show that the energy density of a particle may come from an asymmetry, even if the particle is asymptotically symmetric by nature. To illustrate the different effects of asymmetries, we adopt a model with two dark matter components. We embed it in a multi-componentogenesis scenario that is also able to reproduce neutrino masses and the baryon asymmetry. In some cases, the model predicts an interesting monochromatic neutrino line that may be searched for at neutrino telescopes.

KEYWORDS: Models for Dark Matter, Particle Nature of Dark Matter

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1 Introduction

The nature of Dark Matter (DM), that makes up roughly a quarter of the energy density of our universe, along with the origin of neutrino masses and the baryon asymmetry (BAU), are among the most important open problems that the Standard Model (SM) fails to explain, with overwhelming experimental evidence. Several extensions of the SM have been proposed, where either a single particle or several stable particles, i.e., multi-component DM [1–9], make up the observed DM relic abundance, $\Omega_{\text{DM}} h^2 \sim 0.1$ [10]. The most popular mechanisms to reproduce this relic abundance include thermal freeze-out (FO) of weakly interacting massive particles (WIMPs) [11–16] and freeze-in (FI) of feebly interacting massive particles (FIMPs) [17–19].

WIMPs are initially in thermal equilibrium with the SM and undergo annihilations until they freeze-out when the annihilation rate drops below the Hubble expansion rate; therefore, the DM abundance is inversely proportional to the annihilation rate. On the other hand, FIMPs have a negligible initial abundance and are produced mainly by tiny interactions with particles in the thermal bath so that they never thermalise, and they freeze-in once the mother particle decouples from the bath; thus, the DM abundance is directly proportional to the production rate. Depending on the interactions, a further contribution may come from late decays (LD) of the mother particle, the size of which is model-dependent.

Most of these models assume that the new states are symmetric in nature, i.e., the abundance of the DM particle is the same as that of the antiparticle. However, there exist a wide variety of models that propose that the DM abundance is rather set by an initial asymmetry in the dark sector, in analogy to the visible sector (where $\eta_B = 0.88 \times 10^{-11}$), motivated by the closeness of baryonic and DM energy densities, $\rho_{\text{DM}} \sim 5\rho_B$ [20]. The asymmetry can be first generated in the visible sector and then transferred to dark sector (or vice versa) [21–23], or an asymmetry can be generated simultaneously in both the sectors (*cogenesis*). Such asymmetric dark matter (ADM) models (see refs. [24, 25] for a review) aim to explain the observed baryon asymmetry and DM abundance in a common framework.

An intermediate scenario between the two extremes involves the asymmetric freeze-out of a species, where the DM particle and its antiparticle freeze-out with different number densities, depending on the initial asymmetry, and the dark matter is partially asymmetric [26]. Hence, a further distinction can be made regarding the nature of DM, whether it is symmetric, asymmetric or partially asymmetric. Therefore, the presence of an asymmetry can have significant implications for the mechanisms discussed above in reproducing the DM abundance.

The possibility that asymmetric single-component WIMP DM is produced in a cogenesis scenario has been studied in ref. [27]. Another model has been presented in ref. [28], with the possibility to restore the symmetric nature of DM through late decays of an extra particle. In ref. [29] a model of symmetric multi-component DM, which combines freeze-out and freeze-in, accompanied by the generation of the baryon asymmetry, has been proposed. Finally, refs. [30–32] study models linking the freeze-in production of dark matter with baryogenesis; however, the dark matter in these models is symmetric.

In this work, we aim to generalize this picture, starting from the concreteogenesis scenario realized in ref. [28], by considering an extended dark sector in which we can realize multicomponent DM, so that one DM component, or even both, can be asymmetric. Furthermore, we combine freeze-out and freeze-in production, including the possibility of asymmetric freeze-in. The idea of asymmetric freeze-in was proposed in ref. [33], mainly focusing on simultaneously generating equal and opposite asymmetries in the visible and dark sectors, which are separately in thermal equilibrium at different temperatures. Unlike the case of asymmetric freeze-out, it was shown that this case is not so straightforward and it requires the presence of a richer dark sector, so the dark matter particle is able to thermalise with the dark states even if it is feebly coupled to the SM particles. In ref. [34], the unitarity and CPT constraints for transferring the baryon number to the dark sector via freeze-in were discussed. In ref. [35], it was shown that dark sector asymmetries produced by asymmetric freeze-in can be sizeable and comparable to the baryon asymmetry if the mediator between the two sectors is out-of-equilibrium. However, these works mainly focus on the generation of asymmetry but do not study the details of the annihilation of the symmetric component, which is a necessary ingredient for any ADM mechanism.

To this end, we propose a model in which the dark sector is connected to the visible sector via a mediator particle. The asymmetries in both sectors are generated via cogenesis, which yields a relation between neutrino masses, the baryon asymmetry and the DM relic abundance. In addition to the generation of asymmetries, we also discuss the annihilation of the symmetric component in a complete model with all the necessary ingredients. We propose a feebly interacting particle that can be an asymmetric DM candidate with no sizeable couplings with neither the SM particles nor the dark sector ones. We then investigate the role of the dark sector asymmetry in determining the relic abundance of one/several particles. For example, depending on the model parameters, the DM may either be single-component or multi-component, with either all symmetric or asymmetric components or a mixture of both. Similarly, the production mechanisms may be freeze-out, freeze-in or some via freeze-out and the others via freeze-in. Moreover, we show that in certain scenarios the particle abundance may be set by an asymmetry even if its nature is symmetric. We find that some of them predict the existence of an observable neutrino line.

The paper is structured as follows. In section 2, we elaborate on the framework of generating an asymmetry and the possibility of asymmetric freeze-in. We discuss the complete model, which also generates neutrino masses and the baryon asymmetry, in section 3. The different DM candidates are discussed in section 4. The contribution to the DM relic abundance is studied in section 5. In section 6, we discuss the main phenomenological signatures, with special emphasis on the prediction of a monochromatic neutrino line. In section 7, we discuss a low scale variant of the model (an inverse seesaw). Finally, we give our conclusions in section 8. We show further details in some appendices.

2 General framework

In order to generate a dark sector asymmetry, we focus on the cogenesis of both DM and baryon asymmetries and adopt the two-sector thermal leptogenesis mechanism of ref. [28].

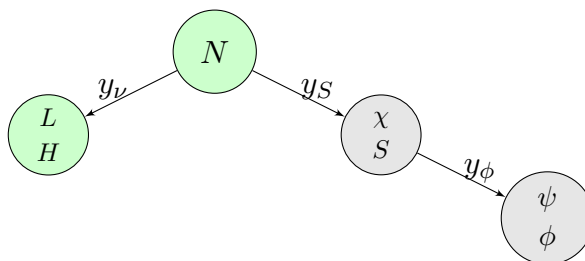


Figure 1. Schematic framework of cogenesis and DM production in the model at $T < M_{N_1}$ via the indicated Yukawa interactions y_i . Here, L and H are the SM lepton and Higgs doublet, respectively, whereas S and ϕ are complex scalars belonging to the dark sector.

In this case, the asymmetries are produced from out-of-equilibrium CP-violating decays of right handed neutrinos (RHNs), which are well-motivated to reproduce neutrino masses. This requires the leptons and some of the dark sector particles to be charged under a lepton symmetry that is broken by Majorana masses of the RHNs, thus satisfying all the conditions for the dynamical generation of an asymmetry [36]. The RHNs are initially in thermal equilibrium, and once the temperature drops below the mass of the lightest RHN, M_{N_1} , the washout and other interactions leading to transfer of asymmetries between the two sectors become inefficient. Subsequently, the asymmetries get frozen in the two sectors. The leptonic asymmetry is partially converted into the baryonic one via sphaleron processes. The dark asymmetry ($\eta_i \equiv Y_i^+ - Y_i^-$, where we introduce the yield as the number density upon entropy density, $Y_i = n_i/s$) is carried by a dark fermion (we denote it by χ) and once its symmetric component gets annihilated away, the asymmetric one sets the relic abundance, $\Omega_{\text{DM}} \propto \eta_\chi m_\chi$. This framework, therefore, connects neutrino mass generation with the baryon asymmetry and the DM relic abundance. See refs. [37–44] for other extensions of seesaw framework that address the three issues under the same umbrella.

In our work, we go a step further and enlarge the dark sector so that the different possibilities discussed in the introduction are feasible. For this purpose, the particle χ in our set-up is not the dark matter but rather decays to another stable fermion (say ψ), which may constitute all or part of the DM abundance. The schematic framework is shown in figure 1. Therefore, the asymmetry in χ can be transferred to ψ via its decays. Indeed, the set-up offers a richer phenomenology as it is now possible to accommodate multi-component DM (in this case, it could be ψ and S), as well as different dynamics thanks to the presence and size of the different interactions. In principle, we can have four cases of equilibration between different sectors, as illustrated also in ref. [45]. In figure 2, we show the four cases of entropy transfer between the various sectors: the SM + N_R (green), χ, S, ϕ (orange) and ψ (gray).¹

The first case involves equilibration among all sectors. It is the multi-component freeze-out scenario that has been widely considered (see for instance ref. [1]), and we will

¹In the scenarios considered below, ϕ , even if it comes from χ decays (see figure 1), is included in the middle blob.

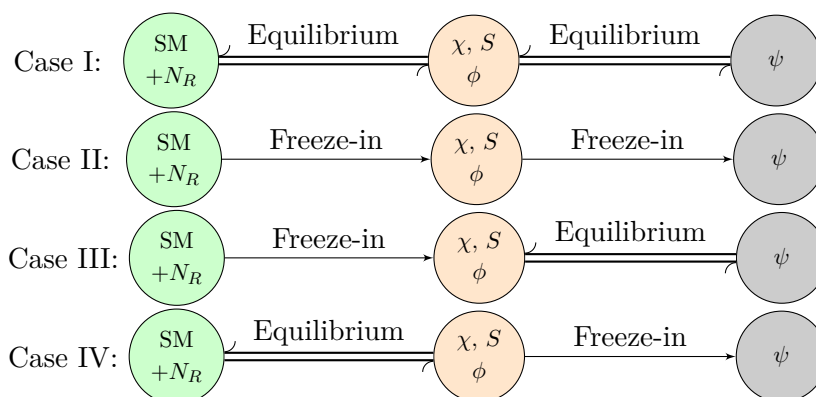


Figure 2. The four scenarios for entropy transfer between the SM sector and the dark sectors formed by χ, S and ψ, ϕ . Similar figure in ref. [45].

not consider it here in the following. Cases II and III correspond to scenarios where χ is not in equilibrium with the $\text{SM} + N_R$ sector, which implies that $y_S \ll 10^{-7}$. In the considered framework where the asymmetries are produced from the decays of heavy right-handed neutrinos, such a small Yukawa coupling would then be unable to generate a dark asymmetry comparable to the visible one, and therefore DM would have to be symmetric. Therefore, we are interested in the last case (Case IV), where:

- i)* the orange sector is in equilibrium with the $\text{SM} + N_R$,
- ii)* the dark sector asymmetry may be generated via co-genesis, of size comparable to the baryonic one, and
- iii)* the ψ sector is not in equilibrium and is produced via freeze-in.

Whether the abundance produced by freeze-in is asymmetric depends on the value of Yukawa y_ϕ , because the ψ population can be symmetric even if the mother particle χ carries an asymmetry. This can be understood as follows. Let us define $x_i \equiv m_i/T$, for a species of mass m_i , where T is the temperature. The asymmetry freezes out at $x_i \sim 20$, whereas the freeze-in from early decays takes place around the mass of the mother particle, i.e., $x_i \sim 1$, when both the mother particle and its antiparticle are in equilibrium. So, if the production from early decays is greater than the asymmetric yield after freeze-out, i.e., $Y_\psi = Y_{\bar{\psi}} > Y_\chi \approx \eta_D$, then the daughter particle will be symmetric in nature, because the production from late-decays (that take place at $x_i \gg 1$) will be sub-dominant and negligible. However, if the production from early decays is smaller than the asymmetry, $Y_\psi < \eta_D$, then the late decays that are active much later after the asymmetry has frozen-in will produce more ψ than $\bar{\psi}$, hence generating an asymmetry in ψ . Realizing this last scenario therefore yields an example of asymmetric freeze-in, which up to our knowledge has not been studied in the literature.

In ref. [46], the authors discussed the idea of having an asymmetry in a messenger particle with a long enough lifetime so that it decays to the dark matter particle once its symmetric component is annihilated away. In that case, however, the non-thermal decay to

dark matter was suppressed by higher-dimensional operators, which require new physics at some higher energy scale. On the other hand, in our scenario the decays are suppressed due to the feeble nature of the interaction, whose strength controls whether the contributions come from either early decays, late decays or both. This in turn determines the nature of ψ population.² Therefore, our renormalizable model does not require any extra new physics.

As we will see, in certain scenarios, the late decays of an asymmetric particle may populate the symmetric component of a species. Hence, the late decays play an important role in determining the final nature of a species. In order to distinguish between the scenarios, we use the notation of ref. [26] to define the asymmetric ratio for a particle species as

$$r_i \equiv Y_i^-/Y_i^+ \quad \text{with} \quad 0 < r_i \leq 1, \quad (2.1)$$

where $+$ ($-$) denotes the particle (antiparticle). Here, the upper (lower) limit in r signifies that the species is completely symmetric (asymmetric) and $\Omega_{\text{DM}} \propto Y_i^+ + Y_i^-$. The asymptotic asymmetric ratio of a species i with mass m_i can be written as [26]

$$r_i^{(\infty)} \simeq \exp \left[-\sqrt{\frac{\pi g_*}{45 x_f}} M_{\text{Pl}} \langle \sigma v \rangle_i \eta_D m_i \right], \quad (2.2)$$

where $\langle \sigma v \rangle_i$ is the thermally-averaged annihilation cross section, $M_{\text{Pl}} \simeq 1.2 \times 10^{19}$ GeV, m_p is the mass of the proton and $x_f \equiv m_i/T_* \sim 20$ is the mass over freeze-out temperature ratio. In the following, we use $r_i^{(\infty)} \equiv r_i$ to alleviate notation. As shown in ref. [1], different regimes appear:

- For $r_i < 10^{-2}$, the behaviour of DM is highly asymmetric (A).
- In the range, $10^{-2} < r_i < 0.9$, DM behaves as partially asymmetric (PA).
- For $r_i > 0.9$, DM is highly symmetric (S).

In the rest of the paper we adopt these ranges to define the nature of DM. Further classification can be made on the basis of the dominant production mechanism that determines the asymptotic nature of the dark matter, be it freeze-out (FO), early decays from freeze-in (FI), or late decays (LD). In this work, we aim to focus exclusively on scenarios where the asymmetry is directly involved in reproducing the relic abundance. For this goal, we propose a complete model that displays the different roles played by the asymmetries.

3 A model for neutrino masses, the baryon asymmetry, and dark matter

Our initial hypothesis for the construction of a model that explains dark matter, neutrino masses and the baryon asymmetry in a common framework is that the symmetry baryon minus lepton number, $B-L$ (under which all SM quarks have charge $1/3$ and all SM leptons have charge -1), plays a key role. This is one of the best-motivated symmetries beyond the SM: it is accidental and anomaly-free in the SM, and when gauged it requires the presence

²In ref. [47], the contributions from early and late decays have been compared for a symmetric DM model.

Field	Spin	U(1) _{B-L}	U(1) _D	U(1) _X
N_R^i	1/2	-1	0	0
σ	0	+2	0	0
χ_0	1/2	-1	1	0
ψ_0	1/2	0	0	+1
S	0	0	-1	0
ϕ	0	+1	-1	+1

Table 1. Particle content of the model and their respective charge assignments under Lorentz and the U(1) groups. The first two states correspond to the sterile neutrino sector, and the last four to the dark sector.

of three sterile neutrinos (which in turn naturally generate active neutrino masses) and is easily embedded in GUTs. When considering a complex asymmetric dark sector, however, this symmetry is not enough, and we require extra U(1)s to forbid Majorana masses and some interaction terms, as well as to annihilate the symmetric components.

Therefore, we augment the SM gauge group by 3 new gauge U(1) symmetries: U(1)_{B-L}, and the dark product U(1)_D ⊗ U(1)_X. We add three right-handed neutrinos N_R (RHNs) to cancel the gauge anomalies of the first one. We also extend the particle content by the addition of two *Dirac dark* fermions, ψ_0 and χ_0 , and 3 scalars σ , S and ϕ . The gauge bosons associated with the three new groups are Z_{B-L}^0, Z_D^0, A'^0 , with gauge couplings g_{B-L}, g_D and g_X , respectively. The quantum numbers of the new fields are summarised in table 1. Note that all the new fields are SM singlets. The new part of the Lagrangian of the model can thus be written as

$$\mathcal{L}_{\text{new}} = \mathcal{L}_{\chi\psi}^0 + \mathcal{L}_{\text{kin}}^0 + \mathcal{L}_{\text{int}} - V(\sigma, S, \phi, H), \quad (3.1)$$

where $\mathcal{L}_{\text{kin}}^0$ includes the kinetic terms of the gauge bosons and the scalars, $V(\sigma, S, \phi, H)$ is the most general scalar potential that one can write given the symmetries of the model, where H is the SM Higgs doublet field, and

$$\begin{aligned} \mathcal{L}_{\chi\psi}^0 &= \bar{\chi}_0(i\not{D} - m_{\chi}^0)\chi_0 + \bar{\psi}_0(i\not{D} - m_{\psi}^0)\psi_0, \\ \mathcal{L}_{\text{int}} &= -y_{\nu}^{\alpha i} \bar{L}^{\alpha} \tilde{H} N_R^i - y_{\sigma}^{ij} \sigma \bar{N}_R^{ic} N_R^j - y_S^i S \bar{N}_R^i \chi_0 - y_{\phi} \phi \bar{\psi}_0 \chi_0 + \text{H.c.} \end{aligned} \quad (3.2)$$

Here $m_{\chi,\psi}^0$ are bare mass terms of the dark fermions, while the indices $\alpha = e, \mu, \tau$ and $i = 1, 2, 3$ run over the generations of leptons and right handed neutrinos, respectively. We use $\tilde{H} = i\sigma_2 H^*$. Note that y_{ν} is a 3×3 general complex matrix, y_{σ} is a 3×3 complex symmetric matrix, y_S is 3-component vector and y_{ϕ} is a complex number. However, several phases are unphysical. Four phases of y_{σ} may be removed by rephasing σ and N_R^i . Two phases of y_S may be removed by rephasing S and χ_0 . y_{ϕ} can be taken real by rephasing ϕ or ψ .

In accordance with the framework discussed above, we take

$$y_{\phi} \ll 1, \quad g_X \ll 1, \quad (3.3)$$

so that ψ_0 cannot thermalise with the SM bath. We further assume a vanishing initial abundance $n_{\psi} = 0$, consistent with an inflationary epoch. On the other hand, the particles

χ_0, σ, S, ϕ and N_R^i reach thermal equilibrium with the SM thermal bath through sizeable new gauge (g_{B-L} and g_D) and scalar interactions.

Note that the Lagrangian also contains the *kinetic mixing* between the U(1) gauge factors. An unavoidable contribution to kinetic mixing among all U(1)s arises at one-loop level with ϕ running in the loop. Also, χ contributes in the case of $U(1)_{B-L} - U(1)_D$ mixing. Hence, the kinetic mixing $-(\kappa/2)Z_{B-L}^{\mu\nu}Z_{D,\mu\nu}$ is naturally of the order of $\kappa \gtrsim g_D g_{B-L}/(16\pi^2) \sim 10^{-3} g_{B-L} g_D$. The kinetic mixing of $U(1)_X$ with the other U(1) factors is $\gtrsim g_X g_i/(16\pi^2)$. Since $g_X \ll 1$ by assumption, this contribution can be safely ignored. Finally, a kinetic mixing among Z_{B-L} and the SM $U(1)_Y$ gauge boson is generated through a loop of SM quarks and leptons, of the order $\gtrsim g_Y g_{B-L}/(16\pi^2) \sum_{i=q,l} Y_i (B-L)_i$. Notice that, given the conservation of $U(1)_{\text{em}}$, the photon does not couple to the $B-L$ current. In any case, the bounds on the kinetic mixing are not relevant for the range of parameters that we consider in the paper.

Finally, let us remark that the necessary ingredients for the model to work could also be achieved with a global $U(1)_{B-L}$ [16, 48], explicitly violated by right-handed neutrino masses, since: i) χ is Dirac in nature because of the gauge $U(1)_D$, ii) it has sizeable interactions with the SM (with N_R, y_S) to thermalise, and iii) it can undergo efficient annihilations due to the $U(1)_D$ gauge group. However, such scenario is not as theoretically appealing as the gauged $B-L$ version that we consider, which demands the existence of 3 right-handed neutrinos. In the following subsections, we discuss the scalar, gauge, fermionic and dark sectors in detail.

3.1 The scalar sector and spontaneous symmetry breaking

The details of the full scalar potential $V(\sigma, S, \phi, H)$ are quite involved as one can write quadratic, quartic and mixed quartic terms for each combination of the scalar fields σ, S, ϕ and the SM Higgs doublet H . However, without entering into the details of the scalar potential, we can safely assume that there is a region of the parameter space in which σ takes a large vev, $\langle \sigma \rangle = v_{B-L} \gtrsim 10^{11}$ GeV, which breaks the $U(1)_{B-L}$ symmetry by 2 units and generates Majorana masses for the sterile neutrinos as well as a mass for the $U(1)_{B-L}$ gauge boson, Z_{B-L} . The value $v_{B-L} \gtrsim 10^{11}$ GeV is chosen so that the lightest RHN mass (M_{N_1}) safely obeys the equivalent Davidson-Ibarra lower bound [49] to achieve thermal Leptogenesis in the model [28], but any larger value will not change the following analysis and conclusions. Therefore, an asymmetry can be generated once the inverse decays of N_1 go out of equilibrium [50, 51]. Note that we take $M_{Z_{B-L}} > M_{N_1}$ and $m_\sigma > M_{N_1}$, so that the heavy Z_{B-L} gauge boson and the radial component of σ are therefore naturally very heavy and decay fast into quark and leptons.

The other scalar, ϕ , takes a much smaller vev than that of σ , i.e., $\langle \phi \rangle = v_\phi \ll v_{B-L}$, which breaks $U(1)_D \otimes U(1)_X \rightarrow U(1)_{X+D}$. After symmetry breaking, we can write $\phi(x) = v_\phi + \varphi(x)/\sqrt{2}$. The overall symmetry breaking pattern of the model can then be represented as

$$U(1)_{B-L} \otimes U(1)_D \otimes U(1)_X \xrightarrow{\langle \sigma \rangle} U(1)_D \otimes U(1)_X \xrightarrow{\langle \phi \rangle} U(1)_{X+D}. \quad (3.4)$$

The only fields that are charged under the unbroken $U(1)_{X+D}$ symmetry are the fermions χ_0, ψ_0 with charge +1 and the scalar S with charge -1. We discuss in section 4 the consequences of this for the DM stability.³

A phenomenologically relevant parameter is the mixing between the SM Higgs boson h (coming from $H = (v_{EW} + h)/\sqrt{2}$) and the scalar φ , which can be characterised by the mixing angle θ . This is directly related to the mixed quartic term $\lambda_{H\phi}\phi^\dagger\phi H^\dagger H$ in the scalar potential. In the rest of the paper, we assume that this mixing angle is small ($\sin\theta \simeq \theta \ll 1$) so we can safely trade h and φ for the mass eigenstates.⁴ Finally, the scalar S does not obtain a vev, $\langle S \rangle = 0$, and may therefore be a DM candidate in the model because of its charge under $U(1)_{X+D}$ (see discussion at the beginning of section 4).

We are interested in the regime

$$M_{N_3}, M_{N_2} \gg M_{N_1} \gg m_\chi^0 \gg m_\psi^0, m_S > m_\phi, \quad (3.5)$$

so that the decay channels shown in figure 1 are kinematically open. We consider values of m_ψ^0 and m_S of similar order of magnitude, in the GeV ballpark, because we are interested in scenarios where both may contribute significantly to the DM relic abundance.

3.2 The gauge sector

The symmetry breaking pattern outlined in eq. (3.4) leads to one massless (two massive) gauge boson(s), which correspond to the unbroken (broken) generator(s). Up to small corrections, the masses are given by

$$m_{A'}^2 = 0, \quad m_{Z_D}^2 = 2g_D^2 v_\phi^2, \quad m_{Z_{B-L}}^2 = 8g_{B-L}^2 v_{B-L}^2, \quad (3.6)$$

where A', Z_D, Z_{B-L} are the mass eigenstates (the full expressions can be found in appendix B.2). The tiny values of the parameters g_X, κ and v_ϕ/v_{B-L} suppress the mixing between the gauge bosons so that the mass eigenstates mostly coincide with the original eigenstates. The massless A'_μ is decoupled from all the other fields (as $g_X \ll 1$) and does not play any role in the following discussion (see appendix B.2 for a discussion on the bounds on a massless dark gauge boson). The mass/kinetic mixing among Z_D and Z_{B-L} induces an interaction of the type $Z_D^\mu J_\mu^{B-L}$, where J_μ^{B-L} is the $B-L$ current,⁵ which may lead to decays of Z_D into SM quarks and leptons. However, the decay width is suppressed by $(v_\phi/v_{B-L})^4$ and it is subleading with respect to other decay channels:

- If $m_{Z_D} > 2m_\chi$, the gauge boson decays into $\bar{\chi}\chi$ or $S^\dagger S$ at tree level, or into $\varphi\varphi$ at one-loop level through a loop of S . This last process depends on the coupling between φ and S . As we will show in appendix B.1, the particles φ must have a fast decay to SM fermions, so this gives a lower bound on their mass, and therefore on the mass of Z_D .

³By construction, the model does not require imposing a discrete symmetry to ensure the stability of the DM components.

⁴In appendix B.1, we briefly discuss the mixing when studying the decays of φ to SM particles. We show that even a very small mixing angle allows efficient decays.

⁵Note that in the absence of the mass mixing there would not be such an interaction even for $\kappa \neq 0$ [52].

- If $2m_S < m_{Z_D} < 2m_\chi$, the gauge boson decays mainly as $Z_D \rightarrow S^\dagger S$. The decays have width $\sim \mathcal{O}(g_D^2 m_{Z_D}/100)$ and are very fast in the relevant temperature regime ($T < v_\phi$).⁶
- If $2m_S > m_{Z_D} > 2m_\varphi$, the gauge boson can only decay into 2φ , see the first point.

In the rest of the paper we focus on the second scenario, namely $2m_S < m_{Z_D} < 2m_\chi$.

3.3 The fermionic sector

We first discuss the generation of tiny neutrino masses via the type-I see-saw mechanism [53–58]. The breaking of $U(1)_{B-L}$ gives masses to the heavy RHNs, $M_{N_i} \sim y_\sigma^i v_{B-L} \gtrsim 10^{11}$ GeV. In the following we use $M_{N_1} \lesssim v_{B-L}$, i.e., we take $y_\sigma^1 \lesssim 1$. The masses for the active neutrinos are given by the seesaw expression,

$$m_\nu = -m_D M_N^{-1} m_D^T, \quad (3.7)$$

with $m_D = y_\nu v_{EW}/\sqrt{2}$. For the dark fermions, once ϕ obtains a vev, a mixing is induced between χ_0 and ψ_0 due to the Yukawa coupling y_ϕ , see eq. (3.2). We define the fermion mixing parameter as

$$\epsilon_f \equiv \frac{y_\phi v_\phi}{m_\chi^0 - m_\psi^0}, \quad (3.8)$$

where $y_\phi \ll 1$ and typically $v_\phi < m_\chi^0$. Thus, the mixing parameter is highly suppressed, $\epsilon_f \ll 1$ (we remind that we are assuming hierarchical masses, i.e., $m_\chi^0 \gg m_\psi^0$). Upon diagonalising the fermionic sector, we have verified that the mass eigenstates (χ, ψ) mostly coincide with the original eigenstates (χ_0, ψ_0) , while the corrections to the masses are $\mathcal{O}(\epsilon_f^2)$ and therefore negligible, see appendix B.3 for more details. Therefore, in the following we drop the subscripts 0, trading the original masses/fields for the physical ones.

Finally, let us mention that no Majorana masses for χ or ψ are generated due to the preserved gauge $U(1)_{X+D}$, due to the fact that S does not take a vev. Higher dimensional operators may be written at dimension 6 and 8 for χ and ψ , such as $\overline{\chi^c} \chi \sigma S S$ and $\overline{\psi^c} \psi \sigma S S \phi^\dagger \phi^\dagger$. However, as S does not take a vev, these operators do not generate Majorana masses.

4 Dark matter components

Recall that the only fields that are charged under the remnant $U(1)_{X+D}$ symmetry are χ, ψ (with charge +1) and S (with charge −1). As $m_\chi \gg m_\psi, m_S$, only ψ, S may be stable. The $X + D$ charge is preserved in decays of the type $\psi \rightarrow S^\dagger + P$ or the opposite $S \rightarrow \bar{\psi} + P$, where P is some uncharged state. Therefore, in principle the lightest among ψ and S is stable and would be the only DM candidate.

However, the decay $\psi \rightarrow S^\dagger + P$ (or the opposite) is suppressed by the masses of the right-handed neutrino N_1 and χ in the propagators. At low energies, $E \ll m_\chi \ll M_{N_1}$, one can integrate out both the particles and study the decay in terms of higher-dimensional

⁶We ignore the possibility $m_{Z_D} \gtrsim 2m_S$ in which the phase space of the decay closes.

operators. As we analyse below in section 6, P corresponds to active neutrinos in the model, i.e., $P = \nu$, yielding a monochromatic neutrino line from ψ decays [59–63]. In that section we show that in a broad region of the parameter space, the decays are suppressed on cosmological timescales and obey current limits, and therefore both particles (ψ and S) contribute significantly to the DM relic abundance. The scenario in which the decay is fast enough and there is only 1 DM candidate has been studied extensively in the literature, so we will not discuss it here. Hence, we focus on the two-DM scenario. The conservation at low energies of the $U(1)_{X+D}$ symmetry yields the constraint

$$\begin{aligned} 0 &= Q_{X+D} = \eta_\psi Q_\psi + \eta_S Q_S \\ &\implies \eta_\psi = \eta_S, \end{aligned} \quad (4.1)$$

where in the last step we used that the $U(1)_{X+D}$ charges are given by $Q_\psi = -Q_S = 1$. First, we discuss the generation of the asymmetry and then the DM production.

4.1 Asymmetric dark matter viaogenesis

In the following, we assume a hierarchical scenario $M_{N_{2,3}} \gg M_{N_1}$, which in turn implies $y_\sigma^{2,3} \gg y_\sigma^1$. The simultaneous generation of lepton and dark sector asymmetries takes place at a high scale ($T \sim M_{N_1}$) via the decays of the lightest RHN, N_1 , into the two channels (the asymmetry generated in the decays of $N_{2,3}$ are washed out by N_1 interactions), as shown in figure 1:

1. $N_1 \rightarrow LH$ generates a lepton asymmetry, η_L , which is later reprocessed into a baryon asymmetry η_B by sphalerons. This is the case of Type-I thermal leptogenesis, well studied in the literature (see ref. [64] for a review), but with extra contributions, see below.
2. $N_1 \rightarrow \chi S$ generates an asymmetry in the dark sector, η_D , analogous to the lepton asymmetry.

Here, we focus on the regime where the asymmetry generated by N_1 decays into S is not washed out, and is comparable to the asymmetry in χ , i.e., $\eta_\chi \sim \eta_S \sim \eta_D$. In order to generate an asymmetry, CP needs to be violated. The CP asymmetry generated in the decays of N_1 can be written as

$$\varepsilon_L = \sum_\alpha \frac{\Gamma_{N_1 \rightarrow L_\alpha H} - \Gamma_{N_1 \rightarrow \bar{L}_\alpha H^\dagger}}{\Gamma_{N_1}}, \quad \varepsilon_\chi = \frac{\Gamma_{N_1 \rightarrow \chi S} - \Gamma_{N_1 \rightarrow \bar{\chi} S^\dagger}}{\Gamma_{N_1}}, \quad (4.2)$$

where $\Gamma_{N_1} = ((y_\nu y_\nu^\dagger)_{11} + |y_S^1|^2) M_{N_1} / (8\pi)$ is the total tree-level decay width of N_1 and y_S^1 is the relevant Yukawa coupling for $N_1 \rightarrow \chi + S$. In the following we refer to it as y_S .

However, due to CPT invariance, no asymmetry can be generated at the tree level. CP violation arises via the interference of tree-level and one-loop level decay (vertex and self-energy corrections) amplitudes, which depends on the imaginary part of the product of the Yukawas involved. Therefore, for this asymmetry to be non-zero, we require at least two distinct phases that come from the Yukawas y_ν and y_S , so the couplings to the

heavier neutrinos $N_{2,3}$ are important. Further, having an imaginary part in the internal loop contribution demands that the would-be decay products can be produced on-shell (i.e., the optical theorem), which is easily satisfied as we take $M_{N_1} \gg m_{\chi,S,L,H}$.

In this cogenesis scenario, it can be seen that both ε_L and ε_χ depend on y_ν and y_S , as the dark sector particles (χ, S) are involved in the one-loop self-energy correction for $N_1 \rightarrow LH$ and vice versa. Thus, the ratio of the decay asymmetries depends on the ratio of the couplings y_ν and y_S and may be correlated with the branching ratio of N_1 decay in each sector (Br_L and Br_χ). This relates neutrino mass generation to the baryon and DM abundances.

In ref. [28] it has been shown that the dark asymmetry η_D can be quite different from the visible one η_B (contrary to the ADM models that predict $\eta_B \sim \eta_D$) due to different branching ratios, washout effects and transfer between the sectors via inverse decays and $2 \leftrightarrow 2$ scatterings. The asymptotic asymmetries for the two sectors can be written as

$$\begin{aligned}\eta_L^\infty &= \varepsilon_L \xi_L Y_{N_1}^{\text{eq}}(T \ll M_{N_1}) \simeq 2.6 \times 10^{-10}, \\ \eta_\chi^\infty &= \varepsilon_\chi \xi_\chi Y_{N_1}^{\text{eq}}(T \ll M_{N_1}),\end{aligned}\tag{4.3}$$

where $\xi_L(\xi_\chi)$ is the leptonic (dark) efficiency parameter characterising the effects of washout and transfer interactions, and $Y_{N_1}^{\text{eq}}(T \ll M_{N_1}) = 135\zeta(3)/4\pi^4 g_*$ is the initial equilibrium N_1 yield, with $g_* \simeq 106.75$, the number of relativistic degrees of freedom. The numerical value of η_L^∞ is selected to match the observed baryon asymmetry, i.e., $\eta_B = (28/79)\eta_L$, generated via sphaleron processes. Note that the presence of $N - Z_{B-L}$ interactions in the model may modify the usual picture; however, the correct value of the asymmetry may be generated [50, 65].

Typically, producing the DM relic abundance in ADM models constrains the DM mass, depending on the ratio η_D/η_B . However, since χ is not the DM candidate in our set-up, it is not constrained to be of the order of few GeVs (for $\eta_D \sim \eta_B$) and thus may be much larger, $\sim \mathcal{O}(\text{TeV})$. In the analysis below, we consider dark asymmetries in the range $\eta_D \sim (0.1-1)\eta_B$, which may be achieved if $\text{Br}_L \sim \text{Br}_\chi$ and the Yukawa couplings of both sectors have similar hierarchies. Thus, we work with DM masses in the GeV ballpark. We refer the reader to figure 6 of ref. [28], where it is shown the order of the asymmetries that can be obtained for a given value of M_{N_1} and different branching ratios.

A key feature of any ADM model is the presence of an interaction that efficiently depletes the symmetric component. In our set-up, in order for the DM component ψ to be asymmetric, we require that the symmetric population of χ is annihilated and only the asymmetric component participates in the late decays into ψ . For example, this can take place via annihilations of the form $\bar{\chi}\chi \rightarrow Z_D Z_D$. In the non-relativistic limit, $s \simeq 4m_\chi^2$, the thermally-averaged cross section (considering s -wave annihilations) for $m_\chi > m_{Z_D}$ is given by

$$\sigma v(\bar{\chi}\chi \rightarrow Z_D Z_D) = \frac{g_D^4}{16\pi m_\chi^2} \left(1 - \frac{m_{Z_D}^2}{m_\chi^2}\right)^{3/2} \left(1 - \frac{m_{Z_D}^2}{2m_\chi^2}\right)^{-2}.\tag{4.4}$$

For $m_\chi > m_S$, an extra contribution comes from annihilations of the form $\bar{\chi}\chi \rightarrow Z_D \rightarrow S^\dagger S$, which opens up the region of parameter space where $m_\chi < m_{Z_D}$. It should be noted

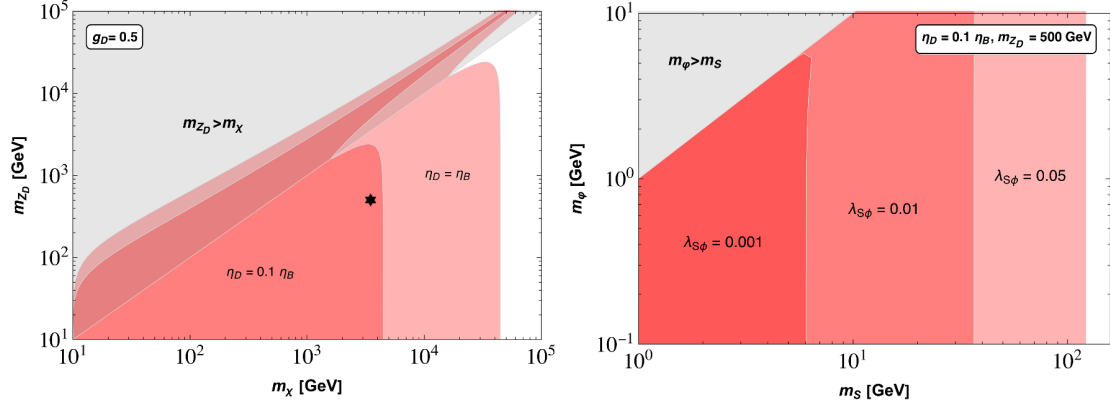


Figure 3. The red regions indicate the regions of m_χ, m_{Z_D} (left) and m_S, m_ϕ (right) where the annihilations of χ, S respectively are strong enough to result in a fractional asymmetry $r_\chi(S) < 10^{-2}$. *Left:* the star represents the benchmark point $\{m_\chi, m_{Z_D}\} = \{3.5 \text{ TeV}, 500 \text{ GeV}\}$. In the gray region the annihilation channel $\bar{\chi}\chi \rightarrow Z_D Z_D$ is closed. *Right:* in the gray region, annihilations to ϕ are closed.

that there is another contribution to annihilations from the channel $\bar{\chi}\chi \rightarrow Z_{B-L} \rightarrow \bar{q}q (\bar{l}l)$, where $q(l)$ is a SM quark (lepton). However, given the large mass of Z_{B-L} , this turns out to be negligible. In the left plot of figure 3 we show the mass ranges of Z_D and χ where the latter is asymmetric, i.e., $r_\chi < 10^{-2}$, see eq. (2.2). One observes how, for larger values of the asymmetry, a larger region of the parameter space has $r_\chi < 10^{-2}$. Note also the presence of the resonance for $m_{Z_D} \simeq 2m_\chi$.

4.2 Contribution of ψ

4.2.1 Production of ψ from χ decays

First, we study the dynamics of ψ , neglecting the contribution to the relic abundance from S . In this limit, χ can only decay to $\psi\phi(\varphi)$. We assume that the $\bar{\chi}\chi$ annihilation processes studied in the previous section are efficient enough so that $r_\chi < 10^{-2}$ and the freeze-out abundance of χ is determined by its asymmetry. The production of ψ is driven by decays of χ : $\chi \rightarrow \psi\phi$ for $T > v_\phi$ and $\chi \rightarrow \psi\varphi$ for $T < v_\phi$. As the $h - \varphi$ mixing is small ($\theta \ll 1$), it is safe to trade φ, h for the mass eigenstates. The production via decays can be divided into two types:

1. While χ is in thermal equilibrium with the SM bath ($T > T_*$, $T_* \simeq m_\chi/20$ being the freeze-out temperature of χ), the decays are symmetric, i.e., ψ and $\bar{\psi}$ are produced in equal amounts from the decays of χ and $\bar{\chi}$, that are symmetric during this period, i.e., $Y_\chi^+ \simeq Y_\chi^- = Y_\chi^{\text{eq}} \gg \eta_D$. This is the usual freeze-in contribution. We also refer to these processes as *early* decays. We denote the abundance of ψ particles produced by freeze-in by Y_{FI} . This symmetric production peaks around $T \sim m_\chi > v_\phi$, so that the channel is $\chi \rightarrow \psi\phi$.
2. Once χ freezes out ($T < T_*$), the population of χ becomes asymmetric, i.e., $Y_\chi^+ \simeq \eta_D \gg Y_\chi^- \simeq r_\chi Y_\chi^+$. Such an asymmetry is then subsequently transferred to ψ via

decays. This is the asymmetric freeze-in contribution of ψ from *late* decays (LD). Since the decays occur late, i.e. $T_D \ll v_\phi$, in this case the channel is $\chi \rightarrow \psi\phi$. The populations are given by

$$Y_{\text{LD}}^+ = \frac{\eta_D}{1 - r_\chi} \simeq \eta_D, \quad Y_{\text{LD}}^- = r_\chi Y_{\text{LD}}^+ \simeq \eta_D r_\chi, \quad (4.5)$$

where in the last step we used $r_\chi < 10^{-2}$.

Therefore, the asymptotic abundance of ψ and $\bar{\psi}$ can be written as

$$Y_\psi^+ = \frac{Y_{\text{FI}}}{2} + Y_{\text{LD}}^+ \simeq \frac{Y_{\text{FI}}}{2} + \eta_D, \quad Y_\psi^- = \frac{Y_{\text{FI}}}{2} + Y_{\text{LD}}^- \simeq \frac{Y_{\text{FI}}}{2} + \eta_D r_\chi, \quad (4.6)$$

where in the last step we used eq. (4.5). We show in the next section that eq. (4.6) gets modified by a multiplicative factor when the contribution of the scalar S is taken into account. The three interesting cases for the asymmetric ratio are

$$r_\psi \simeq \begin{cases} r_\chi & \text{if } \eta_D r_\chi \gg Y_{\text{FI}}, \\ Y_{\text{FI}}/(2\eta_D) & \text{if } \eta_D r_\chi \ll Y_{\text{FI}} \ll \eta_D, \\ 1 & \text{if } Y_{\text{FI}} \gg \eta_D. \end{cases} \quad (4.7)$$

We remind that, according to our definition in section 2, DM is asymmetric if $r_\psi < 10^{-2}$, partially asymmetric if $10^{-2} < r_\psi < 0.9$ and symmetric if $r_\psi > 0.9$. Hence, in order for ψ to be asymmetric, not only we need to require that $r_\chi \ll 1$ but also the symmetric freeze-in contribution to the abundance, $Y_{\text{FI}}/2$, should be suppressed, i.e., $Y_{\text{FI}} \ll \eta_D$. In the opposite regime, where the freeze-in production from early decays is dominant, the final ψ abundance is always symmetric. The decay width is given by

$$\Gamma_{\chi \rightarrow \psi\phi} \simeq \frac{y_\phi^2 m_\chi}{32\pi} \Delta^2(f_\psi, f_\phi), \quad (4.8)$$

where $f_\psi \equiv m_\psi/m_\chi$ and $f_\phi \equiv m_\phi/m_\chi$, and $\Delta(f_\psi, f_\phi)$ is the phase space suppression factor,

$$\Delta^2(f_\psi, f_\phi) = \left(1 - f_\psi^2 - f_\phi^2 + 2f_\psi\right)^2 \left[\left(1 - f_\psi^2 - f_\phi^2\right)^2 - 4f_\psi^2 f_\phi^2 \right], \quad (4.9)$$

with $0 \leq \Delta(f_\psi, f_\phi) \leq 1$. We have neglected the small corrections due to fermion mixing ϵ_f . There is an analogous expression for $\phi \rightarrow \psi$. Here onwards, we omit the arguments of the function Δ and we consider values of the parameters such that $m_\chi \gg m_\psi, m_\phi$. Within this approximation $\Delta \simeq 1$. Furthermore, in this limit the width is practically independent of m_ϕ or m_ψ , so that in the following we do not differentiate among decays into $\psi\phi$ and $\psi\phi$. Notice that the Yukawa y_ϕ also leads to annihilation processes (such as $\bar{\chi}\chi \rightarrow \bar{\chi}\psi$) which would contribute to ψ production. However, as long as $\Delta \simeq 1$, these are subleading with respect to decays.

The freeze-in contribution from early decays can be computed numerically by solving the Boltzmann equations for χ and ψ . However, a very good analytic estimate is given by

$$Y_{\text{FI}} \simeq \frac{135}{8\pi^4} \sqrt{\frac{45}{\pi g_*^3}} \frac{\Gamma_{\chi \rightarrow \psi\phi} M_{\text{Pl}}}{m_\chi^2} \simeq 6 \times 10^{-6} y_\phi^2 \Delta^2 \frac{M_{\text{Pl}}}{m_\chi}, \quad (4.10)$$

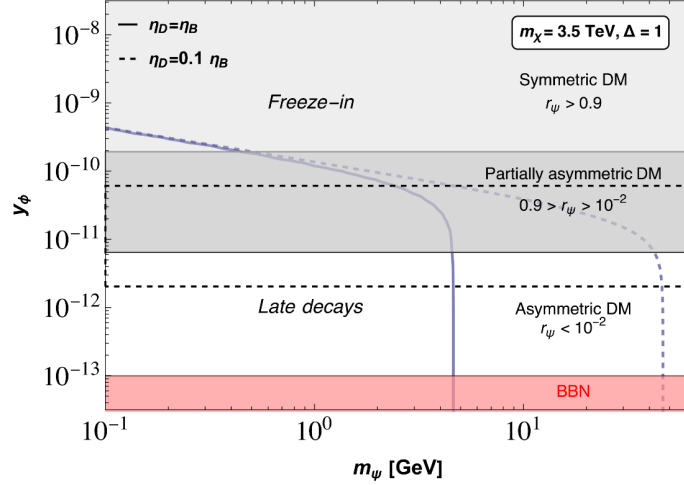


Figure 4. Parameter space for the scenario in which the DM is composed solely of ψ and we assume that S is light enough in each point of the plot so that its contribution to the DM abundance can be safely neglected. We distinguish three different regions depending on the nature of DM (symmetric, asymmetric or partially asymmetric), labelled by the value of r_ψ . In the darker gray region DM is produced through freeze-in and is symmetric. In the white region it is produced via late decays and is asymmetric. The DM relic abundance is reproduced along the blue solid (dashed) line which corresponds to $\eta_D/\eta_B = 1$ (0.1), whereas the horizontal dashed lines indicate the shift in the regions for $\eta_D = 0.1\eta_B$.

where we used that the production peaks around the mass of the heaviest particle involved in the decay, i.e., at $T \simeq m_\chi$. The late decays of χ would instead peak at temperature T_D at which $\Gamma_{\chi \rightarrow \psi\varphi}/H|_{T=T_D} \simeq 1$,

$$T_D \approx y_\phi \Delta \sqrt{\frac{m_\chi M_{\text{Pl}}}{32\pi}} \approx 10 \text{ MeV} \Delta \left(\frac{y_\phi}{10^{-12}} \right) \left(\frac{m_\chi}{3.5 \text{ TeV}} \right)^{1/2}. \quad (4.11)$$

In order to get asymmetric DM, we impose the condition $Y_{\text{FI}} \lesssim 10^{-2}\eta_D$ (here we are assuming values of $\{g_D, m_\chi, m_{Z_D}\}$ such that $r_\chi \lesssim 10^{-2}$) which implies

$$y_\phi \Delta \lesssim 6 \times 10^{-12} \left(\frac{\eta_D}{\eta_B} \right)^{1/2} \left(\frac{m_\chi}{3.5 \text{ TeV}} \right)^{1/2}. \quad (4.12)$$

In figure 4, we show the parameter space in the plane y_ϕ versus m_ψ where the contribution of ψ is dominant. We fix $m_\chi = 3.5 \text{ TeV}$. We highlight the regions where it is symmetric, partially-asymmetric and asymmetric, subject to the constraints discussed below.

4.2.2 Constraints

The gauge coupling g_X is constrained by long-range force experiments, $g_X \lesssim 10^{-8}$, see appendix B.2. In the rest of the paper, we take g_X sufficiently small so that gauge interactions can not drive thermalisation of ψ and therefore g_X never plays a role. On the other side, ψ could thermalise through processes involving the Yukawa coupling y_ϕ . All these processes involve at least one χ particle (or the conjugate). The most important

ones are the decays (and inverse decays) $\chi \rightarrow \psi\varphi$. The condition for non-thermalisation is $\Gamma_{\chi \rightarrow \psi\varphi}/H|_{T \simeq m_\chi} \ll 1$. Indeed, for $T > m_\chi$, the interaction rate/Hubble ratio grows when the temperature decreases, while for $T < m_\chi$ decay and inverse decays become inefficient. Therefore, we evaluate the ratio at its maximal value, i.e., $T \simeq m_\chi$, which gives

$$y_\phi \Delta \ll 30 \sqrt{\frac{m_\chi}{M_{\text{Pl}}}} = 5 \times 10^{-7} \sqrt{\frac{m_\chi}{3.5 \text{ TeV}}}. \quad (4.13)$$

Notice that taking the limit $\Delta \ll 1$ does not help in pushing towards higher values of the Yukawa. Indeed, if Δ is small, the decay channel closes and annihilations become more relevant. Annihilations processes are independent from Δ and give a condition analogous to eq. (4.13) with $\Delta = 1$. However, as we are considering $\Delta \simeq 1$, decays are more important than annihilations.

In the case in which ψ is produced mainly by late decays, we require that these are peaked much after the freeze-out of χ , i.e., $T_D \ll T_*$, which ensures that the population of ψ is asymmetric. This gives a stronger condition,

$$y_\phi \Delta \ll 0.5 \sqrt{\frac{m_\chi}{M_{\text{Pl}}}} = 8.5 \times 10^{-9} \sqrt{\frac{m_\chi}{3.5 \text{ TeV}}}. \quad (4.14)$$

If ψ is produced by early decays, peaked around $T_{\text{prod}} \simeq m_\chi$, it behaves as cold DM as long as $m_\psi \gtrsim \text{keV}$. However, if the production from late decays is significant, we must take into account an additional constraint. Indeed, as we have a heavy particle (χ) decaying late into a lighter stable one (ψ), we must check that the DM free streaming length is smaller than 0.1 Mpc, which gives [66]

$$m_\psi > 3.5 \text{ keV} \langle p/T \rangle_{\text{prod}} \left(\frac{10}{g_*(T_D)} \right)^{1/3} \gtrsim 1.7 \text{ keV} \frac{m_\chi \Delta}{T_D} \left(\frac{10}{g_*(T_D)} \right)^{1/3}, \quad (4.15)$$

where in the second step we used that the typical DM momentum at production is $\langle p \rangle \sim m_\chi \Delta/2$, while $\langle T_{\text{prod}} \rangle = T_D$. Using eqs. (4.11) and (4.15), this translates into a lower bound on the coupling,

$$y_\phi \gtrsim 5.6 \times 10^{-14} \left(\frac{m_\chi}{3.5 \text{ TeV}} \right)^{1/2} \left(\frac{5 \text{ GeV}}{m_\psi} \right) \left(\frac{10}{g_*(T_D)} \right)^{1/3}. \quad (4.16)$$

This bound only applies if *i*) ψ reproduces the DM relic abundance, and *ii*) the production from late decays is dominant, i.e., $Y_{\text{LD}}^+ \gg Y_{\text{FI}}$, corresponding to the vertical blue lines of figure 4. We can use BBN bounds to constrain the lifetime of χ . If χ did not decay, its abundance today, $\Omega_\chi h^2$, would be m_χ/m_ψ times the would-be DM (ψ) abundance. In particular, for $m_\chi \sim \mathcal{O}(\text{TeV})$ and would-be abundance $\eta_D \sim 10^{-11}$, this leads to $\tau_\chi \lesssim (0.1-1) \text{ s}$ [67] for

$$y_\phi \Delta \gtrsim 10^{-13} \left(\frac{3.5 \text{ TeV}}{m_\chi} \right)^{1/2}. \quad (4.17)$$

Here we assumed that the decay of φ into SM radiation (mainly hadrons) occurs instantaneously right after χ decays; otherwise, the bound should read $\tau_\chi + \tau_\varphi \lesssim (0.1-1) \text{ s}$. This is discussed in appendix B.1. The effects of fermion mixing are discussed in appendix B.3.

4.3 Contribution of S

4.3.1 Production of S from freeze-out

The scalar S is produced by N_1 decays at high temperature and shares the same asymmetry of χ , i.e. η_D . Once produced, it thermalises with the SM bath through scalar and gauge interactions and undergoes annihilations. In the following, we focus mainly on masses of S in the 1–50 GeV range, so that its contribution to the relic abundance may be of the same order as that of ψ . In this range of masses, annihilations of S into Z_D may be kinematically forbidden. However, if $m_\varphi < m_S$ (possible by tuning the coefficients of the scalar potential, e.g., $m_\varphi/v_\phi \lesssim 10^{-3}$ and $m_\varphi \sim \mathcal{O}(\text{GeV})$ or lighter), then the annihilations $S^\dagger S \rightarrow \varphi\varphi$ are allowed. Next we consider the minimal option of just the scalar portal $\lambda_{S\phi}|S|^2|\phi|^2$. The non-relativistic cross section for $S^\dagger S \rightarrow \varphi\varphi$ induced by the operator is given by⁷

$$\sigma v(S^\dagger S \rightarrow \varphi\varphi) \simeq \frac{\lambda_{S\phi}^2}{32\pi m_S^2} \left(1 - \frac{m_\varphi^2}{m_S^2}\right)^{1/2} \left(\frac{1 - m_\varphi^2/2m_S^2 - 2\lambda_{S\phi}v_\phi^2/m_S^2}{1 - m_\varphi^2/2m_S^2}\right)^2. \quad (4.18)$$

Even for moderately small values of the coupling, the cross section is significant and annihilations are strong enough to destroy the symmetric population of S . Hence, the symmetric population of S annihilates around $T_*^{(S)} \sim m_S/20$, leaving only the asymmetric population, with the abundance fixed by the dark asymmetry $Y_S^+ \sim \eta_D$, $Y_S^- \sim r_S Y_S \ll Y_S^+$, completely analogous to the computation of χ annihilations. In figure 3 (right panel) we show the region of m_φ and m_S where S is asymmetric for different values of the coupling $\lambda_{S\phi}$ fixing $\eta_D = 0.1\eta_B$ and $m_{Z_D} \sim 500$ GeV.

In principle, other annihilation channels for S are possible, depending on the scalar potential parameters, such as $S^\dagger S \rightarrow h \rightarrow \bar{f}f$ (which is more suppressed for $m_h > m_S$), allowing for a larger set of possibilities. For simplicity, we focus only on annihilations into $\varphi\varphi$, which involve only one coupling $\lambda_{S\phi}$ and need not be very large, $\mathcal{O}(10^{-2})$.

4.3.2 Production of S from χ decays

Integrating out the lightest of the heavy sterile neutrinos, N_1 , induces the effective interactions

$$y_\nu y_S \frac{\bar{L}\tilde{H}S\chi}{M_{N_1}}, \quad y_S^2 \frac{\chi^2 S^2}{M_{N_1}}. \quad (4.19)$$

The former leads to $\chi \rightarrow S^\dagger \nu_L$ and $\chi \rightarrow S^\dagger \nu_L h$ decays, which compete with $\chi \rightarrow \psi\varphi$, eventually populating the $S^\dagger$ sector, whereas, the latter induces $\chi\chi \rightarrow S^\dagger S^\dagger$ annihilations. Focusing on decays, we can use the estimate (see also ref. [63])

$$\Gamma(\chi \rightarrow S^\dagger \nu) \approx \frac{|y_S|^2 m_\chi}{32\pi} \left(\frac{m_\nu}{M_{N_1}}\right) \left(1 - \frac{m_S^2}{m_\chi^2}\right)^2. \quad (4.20)$$

⁷We neglect the diagram containing the self-interaction φ^3 , which enters if the ϕ quartic coupling is not extremely small.

These decays are peaked around $T_D^{(S)}$ defined as $\Gamma_{\chi \rightarrow S^\dagger \nu} / H|_{T=T_D^{(S)}} = 1$ (analogous to the decays into ψ computed earlier) and given by

$$T_D^{(S)} \simeq |y_S| \sqrt{\frac{m_\nu m_\chi M_{\text{Pl}}}{32\pi M_{N_1}}} \simeq \text{MeV} \left(\frac{|y_S|}{10^{-3}} \right) \left(\frac{m_\nu}{0.05 \text{ eV}} \frac{m_\chi}{3.5 \text{ TeV}} \frac{10^{11} \text{ GeV}}{M_{N_1}} \right)^{1/2}. \quad (4.21)$$

Since for the values of the parameters $T_D^{(S)} < T_*$, the only important decay is $\chi \rightarrow S^\dagger \nu_L$, while the conjugate process is irrelevant as the population of $\bar{\chi}$ after freeze-out is negligible (recall that we are in the region of the parameter space in which $r_\chi < 10^{-2}$).⁸

We can parametrise the dominant decay channel of χ by defining the ratio of branching ratios

$$R \equiv \frac{\text{Br}(\chi \rightarrow S^\dagger \nu)}{\text{Br}(\chi \rightarrow \psi \varphi)} \sim \frac{|y_S|^2}{y_\phi^2} \frac{m_\nu}{M_{N_1}}, \quad (4.22)$$

where in the last step we used eqs. (4.8) and (4.20). The decay of χ into ψ is the dominant channel, i.e., $R \ll 1$, for

$$|y_S| \ll \left(\frac{y_\phi}{2 \times 10^{-11}} \right) \left(\frac{M_{N_1}}{10^{11} \text{ GeV}} \right)^{1/2} \left(\frac{0.05 \text{ eV}}{m_\nu} \right)^{1/2}. \quad (4.23)$$

Notice that the contribution of the 3-body decays $\chi \rightarrow \nu h S^\dagger$ should be similar to the 2-body ones, as the decay rate gets suppressed by the phase space factor while at the same time it has an enhancement $(m_\chi/v)^2 \sim 200$ (eventually it can become dangerous for $m_\chi \gg \text{TeV}$). For a generic value of R , the probability to decay into ψ is $1/(1+R)$ while into $S^\dagger$ is $R/(1+R)$. Therefore the abundances of $\psi, \bar{\psi}$ in eq. (4.6) get divided by $(1+R)$.

Concerning S , we can distinguish two possibilities:

1. If $T_D^{(S)} < T_*^{(S)} < T_*$, at the time where decays into $S^\dagger$ peak, the latter has already decoupled from the thermal bath with an asymmetric abundance $Y_S = \eta_D$ and therefore the S and $S^\dagger$ population can not annihilate. Then the abundances of $S, S^\dagger$ are

$$Y_S \equiv Y_S^+ = \eta_D, \quad Y_{S^\dagger} \equiv Y_S^- = \frac{R}{1+R} \eta_D, \quad (4.24)$$

where we assume that $r_S < 10^{-2}$ and we have ignored it for simplicity. This corresponds to

$$|y_S| \lesssim 0.1 \left(\frac{m_S}{\text{GeV}} \right) \left(\frac{M_{N_1}}{10^{11} \text{ GeV}} \right)^{1/2} \left(\frac{0.05 \text{ eV}}{m_\nu} \right)^{1/2} \left(\frac{3.5 \text{ TeV}}{m_\chi} \right)^{1/2}. \quad (4.25)$$

2. If $T_*^{(S)} < T_D^{(S)} < T_*$, the decays produce a population of $S^\dagger$, while $S^\dagger S$ annihilations are still efficient. As a result there is a partial washout of η_S , which gets reduced to $\eta_D/(1+R)$. Therefore, at $T < T_*^{(S)}$ a population of S decouples with abundance

$$Y_S^+ = \frac{1}{1+R} \eta_D, \quad Y_S^- \ll Y_S^+, \quad (4.26)$$

⁸Also in this case we can distinguish between the late decays, peaked at $T_D^{(S)}$, and the early decays at $T > T_*$. The latter produce a symmetric population of S and $S^\dagger$ from the decays of χ and $\bar{\chi}$ (the analogue of the freeze-in population of ψ , peaked at $T \simeq m_\chi$). However, this symmetric population thermalises and undergoes annihilations leaving no imprint in the final abundance.

where we assume $r_S < 10^{-2}$. Notice that for $R \ll 1$ eqs. (4.24) and (4.26) are equivalent as the decays are irrelevant.

4.3.3 Constraints

We can constrain the value of $|y_S|$ for the case in which $R > 1$. As for this choice χ decays mostly into $S^\dagger$, we must impose that the decay occurs before BBN, in analogy with section 4.2.2. Using the decay rate in eq. (4.20) we find that the BBN bound translates into the constraint

$$|y_S| \gtrsim 10^{-3} \left(\frac{0.05 \text{ eV}}{m_\nu} \frac{M_{N_1}}{10^{11} \text{ GeV}} \frac{3.5 \text{ TeV}}{m_\chi} \right)^{1/2}. \quad (4.27)$$

For $R < 1$ we can choose smaller y_S while BBN constrains the value of y_ϕ , see eq. (4.17).

We also find that for $R > 1$ the model is characterised by an interesting feature: the decays $\chi \rightarrow S^\dagger \nu_L$, which occur before BBN and neutrino decoupling (around $T \sim \mathcal{O}(10 \text{ MeV})$ for the choice of parameters we adopted and $|y_S| \sim 10^{-2}$), generate also an asymmetric population of neutrinos,

$$\Delta\eta_\nu = \frac{R}{(1+R)} \eta_D, \quad (4.28)$$

which is maximal ($\approx \eta_D$) for large R . Therefore, the neutrino population is more asymmetric than in the standard case, as this new contribution sums up to the usual leptonic asymmetry generated earlier on by leptogenesis. Note, however, that since these decay processes occur below the scale of electroweak symmetry breaking, this leptonic asymmetry is not transferred to the baryonic one.

Finally, if $R > 1$ (and $T_D^{(S)} < T_*^{(S)}$), the $S^\dagger$ population that arises from late decays of χ may be warm. In such a case, if a significant fraction of the DM was made by this $S^\dagger$ population, constraints from free streaming length (see eq. (4.15) with $m_\psi \rightarrow m_S$ and $T_D \rightarrow T_D^{(S)}$) would give $|y_S| \gtrsim 5 \text{ MeV}/m_S$. However, even in the cases in which the contribution of ψ DM is negligible, the DM is composed by a mixture of S (produced by freeze-out, always cold) and $S^\dagger$, where $Y_S^+ \geq Y_S^-$ (the equality applies if $R \gg 1$). This corresponds to a mixture of cold/warm DM, which in general is less constrained than the warm DM case [68].

In the following, we discuss the scenarios arising due to the different dynamics of the dark components in the model. It can be checked that eq. (4.1) is always fulfilled.

5 Dark matter relic abundance

In this section, we discuss the possible scenarios that can be distinguished depending on the values of the parameters y_ϕ and R (which can be rephrased as a function of y_ϕ and $|y_S|$), while satisfying all the constraints. We restrict our analysis to the region of the parameter space in which:

1. Both ψ and S are stable, see the discussion in section 6.

2. $\bar{\chi}\chi \rightarrow Z_D Z_D$ annihilations are efficient enough so that $r_\chi < 10^{-2}$, i.e., the late decays of χ are always asymmetric.
3. $S^\dagger S \rightarrow \varphi\varphi$ annihilations are efficient enough so that $r_S < 10^{-2}$, i.e., only the asymmetric population of S survives the annihilations.

Regarding the Yukawa parameter y_ϕ , we consider the range $10^{-13} \lesssim y_\phi \lesssim 5 \times 10^{-7}$, where the lower bound comes from BBN (only applicable if $R < 1$), while the upper one comes from the requirement that ψ does not thermalise. Regarding the DM component ψ , we identify its nature depending on the value of y_ϕ . In the two limiting cases, we have

- *Asymmetric ψ* : if $y_\phi \lesssim 6 \times 10^{-12} \sqrt{\eta_D/\eta_B}$, the freeze-in population of ψ from early decays is negligible, while the late decays are dominant, resulting in an asymmetric population $Y_\psi = \eta_D \gg Y_{\text{FI}}$.
- *Symmetric ψ* : if $y_\phi \gtrsim 2 \times 10^{-10} \sqrt{\eta_D/\eta_B}$, the ψ population is dominated by freeze-in (early decays), $Y_{\text{FI}} \gg \eta_D$. The resulting ψ population is symmetric.

At the same time, we identify the nature of the DM component S according to the value of the parameter R . The two limiting cases are:

- *Asymmetric S* : if $R \ll 1$ (corresponding to $|y_S| \ll y_\phi/(2 \times 10^{11})$), χ decays into $S^\dagger$ are negligible and the abundance of S is determined by the asymmetry via freeze-out, i.e., $Y_S = \eta_D$.
- *Symmetric S* : if $R \gtrsim \mathcal{O}(10)$, χ mostly decays into $S^\dagger$, so that the final abundance of $S, S^\dagger$ population is the same, $Y_S = Y_{S^\dagger} = \eta_D$ and DM is symmetric. The interesting feature of this scenario is that the DM abundance is set by the asymmetry η_D even though the nature of DM is symmetric. Furthermore, given the large value of R , the production of ψ is negligible. Therefore the scenario leads to practically only 1 DM component with abundance completely fixed by the asymmetry: $Y_S + Y_{S^\dagger} = 2\eta_D$, leading to the prediction

$$m_S \simeq 2.5 \text{ GeV} \left(\frac{\eta_B}{\eta_D} \right). \quad (5.1)$$

This scenario (scenario 6 according to our classification in table 2) is phenomenologically interesting as it leads to an enhanced indirect detection signal compared to the usual freeze-out case.

The intermediate regimes give rise to new possibilities, such as partially-asymmetric DM for the different components. The classification of scenarios depending on the nature of the ψ/S population and the values of y_ϕ and R is shown in table 2. The scenarios are discussed in detail in appendix A, where we also provide the classification on the basis of dominant production mechanism and asymptotic nature of both dark matter components in table 4. For all the scenarios, one can use the following expressions for the individual and total relic abundance,

$$\frac{\Omega_\psi}{\Omega_S} = \frac{m_\psi(\eta_D + Y_{\text{FI}})}{\eta_D m_S f(R)}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} = \frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D m_S f(R)}{\eta_B(1+R)m_p}, \quad (5.2)$$

Sc.	ψ population	S population	$10^{-10} y_\phi / \sqrt{\eta_D / \eta_B}$	R	$T_D^{(S)} / T_*^{(S)}$
1	Asymmetric	Asymmetric	≤ 0.06	$\ll 1$	Any
2	Asymmetric	Partially Asymmetric	≤ 0.06	$\mathcal{O}(1)$	< 1
1-2	Asymmetric	Asymmetric	≤ 0.06	$\mathcal{O}(1)$	> 1
3	Partially Asymmetric	Asymmetric	$0.06 - 2$	$\ll 1$	Any
4	Partially Asymmetric	Partially Asymmetric	$0.06 - 2$	$\mathcal{O}(1)$	< 1
3-4	Partially Asymmetric	Asymmetric	$0.06 - 2$	$\mathcal{O}(1)$	> 1
5	Symmetric	Asymmetric	$\gtrsim 2$	$\ll 1$	Any
6	Negligible	Symmetric	$y_\phi \lesssim 5 \times 10^{-7}$	$\gtrsim \mathcal{O}(10)$	< 1

Table 2. Classification of the different scenarios depending on the nature of ψ/S population and the values of y_ϕ (in units of $\sqrt{\eta_D / \eta_B}$) and R . The values of the parameters are fixed: $m_\chi = 3.5$ TeV, $m_{Z_D} = 500$ GeV, $g_D = 0.5$, $M_{N_1} = 10^{11}$ GeV, and $m_\nu = 0.05$ eV. We take $y_\phi < 5 \times 10^{-7}$ to avoid thermalisation of ψ . In Scenarios 2, 4 and 6 the Yukawa satisfies $|y_S| \lesssim 0.1(m_S/\text{GeV})$. In the Mixed Scenarios 1-2 and 3-4 we have $|y_S| > 0.1(m_S/\text{GeV})$. Scenarios 1, 3 and 5 give the same result independently of this condition.

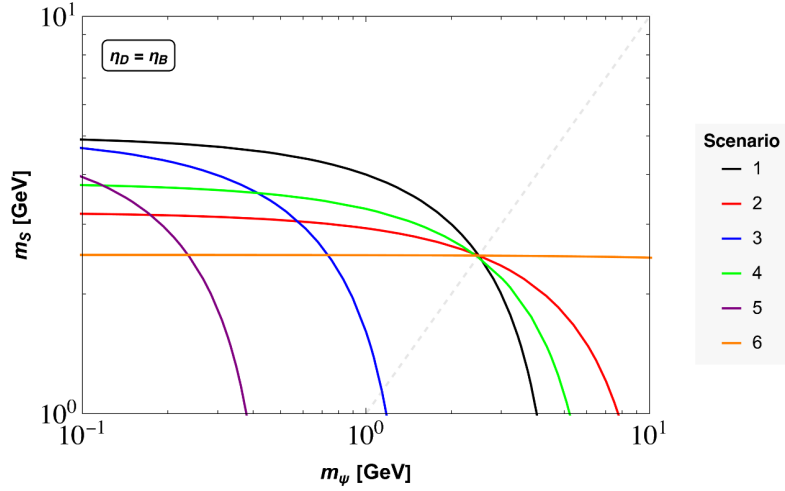


Figure 5. Contours of DM relic abundance in the m_ψ - m_S plane corresponding to $\eta_D = \eta_B$ for the different scenarios discussed above. We fix the values of the yukawa couplings ($y_\phi, |y_S|$) for each scenario: *Scenario 1*: $(10^{-12}, 10^{-3})$, *Scenario 2*: $(10^{-12}, 5 \times 10^{-2})$, *Scenario 3*: $(10^{-10}, 10^{-3})$, *Scenario 4*: $(6.4 \times 10^{-12}, 10^{-1})$, *Scenario 5*: $(2 \times 10^{-10}, 2 \times 10^{-4})$, *Scenario 6*: $(10^{-13}, 5 \times 10^{-2})$. The gray dashed line corresponds to $m_\psi = m_S$.

where

$$f(R) \equiv \begin{cases} 1 + 2R & \text{if } T_D^{(S)} < T_*^{(S)} \\ 1 & \text{if } T_D^{(S)} > T_*^{(S)}. \end{cases} \quad (5.3)$$

In figure 5 we show contours of correct relic abundance in the m_S versus m_ψ plane for the different scenarios. For definiteness, we consider the following benchmark point for all our plots: $m_\chi = 3.5$ TeV, $m_{Z_D} = 500$ GeV, $g_D = 0.5$, $M_{N_1} = 10^{11}$ GeV and $m_\nu = 0.05$ eV. We take $\eta_D = \eta_B$ and for each scenario we fix an appropriate value for the Yukawa couplings $|y_S|$ and y_ϕ (Scenarios 1, 2 and 5 are also shown in figures 8 and 9, respectively). DM could

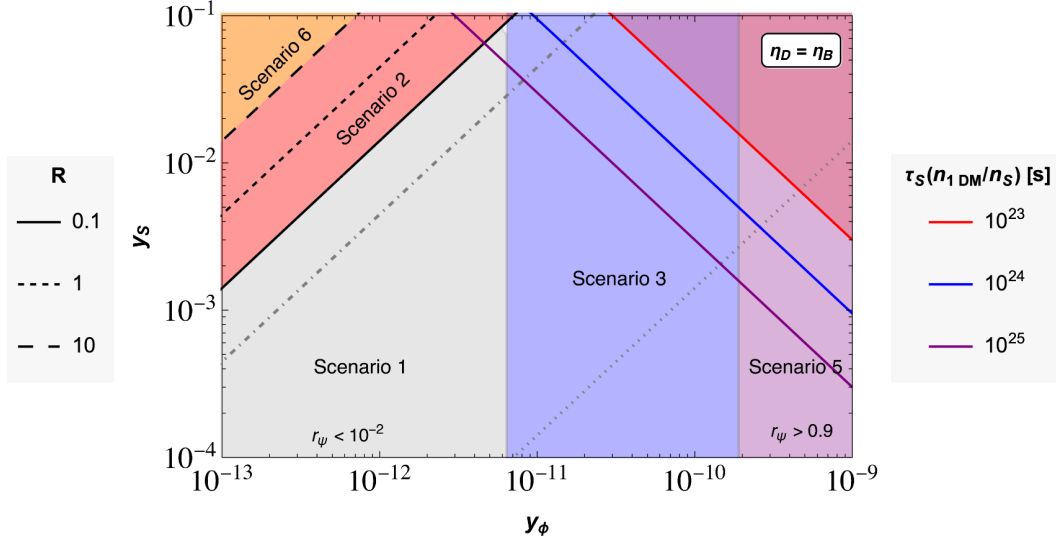


Figure 6. Parameter space of the scenarios discussed above. We fix $m_\chi = 3.5$ TeV, $m_{Z_D} = 500$ GeV, $g_D = 0.5$, $M_{N_1} = 10^{11}$ GeV, $m_\nu = 0.05$ and $\eta_D = \eta_B$. Scenario 4 is not visible in the plot but would appear in the intersection between Scenario 2 (red region) and Scenario 3 (blue region). Depending on the value of m_S , a portion of the region in Scenario 2 (3) could convert into the Mixed Scenario 1-2 (3-4); however, for $m_S \gtrsim \text{GeV}$, this requires $|y_S| \gtrsim 0.1$. The gray dot-dashed (dotted) line corresponds to $R = 1$ for $M_{N_1} = 10^9$ (10^4) GeV. Above the red (blue) [purple] line, the heavier between S and ψ , for fixed mass $\max(m_S, m_\psi) = 3$ GeV, has a lifetime, re-scaled by its relative number density, of 10^{23} (10^{24}) [10^{25}] s. Therefore, the region above the red line is excluded, see section 6.

be mainly composed by ψ , with m_ψ in between hundreds of MeV and tens of GeV, mainly by the scalar S with $m_S \sim \text{GeV}$, or by a combination of them (both GeV-ish). For fixed R , ψ DM is heavier when asymmetric (Scenario 1) and lighter when symmetric (Scenario 5). For fixed r_ψ , ψ gets heavier as R gets larger. On the contrary, S DM gets lighter while R grows. The minimal value of m_S is fixed by eq. (5.1), corresponding to Scenario 6. Notice the presence of a four-fold degeneracy between scenarios 1, 2, 4, 6. This corresponds to the case $m_S = m_\psi$, in which the mass of both the DM particles is fixed by eq. (5.1) (the relation is exact for Scenarios 1, 2 and 6, while for Scenario 4 it is approximately valid if $Y_{\text{FI}} < \eta_D$). Choosing another value for η_D leads to a rescaling of DM masses by a factor η_B/η_D .

The main results of the paper are provided in table 2 and figure 6 where we show the requirements on y_ϕ and R for each scenario, and show the allowed parameter space in the plane $|y_S|$ versus y_ϕ , respectively. We restrict our analysis to the region $10^{-13} \lesssim y_\phi \lesssim 5 \times 10^{-7}$ (no thermalisation of ψ and BBN bound on y_ϕ fulfilled for $R < 1$). The Yukawa $|y_S|$ is large enough to generate a sizeable dark asymmetry (and respects the BBN bound in the region in which $R > 1$, i.e., $|y_S| \gtrsim 10^{-3}$). The gray dot-dashed (dotted) line corresponds to $R = 1$ for $M_{N_1} = 10^9$ (10^4) GeV. Clearly, as N_1 gets lighter, χ preferably decays into $S^\dagger$. Notice that in the plot the masses of the DM particles are not fixed, but at every point there are always some values of the latter for which the correct DM relic abundance is reproduced.

6 Phenomenological signals

In principle, the models considered in this work may be difficult to test and disentangle in their current version, because of several reasons:

- i)* Freeze-in scenarios invoke very small couplings, $g_X \ll 1, y_\phi \ll 1$;
- ii)* The considered scale of $B-L$ breaking is very large, $v_{B-L} \gg v_{EW}$, so collider searches are not an option;
- iii)* Asymmetric DM yields suppressed indirect detection signals in general. Moreover, in our scenarios, the symmetric component is typically erased into the dark sector, via $\bar{\chi}\chi \rightarrow Z_D Z_D$ and $S^\dagger S \rightarrow \varphi\varphi$, see figure 3.
- iv)* The mixing of φ with the Higgs, generated by $\lambda_{H\phi} H^\dagger H \phi^\dagger \phi$, was taken to be very small.

However, there are a few distinctive signals of S , through the usual Higgs portal, $\lambda_{HS} H^\dagger H S^\dagger S$:

- Direct detection in the case in which the DM is mainly composed of S , which currently sets the limits $\lambda_{HS} \lesssim 0.01$ [69].
- Higgs invisible decays, for $m_S < m_h/2$, which currently sets the limits $\lambda_{HS} \lesssim 0.01$ [69, 70]. In this case, it could be that S was produced via Higgs decays, but it did not constitute a dominant part of the DM.
- Similarly, indirect detection signals from annihilations are present if the final abundance is composed of S and partially-asymmetric (Scenarios 2 and 4) or symmetric (Scenario 6). In this case, on-shell S -annihilations into muons, pions and electrons are possible. Typically, such light thermal DM is severely constrained by its energy injection in the CMB. In particular, Scenario 6 is characterized by an enhanced indirect detection signal due to larger than usual annihilation rates. Furthermore, it is very predictive, with a DM mass of 2.5 GeV, see eq. (5.1). Therefore, it may be interesting to further investigate for which values of λ_{HS} is Scenario 6 allowed.

Scenario 6 is also interesting as the $S^\dagger$ population arising from late decays of χ , results in a mixture of cold/warm DM, where S particles coming from the thermal plasma, represent the cold component, while their anti-particles, coming from late decays, the warm one, with a possible impact on structure formation [28]. Furthermore, the additional contribution to the asymmetric background neutrino population is maximal in this case, $\Delta Y_\nu \approx \eta_D$.

Last but not least, there is a very interesting phenomenological signal of our model: the presence of a monochromatic flux of neutrinos coming from the late decay of the heaviest of the two DM components. Let us assume $m_S > m_\psi$ (the opposite case $m_\psi > m_S$ is

completely analogous). At low energies, $E \ll m_\chi \ll M_{N_1}$, the decay $S \rightarrow \bar{\psi} + \nu_L$ is generated by the dimension-6 operator

$$\mathcal{O}_6 = \bar{L}\tilde{H}S\phi^\dagger\psi. \quad (6.1)$$

This operator arises by first integrating the right-handed neutrino field N_R , which generates the interaction given in eq. (4.19), and then at the lower scale the fermion χ , giving rise to the interaction⁹

$$\frac{y_S y_\phi y_\nu}{M_{N_1} m_\chi} \mathcal{O}_6. \quad (6.2)$$

Once H and ϕ acquire vevs, the decay $S \rightarrow \bar{\psi} + \nu_L$ is generated. The decay width reads

$$\Gamma(S \rightarrow \bar{\psi} + \nu_L) \approx \frac{|y_S|^2 y_\phi^2 m_S}{32\pi} \left(\frac{v_\phi}{m_\chi} \right)^2 \left(\frac{m_\nu}{M_{N_1}} \right) \left(1 - \frac{m_\psi^2}{m_S^2} \right). \quad (6.3)$$

To guarantee that both S and ψ are cosmologically stable and contribute to the DM abundance, the lifetime of the heavier particle needs to be larger than the age of the Universe, $\tau_S > t_U \sim 4 \times 10^{17}$ s. However, there are stronger constraints if the decay products include an active neutrino, $\tau_S \gtrsim 10^{23}$ s (for a GeV-ish single-component DM) [59, 60]. If S decays at late times, it leads to a very distinctive signature: a neutrino line peaked at $m_S/2 \sim \mathcal{O}(\text{GeV})$. Therefore, depending on the relative abundance of ψ and S , some region of the parameter space could be excluded, see refs. [61–63]. To quantify this, we compare the experimental bound with the re-scaled lifetime of the particle, $\tau_S^{\text{re-sc}} = \tau_S (n_{\text{1DM}}/n_S) > 10^{23}$ s, where n_{1DM} (n_S) is the single-component DM (S) number density. This constraint gives a condition on the parameters,

$$y_\phi |y_S| \lesssim 3 \times 10^{-12} \left(\frac{3 \text{ GeV}}{m_S} \right)^{1/2} \left(\frac{700 \text{ GeV}}{v_\phi} \right) \left(\frac{m_\chi}{3.5 \text{ TeV}} \right) \left(\frac{0.05 \text{ eV}}{m_\nu} \right)^{1/2} \left(\frac{M_{N_1}}{10^{11} \text{ GeV}} \right)^{1/2} \left(\frac{\Omega_{\text{DM}}}{\Omega_S(y_\phi, |y_S|, m_S, \eta_D)} \right)^{1/2}, \quad (6.4)$$

valid if $m_S > m_\psi$. In general the condition is not linear in the parameters. If $\Omega_S/\Omega_{\text{DM}} \ll 1$, the constraint is extremely weak: DM is made only by ψ and there are no decays. On the contrary, if $\Omega_S/\Omega_{\text{DM}} \sim \mathcal{O}(1)$, the condition becomes linear in the Yukawas. In the opposite regime, $m_\psi > m_S$, there is an equivalent condition to that of eq. (6.4), with $m_S \rightarrow m_\psi$ and $\Omega_S \rightarrow \Omega_\psi$. If the two states are almost degenerate $m_S \simeq m_\psi$ there is a strong phase space suppression so that the constraint gets weaker.

Notice that, even taking into account that the lifetime is re-scaled by the relative number density, this constraint is much stronger than the one coming from $\tau_S > t_U$, so that it automatically guarantees the stability of both S and ψ .

The neutrino line emerging from $S \rightarrow \bar{\psi}\nu_L$ ($\psi \rightarrow S^\dagger\nu_L$) decay is the main smoking gun of our scenarios. As an illustrative example, in figure 6 we show three different lines

⁹Operator $\mathcal{O}_6$ may also be generated by a UV completion of our model at scales above v_{B-L} . This produces additional constraints, weaker than eq. (6.4). We discuss them in appendix C.

Field	Spin	U(1) _{B-L}	U(1) _D	U(1) _X
S_L	1/2	0	0	0
σ'	0	+1	0	0

Table 3. Quantum numbers of the new states in the inverse seesaw variant.

in red, blue and purple corresponding respectively to $\tau_S^{\text{re-sc}} = 10^{23}, 10^{24}, 10^{25}$ s, for fixed $\max(m_S, m_\psi) = 3$ GeV. For this choice of mass, as we can see in figure 5, $\Omega_S/\Omega_{\text{DM}} \sim \Omega_\psi/\Omega_{\text{DM}} \sim \mathcal{O}(1)$; therefore, eq. (6.4) gives a simple linear constraint on the Yukawas and the DM masses. The corresponding region above the red line is excluded. In the region between the red and the purple lines the signal is close to the experimental sensitivity and could lead to the observation of a neutrino line in existing or near-future neutrino telescopes (see refs. [71–76]). From the figure we deduce that Scenarios 3 and 5 are therefore the most testable ones. Notice that, for different choices of $\max(m_S, m_\psi)$, the picture can be more complicated because eq. (6.4) becomes non-linear. By performing a numerical scan over the relevant Yukawa and DM mass ranges, we have checked that, as expected, there is always an excluded region in the upper right corner of figure 5, which corresponds to $y_\phi \gtrsim 10^{-10}$, $|y_S| \gtrsim 0.03$.

7 A low-energy variant: the inverse seesaw

So far, we have assumed that active neutrinos get masses through the Type-I seesaw at a very high scale. In this framework, the lepton and dark asymmetries are generated in a similar way as in high-scale leptogenesis. In the following, we discuss the possibility to embed our models into a low-scale leptogenesis scenario, trying to lower the scale of $B - L$ of the scenarios considered in this work. This is an alternative path, which may yield interesting phenomenological signals. In this case, there can be DM interactions with the SM, mediated by Z_{B-L} . Apart from direct and indirect detection signals, χ may be produced at colliders via $\bar{q}q \rightarrow Z_{B-L} \rightarrow \bar{\chi}\chi$ and then decay, $\chi \rightarrow \psi\varphi$ or $\chi \rightarrow S^\dagger\nu$, yielding missing energy.¹⁰ For $R \lesssim 1$, if the mixing of φ with the Higgs is in the correct range, displaced vertices may be produced from $\varphi \rightarrow \text{SM SM}$. Searches for long-lived particles are very active, with lots of experiments running or designed for the following years [77–79].

It has been shown that leptogenesis at a low scale is possible, for example via resonant leptogenesis, with $\mathcal{O}(10 \text{ TeV})$ right-handed neutrino masses [80–82]. Another possibility is to adopt an inverse see-saw (ISS) mechanism to give mass to neutrinos [83, 84]. Next we consider this option within a $B - L$ set-up, see also refs. [85–89] and the review [90]. We can modify our model by replacing the scalar σ with two new fields: a scalar σ' and three copies of a fermion S_L , with the quantum numbers outlined in table 3.

The Lagrangian is the same as eq. (3.1) (excluding the terms involving σ), with the addition of

$$\mathcal{L}_{\text{ISS}} = \overline{S}_L i \not{\partial} S_L - \sigma' \overline{S}_L y_{\sigma'} N_R - \frac{1}{2} \overline{S}_L \mu S_L^c + \text{H.c.}, \quad (7.1)$$

¹⁰Notice that there is no $Z_{B-L} S^\dagger S$ vertex in the Lagrangian.

where μ is a 3×3 complex symmetric matrix which can be taken to be real and diagonal without loss of generality, and $y_{\sigma'}$ is a 3×3 complex matrix. If the new scalar takes a vev, $\langle \sigma' \rangle = v_{B-L}$, $B-L$ is spontaneously broken and S_L and N_R form a pseudo-Dirac pair, with mass $M_D = y_{\sigma'} v_{B-L}$. In this case, neutrinos get a mass through the inverse see-saw mechanism. We focus on the range $M_D > m_D \gg \mu$. The mass of the active neutrinos is given by

$$m_\nu \simeq m_D M_D^{-1} \mu (M_D^{-1})^T m_D^T, \quad (7.2)$$

so that light neutrino masses may be reproduced with small values of $M_D \sim v_{B-L} \sim \mathcal{O}(\text{TeV})$ for small enough values of μ .

Both low-scale variants result into a massive gauge boson, Z_{B-L} , with a mass in the 5–10 TeV range, allowed by experimental constraints [91]. In this case, the larger $B-L$ gauge interactions allow to erase the symmetric χ population. Therefore, the $U(1)_D$ gauge interactions are not needed, and in the limit $g_D \rightarrow 0$, $U(1)_D$ acts as a global symmetry.¹¹ Notice that the presence of a global $U(1)_D$ is still crucial to stabilise the DM particles and forbid dangerous operators, such as $\bar{\psi}\phi N_R$, which would generate $\psi-N_R$ mixing. However, $U(1)_D$ could also be replaced by a Z_2 symmetry, under which all fields but χ , S and ϕ are even. In such a case, the symmetry stabilising the DM particles is a Z'_2 , coming from the combination of the broken $U(1)_X$ and the original Z_2 .

Thanks to the low scale, the annihilations $\bar{\chi}\chi \rightarrow Z_{B-L} \rightarrow \bar{q}q (\bar{l}l)$ are strong enough to erase the symmetric population of χ and make the model more testable. The cross section for the $\bar{\chi}\chi \rightarrow Z_{B-L} \rightarrow \bar{f}f$ process (assuming we are far enough from the resonance, i.e., $4m_\chi^2 - m_{Z_{B-L}}^2 \gg \Gamma_{Z_{B-L}} m_{Z_{B-L}}$), reads

$$\sigma v(\bar{\chi}\chi \rightarrow \bar{f}f) = \frac{N_C(f) g_{B-L}^4 q_{B-L}^2(f)}{2\pi} \sqrt{1 - \frac{m_f^2}{m_\chi^2} \frac{2m_\chi^2 + m_f^2}{(4m_\chi^2 - m_{Z_{B-L}}^2)^2}}, \quad (7.3)$$

where $N_C(f) = 3(1)$ and $q_{B-L}(f) = 1/3(-1)$ for quarks (leptons). The parameter space where these annihilations lead to $r_\chi < 10^{-2}$ is shown in figure 7, with $g_D = 0.5$ and $\eta_D = \eta_B$.

The different scenarios for DM are similar to the ones studied in the previous sections. However, the decay $\chi \rightarrow S^\dagger \nu_L$ is now enhanced (with respect to the previous case) by the larger value of (m_ν/M_{N_1}) . Therefore, we need smaller values of $|y_S|$ to realise a scenario with $R \ll 1$. However, too small $|y_S|$ may be problematic for the generation of the dark asymmetry. A precise lower bound is not well established and could be computed in future works. Assuming $|y_S| \gtrsim 10^{-4}$, the most natural scenario is that the decays of χ into $S^\dagger$ occur while the latter are still in equilibrium, at $T_D^{(S)} > T_*^{(S)}$, (partially) washing-out the

¹¹In the limit $g_D \rightarrow 0$, the unbroken $U(1)_{X+D}$ symmetry is now global. The would-be Goldstone boson of the broken $U(1)_{X-D}$ is now eaten by A' , which becomes massive. A' is light because its mass is proportional to the gauge coupling g_X , $m_{A'} \sim g_X v_\phi \ll v_\phi$ (recall that g_X is tiny by assumption to avoid ψ thermalisation), and it does not thermalise because it has only gauge interactions driven by g_X . A' couples to the $B-L$ current proportionally to $(v_\phi^2/v_{B-L}^2)g_X$. The interactions of a massive gauge boson are less constrained than those of a massless one. However, as now $v_\phi \lesssim v_{B-L}$, one obtains a constraint of the order of $g_X \lesssim 10^{-12}$.

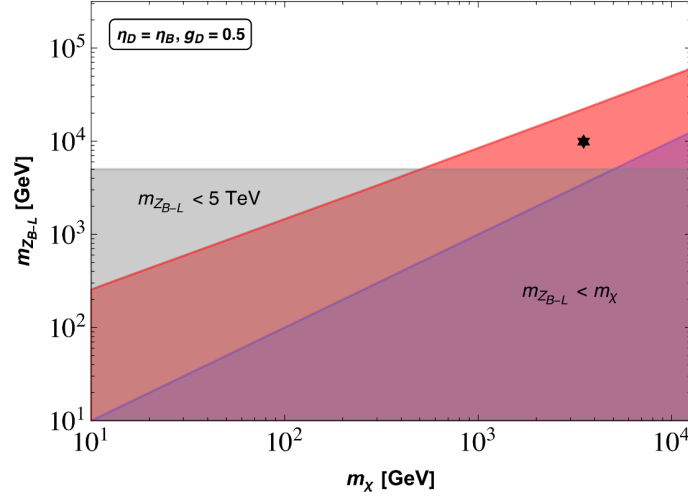


Figure 7. Parameter space for the low-scale variant with the gauge coupling fixed at 0.5 and $\eta_D = \eta_B$. In the red region the fractional asymmetry $r_\chi < 10^{-2}$. The gray region is excluded because Z_{B-L} is too light whereas in the purple region Z_{B-L} is lighter than χ . Since we expect $m_{Z_{B-L}} \gtrsim m_{N_1}$, in the purple region $N_1 \rightarrow \chi S$ decays are forbidden. The star indicates the benchmark point $\{m_\chi, m_{Z_{B-L}}\} = \{3.5 \text{ TeV}, 10 \text{ TeV}\}$.

asymmetry η_S . Therefore, if y_ϕ is small enough such that $R > 10$, χ mostly decay into $S^\dagger$ and η_S gets completely erased. Subsequently S undergoes a standard symmetric freeze-out and its abundance is determined by the annihilation cross section $\sigma v_{S^\dagger S}$.

If y_ϕ is larger, other scenarios can be realised, such as Scenarios 3, Mixed 3-4 or 5. Eventually, if both $y_\phi > 2 \times 10^{-10} \sqrt{\eta_D/\eta_B}$ and $R > 1$, a possibility not realised when $v_{B-L} \gg \text{TeV}$, a new scenario appears: ψ is symmetric and is produced mainly by freeze-in with abundance $Y_\psi^+ \simeq Y_\psi^- = Y_{\text{FI}}/(2(1+R))$ while the decays into $S^\dagger$ partially washout η_S , leaving an asymmetric abundance of S , $Y_S^+ = \eta_D/(1+R) \gg Y_S^-$. This is basically a generalisation of Scenario 5 with all the abundances rescaled by a $1+R$ factor. Notice that, in any case, in order to have multicomponent DM, S must be light enough to satisfy the neutrino constraints, corresponding to the decay $S \rightarrow \psi \nu_L$ discussed in section 6, which get stronger the smaller the v_{B-L} scale. Interestingly, in this case regions of smaller y_ϕ and $|y_S|$ could be probed.

8 Conclusions

The dark sector may be very rich. If this is the case, there is a plethora of possibilities regarding the number of stable particles, their nature, and their production mechanisms. In this work we classify and analyse the different options by adopting aogenesis model that simultaneously explains neutrino masses, the baryon asymmetry and the DM relic abundance. While neutrino masses and the baryon asymmetry are produced via the standard Type-I seesaw mechanism and leptogenesis (with some extra contributions), respectively, we find that in such a framework there is a variety of viable scenarios for explaining the nature and abundance of dark matter.

Once the decays of right handed neutrinos into the visible and dark sectors generate the asymmetries, some dark sector particles undergo asymmetric freeze-out and others are produced via freeze-in. The model has two potential DM candidates, and we focus on the parameter space where both particles are stable. However, whether they both contribute to the DM abundance similarly (two-component DM) or one has a negligible contribution (one-component DM) depends on their asymptotic asymmetries, where the late decays play an important role. Such decays may significantly populate the asymmetric or symmetric component at later times, thereby restoring annihilations, which may lead to enhanced signals in DM indirect detection. In this case, even though the DM is symmetric at the end, its abundance is still set by the asymmetry, and is thus independent of the annihilation rate, contrary to the usual WIMP scenario.

We have analysed the range of model parameters that control the contribution of each component to the DM abundance, and outline the possible scenarios, classified according to the nature and production mechanism of each particle. We consider DM masses in the GeV ballpark and dark asymmetries similar to the baryonic one; however, the set-up can easily accommodate lighter (heavier) DM for a larger (smaller) dark sector asymmetry if the branching ratio of right handed neutrinos into the dark sector is smaller (larger). We have found that one of the main distinctive signatures is a neutrino line from S (or ψ) decays. This would constitute a smoking gun of our model, within reach of existing or near-future neutrino telescopes for a significant region of the parameter space of some of the scenarios.

We conclude that having an initial asymmetry in the dark sector does not necessarily predict completely asymmetric dark matter, with its mass constrained by the dark asymmetry. In extended models, it allows the DM component to be partially asymmetric or symmetric, leading to more flexibility regarding the DM mass as well as the phenomenological implications. Finally, although in this work we focused and extended theogenesis scenario, which relates neutrino physics and dark matter, it would be interesting to consider other frameworks in which the different possibilities outlined here may also be present.

Acknowledgments

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A Details of the different scenarios

Scenario 1. In this scenario, both components are asymmetric as we take $y_\phi \lesssim 6 \times 10^{-12} \sqrt{\eta_D/\eta_B}$ and $R \ll 1$, so that the freeze-in population of ψ from early decays is negligible, while the late decays are dominant, resulting in asymmetric ψ with abundance $Y_\psi^+ = \eta_\psi = \eta_D$, whereas S freezes-out once all the symmetric population has annihilated. The population of $S^\dagger$ produced by the late decays of χ is negligible. The abundance is determined by its asymmetry, i.e. $Y_S^+ = \eta_S = \eta_D$. Hence, both the components have

Sc.	ψ	S	$\Omega_{\text{DM}}/\Omega_B$	Ω_S/Ω_ψ
1	Asymmetric LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = \eta_D$ $Y_\psi^- \ll Y_\psi^+$	Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ $Y_S^+ = \eta_D$ $Y_S^- \ll Y_S^+$	$\frac{\eta_D}{\eta_B} \frac{m_\psi + m_S}{m_p}$	$\frac{m_\psi}{m_S}$
2	Asymmetric LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = \eta_D/(1+R)$ $Y_\psi^- \ll Y_\psi^+$	Partially asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ + LD $\chi \rightarrow S^\dagger \nu_L$ $Y_S^+ = \eta_D$ $Y_S^- = \eta_D R/(1+R)$	$\frac{\eta_D}{\eta_B} \frac{m_\psi + (1+2R)m_S}{(1+R)m_p}$	$\frac{m_\psi}{m_S(1+2R)}$
1-2	Asymmetric LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = \eta_D/(1+R)$ $Y_\psi^- \ll Y_\psi^+$	Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ + LD $\chi \rightarrow S^\dagger \nu_L$ $Y_S^+ = \eta_D/(1+R)$ $Y_S^- \ll Y_S^+$	$\frac{\eta_D}{\eta_B} \frac{m_\psi + m_S}{(1+R)m_p}$	$\frac{m_\psi}{m_S}$
3	Partially asymmetric FI + LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = Y_{\text{FI}}/2 + \eta_D$ $Y_\psi^- = Y_{\text{FI}}/2$	Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ $Y_S^+ = \eta_D$ $Y_S^- \ll Y_S^+$	$\frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D m_S}{\eta_B m_p}$	$\frac{m_\psi(\eta_D + Y_{\text{FI}})}{m_S \eta_D}$
4	Partially Asymmetric FI + LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = (Y_{\text{FI}}/2 + \eta_D)/(1+R)$ $Y_\psi^- = Y_{\text{FI}}/(2(1+R))$	Partially Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ + LD $\chi \rightarrow S^\dagger \nu_L$ $Y_S^+ = \eta_D$ $Y_S^- = \eta_D R/(1+R)$	$\frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D(1+2R)m_S}{\eta_B(1+R)m_p}$	$\frac{m_\psi(\eta_D + Y_{\text{FI}})}{m_S \eta_D(1+2R)}$
3-4	Partially Asymmetric FI + LD $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = (Y_{\text{FI}}/2 + \eta_D)/(1+R)$ $Y_\psi^- = Y_{\text{FI}}/(2(1+R))$	Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ + LD $\chi \rightarrow S^\dagger \nu_L$ $Y_S^+ = \eta_D/(1+R)$ $Y_S^- \ll Y_S^+$	$\frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D m_S}{\eta_B(1+R)m_p}$	$\frac{m_\psi(\eta_D + Y_{\text{FI}})}{m_S \eta_D}$
5	Symmetric FI $\chi \rightarrow \psi\varphi$ $Y_\psi^+ = Y_{\text{FI}}/2 + \eta_D \simeq Y_{\text{FI}}/2$ $Y_\psi^- = Y_{\text{FI}}/2$	Asymmetric FO $S^\dagger S \rightarrow \varphi\varphi$ $Y_S^+ = \eta_D$ $Y_S^- \ll Y_S^+$	$\frac{\eta_D}{\eta_B} \frac{m_\psi(Y_{\text{FI}}/\eta_D) + m_S}{m_p}$	$\frac{m_\psi Y_{\text{FI}}}{m_S \eta_D}$
6	Negligible production	Symmetric FO $S^\dagger S \rightarrow \varphi\varphi$ + LD $\chi \rightarrow S^\dagger \nu_L$ $Y_S^+ = \eta_D$ $Y_S^- = \eta_D$	< 1	$\frac{\eta_D}{\eta_B} \frac{2m_S}{m_p}$

Table 4. Classification of scenarios in our model on the basis of dominant production mechanism and asymptotic nature of both dark matter components, ψ and S .

individual abundances η_D and we have

$$\frac{\Omega_\psi}{\Omega_S} = \frac{m_\psi}{m_S}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{\eta_D(m_\psi + m_S)}{m_p \eta_B}. \quad (\text{A.1})$$

In figure 8 (black lines), we show the region of the parameter space in which the correct

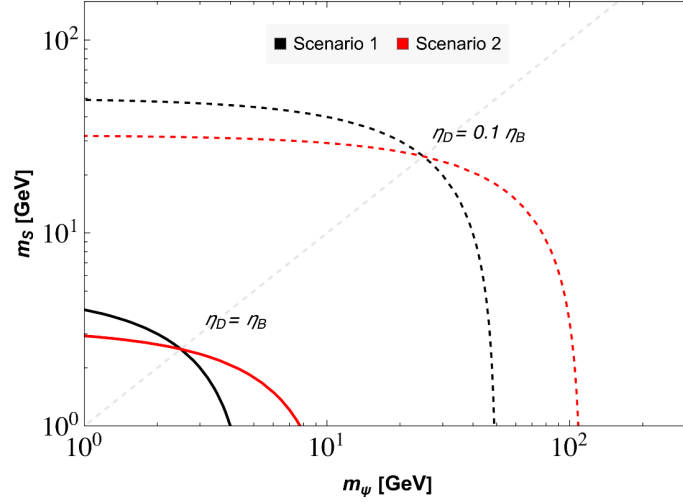


Figure 8. Values of m_ψ and m_S for which the correct relic abundance can be reproduced for Scenario 1 (black lines) and Scenario 2 (red lines). We show $\eta_D \simeq (0.1)\eta_B$ using solid (dashed) lines. The relative abundance of ψ with respect to S increases when the curves are followed clockwise. The gray dashed line corresponds to $m_\psi = m_S$. We fix $m_\chi = 3.5$ TeV, $m_{Z_D} = 500$ GeV, $g_D = 0.5$, $M_{N_1} = 10^{11}$ GeV, $m_\nu = 0.05$ eV. Furthermore, we fix $y_\phi = 10^{-12}$. For Scenario 1 (2): $|y_S| = 10^{-3}$ (5×10^{-2}), so that $R \simeq 0.0005$ (1.25). In both cases $Y_{\text{FI}} \simeq 10^{-4}\eta_B$.

DM relic abundance is reproduced in the plane m_S versus m_ψ , for two values of the dark asymmetry η_D . Ω_ψ/Ω_S increases when the curves are followed clockwise. We can see that for $\eta_D \simeq 0.1(1)\eta_B$, dashed (solid) black line, we have $m_\psi \simeq m_S \simeq 30(3)$ GeV if both species contribute similarly. Alternatively, we can push the mass of S down to GeV in such a way that its contribution to DM abundance is negligible and reproduce the relic abundance for $m_\psi \simeq 50(5)$ GeV for $\eta_D \simeq 0.1(1)\eta_B$, or vice versa.

Scenario 2. Similar to the previous scenario, here also we take $y_\phi \lesssim 6 \times 10^{-12} \sqrt{\eta_D/\eta_B}$ but $R \sim \mathcal{O}(1)$ (i.e., $|y_S| \sim 0.5 \times 10^{11} y_\phi$). We also assume that $|y_S|$ is small enough so that eq. (4.25) is satisfied, i.e., the decays of χ to $S^\dagger$ occur when the latter has already decoupled from the thermal bath, $T_D^{(S)} < T_*^{(S)}$ and $S^\dagger S$ annihilations are not active.

Due to $R \sim \mathcal{O}(1)$, χ partially decays into ψ and partially into $S^\dagger$ with probabilities $1/(1+R)$ and $R/(1+R)$, respectively. ψ is highly asymmetric with abundance $Y_\psi^+ = \eta_\psi \sim \eta_D/(1+R)$, whereas S becomes partially asymmetric because a population of $S^\dagger$ is produced by $\chi \rightarrow S^\dagger \nu$ decays. Therefore, $\eta_S = \eta_D - R\eta_D/(1+R) = \eta_D/(1+R)$, and its total abundance is determined by the sum $Y_S^+ + Y_S^- = (1+2R)\eta_D/(1+R)$. The contribution to the DM abundance becomes

$$\frac{\Omega_\psi}{\Omega_S} \approx \frac{m_\psi}{m_S} \frac{1}{(1+2R)}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{\eta_D(m_\psi + (1+2R)m_S)}{(1+R)m_p\eta_B}. \quad (\text{A.2})$$

As discussed earlier, this scenario leads to an enhanced background of an asymmetric neutrino population. In the limit $R = 0$, all χ decay into ψ and we recover eq. (A.1). For $R = 1$ the relative ψ/S abundance is $m_\psi/(3m_S)$. The factor 3 comes from the fact

that while we only have ψ and no $\bar{\psi}$, we have both S and $S^\dagger$ with $Y_\psi^+ = Y_S^- = \eta_D/2$ and $Y_S^+ = \eta_D$. The allowed parameter space is shown in figure 8 by red lines. In this case, the shape of the curves is not symmetric as in Scenario 1; in the limiting case where ψ (S) dominates the abundance one needs $m_\psi \simeq 100$ GeV ($m_S \simeq 30$ GeV) for $\eta_D \simeq 0.1 \eta_B$, and masses roughly one order of magnitude smaller for $\eta_D \simeq \eta_B$, as expected.

Mixed Scenario 1-2. We consider the same range for y_ϕ and R but we assume that $|y_S|$ is large enough to violate eq. (4.25), i.e., χ decays to $S^\dagger$ while $S^\dagger S$ annihilations are still efficient, $T_D^{(S)} > T_*^{(S)}$. For ψ we find the same results of the previous section. The $S^\dagger$ population produced by χ decays partially washes-out η_S , leaving an asymmetric population of S with abundance $Y_S^+ = \eta_D/(1+R)$, while the $S^\dagger$ population gets erased by $S^\dagger S$ annihilations. The contribution to the DM abundance is now

$$\frac{\Omega_\psi}{\Omega_S} \simeq \frac{m_\psi}{m_S}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{\eta_D(m_\psi + m_S)}{(1+R)m_p\eta_B}. \quad (\text{A.3})$$

This scenario is a mixture between Scenario 1 (the nature of the DM particles is the same, both asymmetric, and the abundances are the same ones rescaled by $(1+R)$) and Scenario 2 (the ranges of y_ϕ and R are identical and the same processes determine the final abundance).

Scenario 3. For larger values of the Yukawa, $6 \times 10^{-12} \sqrt{\eta_D/\eta_B} \lesssim y_\phi \lesssim 2 \times 10^{-10} \sqrt{\eta_D/\eta_B}$, the freeze-in contribution to ψ production from early decays grows and becomes comparable to the contribution from late decays. So ψ is partially asymmetric with $Y_\psi^+ \simeq Y_{\text{FI}}/2 + \eta_D$ and $Y_\psi^- \simeq Y_{\text{FI}}/2$. We take $R \ll 1$, so that all χ decay into ψ , partially while being in thermal equilibrium (freeze-in) and partially at a later time (late decays). The ψ abundance is $Y_\psi^+ + Y_\psi^- = Y_{\text{FI}} + \eta_D$. On the other hand, the abundance of S is determined by thermal freeze-out, once the symmetric population is annihilated away. Thus, S freezes-out with an asymmetry $Y_S^+ = \eta_D$. In this case the relative abundance between the two DM component and the total DM abundance are given by

$$\frac{\Omega_\psi}{\Omega_S} = \frac{m_\psi(\eta_D + Y_{\text{FI}})}{m_S\eta_D}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D m_S}{\eta_B m_p}. \quad (\text{A.4})$$

Scenario 4. If $R \sim \mathcal{O}(1)$ for the values of Yukawa y_ϕ considered in the previous scenario (assuming that eq. (4.25) is satisfied), the population of $\psi, \bar{\psi}$ is produced partially by freeze-in and partially by late decays, whereas the decays into $S^\dagger$ washout the asymmetry in S , making it partially asymmetric. Therefore, in this scenario, the late decays determine the asymptotic nature of both DM components. The corresponding expressions for the DM abundance are

$$\frac{\Omega_\psi}{\Omega_S} = \frac{m_\psi}{m_S} \frac{\eta_D + Y_{\text{FI}}}{\eta_D(1+2R)}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} = \frac{m_\psi(\eta_D + Y_{\text{FI}}) + \eta_D(1+2R)m_S}{\eta_B(1+R)m_p}. \quad (\text{A.5})$$

Mixed Scenario 3-4. Here we consider the same range of y_ϕ and R as in the previous scenario but a larger $|y_S|$, which violates eq. (4.25). The discussion for ψ is the same as in Scenario 4. In analogy with the Mixed Scenario 1-2, $S^\dagger S$ annihilations erase the $S^\dagger$ from

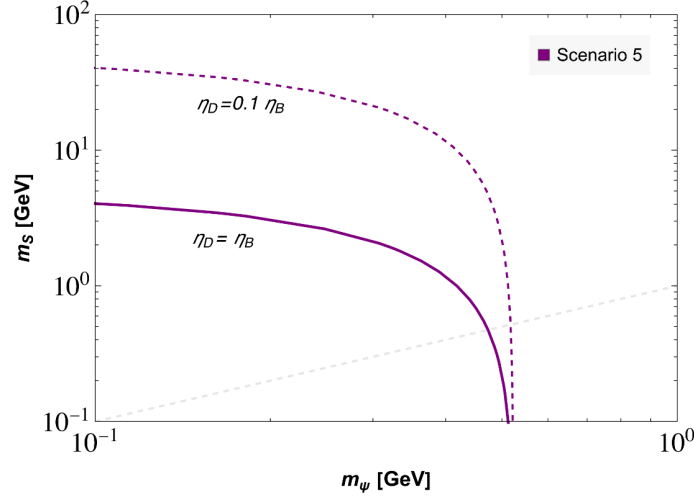


Figure 9. Values of m_ψ and m_S for which the correct relic abundance can be reproduced for Scenario 5. The gray dashed line corresponds to $m_\psi = m_S$. We fix the parameters $|y_S| = 2 \times 10^{-4}$ and $y_\phi = 2 \times 10^{-10}$. With this choice, $R \simeq 5 \times 10^{-10}$ and $Y_{\text{FI}} \simeq 10\eta_B$.

the thermal bath, leaving an asymmetric population of S which survives the annihilations, so that

$$\frac{\Omega_\psi}{\Omega_S} \simeq \frac{m_\psi}{m_S} \frac{Y_{\text{FI}} + \eta_D}{\eta_D}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{m_\psi(Y_{\text{FI}} + \eta_D) + m_S \eta_D}{(1 + R)m_p \eta_B}. \quad (\text{A.6})$$

This scenario is a mixture between Scenarios 3 and 4.

Scenario 5. For even larger Yukawas, in the range $2 \times 10^{-10} \sqrt{\eta_D/\eta_B} \lesssim y_\phi \lesssim 5 \times 10^{-7}$, the ψ sector is mainly populated during freeze-in, while late decays only produce a subdominant component, i.e., $Y_{\text{FI}} \gg \eta_D$. Therefore, the ψ population is (almost) symmetric, $Y_\psi^+ = Y_{\text{FI}}/2 + \eta_D \simeq Y_{\text{FI}}$. Concurrently, R is typically small because of the larger value of y_ϕ and therefore S freezes with an asymmetric abundance, $Y_S^+ = \eta_D$.

Therefore, DM is mostly symmetric in ψ and asymmetric in S , produced by freeze-in and freeze-out, respectively. The relative abundance of the two species and the total DM abundance are

$$\frac{\Omega_\psi}{\Omega_S} = \frac{m_\psi}{m_S} \frac{Y_{\text{FI}}}{\eta_D}, \quad \frac{\Omega_{\text{DM}}}{\Omega_B} \simeq 5 = \frac{m_\psi Y_{\text{FI}} + m_S \eta_D}{m_p \eta_B}. \quad (\text{A.7})$$

In figure 9 we show the mass ranges of ψ and S for which the DM relic abundance can be reproduced for different values of η_D . Notice that if S is subdominant, the mass of ψ is fixed, irrespective of the value of asymmetry. This is due to the fact that the DM mass is fixed by the freeze-in contribution in eq. (4.10), which is (almost) independent of m_ψ . On the other hand, it depends on the mass of χ . For the values used in the figure ($0.1\eta_B < \eta_D < \eta_B$, etc.), we see that DM could be mainly composed by light symmetric ψ of mass around 500 MeV, mainly by asymmetric GeV-ish S , or by a combination of them.

As we have underlined above, for these values of y_ϕ , R is typically small. However, if we lower the v_{B-L} scale it is possible to reach $R > 1$ even in this scenario. We briefly discuss this possibility at the end of section 7.

Scenario 6. Finally, when $R \gtrsim \mathcal{O}(10)$, for whatever value of y_ϕ compatible with it, the majority of the χ population decays into $S^\dagger$ after freeze-out of S (but before BBN if $10^{-3} < |y_S| < 0.1(m_S/\text{GeV})$) and the asymmetry is completely washed-out. However, the populations of S and $S^\dagger$ survive independently, as the rate of $S^\dagger S$ annihilations drops below the Hubble expansion rate by that time. Hence, the DM relic is almost completely made up by the symmetric population of S and $S^\dagger$, i.e., $Y_S^+ = Y_S^-$, while there is a negligible abundance of ψ produced from early or late decays and leads to $m_S \simeq 2.5 \text{ GeV}$ (η_B/η_D). The late restoration of the symmetric component would thus result in the enhancement of indirect detection signals.

Notice that for larger $|y_S|$ the decays take place while S and $S^\dagger$ are still in equilibrium so that $S^\dagger S$ annihilations washout the asymmetry and one recovers the standard scenario of symmetric freeze-out in which the S abundance is determined by the annihilation cross section $\sigma v_{S^\dagger S}$ instead of the asymmetry. Depending on m_S and $\sigma v_{S^\dagger S}$ it may be possible to reproduce the correct relic abundance. We did not consider this possibility in the main text.

B Additional constraints

B.1 Decays of φ to SM particles

We ensure that φ particles produced from χ decay fast enough into SM particles, so that the energy transfer from χ to SM radiation occurs before BBN. A rigorous bound arises from imposing $\tau_\chi + \tau_\varphi \lesssim 1 \text{ sec}$. However, it is sufficient to check that at the time of χ decays, corresponding to $T = T_D$, the decays of φ are fast compared to the Hubble rate and the bound in eq. (4.17) applies (corresponding to $\tau_\chi \lesssim 1 \text{ sec}$). The decay $\varphi \rightarrow \text{SM}$ can occur through the Higgs portal operator $\lambda_{H\phi}|\phi|^2|H|^2$ as both the scalars take a vev. Through this portal φ decays into SM fermions, mainly (if kinematically allowed) into $\bar{b}b$ or $\bar{c}c$ or lighter species if $m_\varphi < \text{GeV}$. The decay rate is $\Gamma_\varphi = \sin^2 \theta \times \Gamma_{h_\varphi \rightarrow \bar{f}f}$, where $\sin^2 \theta$ is the mixing among φ and the SM Higgs boson h while $\Gamma_{h_\varphi \rightarrow \bar{f}f} \sim (m_\varphi/32\pi)(m_f/v_{\text{EW}})^2$ is the decay width of a SM Higgs boson with mass m_φ into SM fermions.

We require that the decay is fast at $T = T_D$ (when φ are produced via χ decays), i.e., $\Gamma_\varphi/H|_{T=T_D} > 1$. For $m_\chi \sim \text{TeV}$, $m_\varphi \sim \text{few GeV} > 2m_c$ (or eventually $2m_b$) and $y_\phi \sim 10^{-12}$ the ratio Γ_φ/H at $T = T_D$ is $\gg 1$ even for small mixing angle, $\sin \theta \gtrsim 10^{-8}$. Even lighter $m_\varphi \gtrsim \text{MeV}$ is allowed as φ can decay to $\bar{u}u, \bar{d}d, \bar{e}e$, with a larger mixing angle $\sin \theta > 5 \times 10^{-4}$. Bounds from LEP constrain the mixing of a GeV-ish scalar to the SM Higgs to be $\sin \theta < 0.1$ [70]. Therefore, the condition of eq. (4.17) is valid. m_φ lighter than MeV cannot decay to any SM fermion. However it could decay into 2 photons through the effective Higgs-photon interactions. We do not study this possibility and we consider $m_\varphi > \text{MeV}$. This also implies $m_S \gtrsim \text{MeV}$.

B.2 Gauge boson spectrum and constraints on massless A'_μ

The gauge boson masses after symmetry breaking are given by

$$\begin{cases} m_{A'}^2 = 0, \\ m_{Z_D}^2 = 2g_D^2 v_\phi^2 \left(1 + \mathcal{O} \left(g_X^2, (v_\phi/v_{B-L})^2, \kappa^2 \right) \right), \\ m_{Z_{B-L}}^2 = 8g_{B-L}^2 v_{B-L}^2 \left(1 + \mathcal{O} \left(g_X^2, (v_\phi/v_{B-L})^2, \kappa^2 \right) \right), \end{cases} \quad (\text{B.1})$$

where we used that $g_X \ll 1$, $\kappa \ll 1$ and $v_\phi \ll v_{B-L}$. The mixing among them at the leading order in the small expansion parameters (omitting Lorentz indices) can be expressed as

$$\begin{cases} A'^0 \simeq A' - g_X/g_D Z_D, \\ Z_D^0 \simeq Z_D + g_X/g_D A' - \left[(g_D/4g_{B-L})(v_\phi^2/v_{B-L}^2) + \kappa \right] Z_{B-L}, \\ Z_{B-L}^0 \simeq Z_{B-L} + (g_D/4g_{B-L})(v_\phi^2/v_{B-L}^2)(1 + \kappa g_{B-L}/g_D) Z_D' - (g_X/4g_{B-L})(v_\phi^2/v_{B-L}^2) A'. \end{cases}$$

The massless gauge bosons A'_μ only interact with the fermion ψ_0 . Taking into account the fermion mixing, the Lagrangian contains the interaction terms $\bar{\psi}\psi A'$ (suppressed by g_X), $\bar{\psi}\chi A' + \text{H.c.}$ (suppressed by $g_X\epsilon_f$) and $\bar{\chi}\chi A'$ (suppressed by $g_X\epsilon_f^2$). These vertices give rise to scattering processes as $\bar{\chi}\chi \rightarrow A'A'$, $\bar{\psi}\psi \rightarrow A'A'$, ... or decays $\chi \rightarrow A'\psi$. However, given the smallness of g_X (and ϵ_f) these processes are extremely suppressed and no sizeable population of A' (which in principle could contribute to dark radiation) is produced.

We can give an upper bound on the value of the gauge coupling coming from long-range force experiments. Indeed the mixing between A' and Z_{B-L} induces an interaction $g_{\text{eff}} A'_\mu J_{B-L}^\mu$ with $g_{\text{eff}} \simeq (g_X/g_{B-L})(v_\phi/v_{B-L})^2$. For the reference values of this work, $v_\phi \sim \text{TeV}$ and $v_{B-L} \sim 10^{11} \text{ GeV}$, this corresponds to an effective coupling $g_{\text{eff}} \sim 10^{-16} g_X/g_{B-L}$. Long-range force experiments constrain the coupling to the $B-L$ current to be $g_{\text{eff}} \lesssim 10^{-24}$ [92], which implies $g_X \lesssim 10^{-8}$ (for $g_{B-L} \sim \mathcal{O}(1)$). This is comparable with the condition for ψ not to reach thermal equilibrium.

B.3 Implications of fermion mixing

Diagonalization of the fermionic sector at leading order in ϵ_f leads to the following masses for the mass-eigenstates χ and ψ ,

$$\begin{cases} m_\psi = m_\psi^0 - \epsilon_f^2(m_\chi^0 - m_\psi^0) + \mathcal{O}(\epsilon_f^3), \\ m_\chi = m_\chi^0 + \epsilon_f^2(m_\chi^0 - m_\psi^0) + \mathcal{O}(\epsilon_f^3). \end{cases} \quad (\text{B.2})$$

The fields in the mass basis are related to the original fields as

$$\begin{cases} \psi = (1 - \epsilon_f^2/2) \psi_0 - \epsilon_f \chi_0 + \mathcal{O}(\epsilon_f^3), \\ \chi = (1 - \epsilon_f^2/2) \chi_0 + \epsilon_f \psi_0 + \mathcal{O}(\epsilon_f^3). \end{cases} \quad (\text{B.3})$$

Fermion mixing induces new interactions due to the $\bar{\psi}Z_D\chi$ coupling, suppressed by ϵ_f , leading to new production processes such as $\bar{\chi}\chi \rightarrow \bar{\chi}\psi$. However, these scattering processes are typically sub-dominant to decays due to suppression by y_ϕ and ϵ_f as well as a phase space suppression (taking $\Delta = 1$), and thus can be neglected.

In addition to the above, the mixing could also lead to an additional contribution to ψ production. For $T > v_\phi$, the particles in the thermal bath are χ_0 and ψ_0 and there is no mixing. For $T < v_\phi$, the vev of ϕ induces the mixing so that the states χ_0 contain a small ψ component, proportional to ϵ_f . Therefore, $\chi_0 \rightarrow \psi$ conversions contribute to the final ψ abundance (this is analogous to the production of sterile neutrino from active neutrino mixing). This is also freeze-in process as the population of ψ is produced non-thermally from a small coupling (the fermion mixing). In analogy with sterile neutrino DM, the

interaction rate of ψ is $\Gamma_\psi \simeq \epsilon_f^2 \Gamma_\chi$, with $\Gamma_\chi = n_\chi \sigma_{\bar{\chi}\chi}$. If $v_\phi \lesssim m_\chi$ the χ particles are non-relativistic but still in equilibrium when the mixing is generated. So, $n_\chi = n_\chi^{\text{eq}}$ and these processes contribute to the symmetric component of ψ . Using the number density at equilibrium we checked that for the typical values of our parameters ($m_\chi \simeq 3.5 \text{ TeV}$, $g_D \simeq 0.5$, $v_\phi \sim \text{TeV}$) this contribution is at most comparable (but not larger) than the one from decays. A detailed study of this contribution, solving the full Boltzmann Equations (or using the density matrix formalism) is beyond the scope of this work. Therefore, we neglect this contribution, having in mind that it would change the total (symmetric) freeze-in contribution by at most an $\mathcal{O}(1)$ factor.

C Contributions to operator $\mathcal{O}_6 = \bar{L}\tilde{H}S\phi^\dagger\psi$

In section 6, we discussed the stability of the two DM components S and ψ and the possible observation of a neutrino line from the decay of one component into the other. As we saw, the decay is mediated by the dimension-6 operator

$$\mathcal{O}_6 = \bar{L}\tilde{H}S\phi^\dagger\psi, \quad (\text{C.1})$$

which is generated at low energy by first integrating the right-handed neutrino field and then the fermion χ , as discussed in the main text. A second contribution to $\mathcal{O}_6$ may arise if we assume that the theory contains the dimension-5 operator

$$\mathcal{O}_5 = \frac{\bar{N}_R(S\phi^\dagger)\psi}{\Lambda_{\text{UV}}}, \quad (\text{C.2})$$

where Λ_{UV} is a cut-off scale parametrising the UV completion of our model at scales above v_{B-L} . Integrating out N_R gives rise to

$$\frac{y_\nu}{\Lambda_{\text{UV}}M_{N_1}}\mathcal{O}_6. \quad (\text{C.3})$$

The condition on the lifetime $\tau_S^{\text{re-sc}} > 10^{23} \text{ s}$ (taking $m_S > m_\psi$), is now satisfied if

$$\Lambda_{\text{UV}} \gtrsim 10^{16} \text{ GeV} \left(\frac{m_\nu}{0.05 \text{ eV}} \right)^{1/2} \left(\frac{10^{11} \text{ GeV}}{M_{N_1}} \right)^{1/2} \left(\frac{v_\phi}{\text{TeV}} \right) \left(\frac{m_S}{50 \text{ GeV}} \right)^{1/2}. \quad (\text{C.4})$$

This scale must be at most Planckian, i.e., $\Lambda_{\text{UV}} \lesssim M_{\text{Pl}}$, which implies a bound on the mass of the heaviest of the 2 DM particles

$$m_S \lesssim 10^4 \text{ TeV} \left(\frac{0.05 \text{ eV}}{m_\nu} \right) \left(\frac{M_{N_1}}{10^{11} \text{ GeV}} \right) \left(\frac{700 \text{ GeV}}{v_\phi} \right)^2. \quad (\text{C.5})$$

Finally, $\mathcal{O}_6$ could be directly generated by UV physics as $\mathcal{O}_6/\Lambda_{\text{UV}}'$. In this case we just need

$$\Lambda_{\text{UV}}' \gtrsim 10^{15} \text{ GeV} \left(\frac{v_\phi}{\text{TeV}} \right)^{1/2} \left(\frac{m_S}{50 \text{ GeV}} \right)^{1/4}, \quad (\text{C.6})$$

which gives a weaker constraint. Notice that the same bounds apply on m_ψ if $m_\psi > m_S$. Summarising, eqs. (6.4) and (C.5) set an upper bound on the mass of the heavier between ψ and S . Therefore, this shows that it is quite natural that both ψ and S are stable on cosmological scales, and that they both contribute to the DM relic abundance and respect current limits from neutrinos.

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Multi-component dark sectors: symmetries, asymmetries and conversions

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ABSTRACT: We study the relic abundance of several stable particles from a generic dark sector, including the possible presence of dark asymmetries. After discussing the different possibilities for stabilising multi-component dark matter, we analyse the final relic abundance of the symmetric and asymmetric dark matter components, paying special attention to the role of the unavoidable conversions between dark matter states. We find an exponential dependence of the asymmetries of the heavier components on annihilations and conversions. We conclude that having similar symmetric and asymmetric components is a natural outcome in many scenarios of multi-component dark matter. This has novel phenomenological implications, which we briefly discuss.

KEYWORDS: Models for Dark Matter, Particle Nature of Dark Matter

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1 Introduction

Several observations indicate the existence of non-visible matter termed as Dark Matter (DM), that comprises roughly 25% of the total energy budget of the universe, and is therefore almost five times more abundant than the visible matter [1]. Despite this significant contribution, hardly anything is known about its composition, nature or dynamics. Among existing efforts and speculations made regarding the particle nature of DM, Weakly Interacting Massive Particles (WIMPs) are one of the simplest and most popular particle DM candidates, contributing to the energy density once they undergo thermal freeze-out. On the other hand, Asymmetric Dark Matter (ADM) models are motivated by the close relation between DM and baryonic energy densities, with $\rho_{\text{DM}} \sim 5\rho_{\text{B}}$ [2]. Given that

the baryon density is set by the baryon asymmetry of the universe, $\eta_B \sim 6 \times 10^{-10}$, the DM abundance could stem from an asymmetry in the dark sector in a similar manner (see refs. [3, 4] for a review). In the case of WIMPs, the value of the abundance is set by the thermally averaged annihilation cross section, whereas for ADM it is set by the initial asymmetry and the DM mass. Moreover, as discussed in ref. [5], there is also the possibility of an intermediate regime between the two extremes discussed above, where the DM is partially asymmetric and there is an exponential dependence on the annihilation cross section.

However, most of these scenarios assume that the entire dark sector is composed of just one DM state. Given the plethora of stable particles in the Standard Model (SM) (electrons, protons, nuclei, neutrinos, photons...), it is reasonable to think that the dark sector could be made up of more than one stable component. Efforts are being made towards the study of such multi-component DM models, where the dark sector is comprised by several stable particles. For example, in refs. [6–8] the influence of DM conversions (annihilation of one DM state into another) in the final relic abundance was emphasised. Similarly, semi-annihilations can affect the final abundance [9]. Finally, multi-component DM, with its many degrees of freedom, leads to a richer phenomenology (see for instance refs. [10–12]), and can help to relax the existing bounds from direct detection, indirect detection and collider searches. Another proposal for multi-component dark sectors is Dynamical Dark Matter [13]. In this framework, the dark states come in large numbers and are unstable, with lifetimes proportional to their individual abundance. For a recent review of their possible collective effects see ref. [14].

The numerous states and interactions with SM and among themselves in a multi-component DM setup allow us to treat it as a complex system, which exhibits a behaviour that can be characterised by different scalings: power law and exponential. Such scaling laws are common in Physics and in other branches of science. Understanding this behaviour of multi-component DM subject to different interactions (annihilations and conversions) and conditions (symmetric/asymmetric) is one of the main goals of this paper. For that purpose, we will extend the analysis performed for one component asymmetric WIMP in ref. [5] to two or more components with conversions among them.

The paper is structured as follows. In section 2, we discuss the symmetries and combinations that are required to stabilise multiple states in the dark sector. In section 3, we briefly discuss how DM asymmetries in different components can be generated from a $B - L$ asymmetry. In section 4, we develop the multi-component framework for both symmetric and asymmetric components. First, we study the final relic abundance of multi-component DM taking into account the unavoidable DM conversions. Then, in section 5, we parameterise the possible particle physics models using effective interactions. We discuss the phenomenological implications of these scenarios in section 6 and provide some conclusions in section 7.

2 Dark sector stabilisation symmetries

In this section, we briefly discuss the simplest possible symmetries to have multi-component DM. In DM models, usually a discrete and/or continuous global symmetry is imposed

to stabilise the DM candidate. As pointed in ref. [15], continuous gauge symmetries alone cannot stabilise the DM and need to be spontaneously broken. However, they can motivate the origin of discrete gauge symmetries as a remnant of spontaneously broken gauge symmetries [16], such as matter parity (and R -parity in supersymmetry) from gauged $B-L$ [17]. In this work, we consider a multi-component DM framework that can be stabilized by adopting different symmetries, with the presence of both symmetric and asymmetric components. For symmetric DM, it is more natural to have a Majorana fermion or real scalar, stabilised by discrete symmetries. On the other hand, for asymmetric DM, the candidate needs to be a Dirac fermion or a complex scalar, stabilised by continuous symmetries.

Let us start with the case of discrete symmetries. Most popular DM models employ a single Z_2 symmetry which guarantees the stability of the lightest Z_2 odd particle, allowing only one DM candidate in the model. In order to have multi-component DM, we would require a larger discrete symmetry or consider products of more than one discrete symmetry. We first investigate the number of stable DM candidates that are allowed for $Z_{\mathcal{N}>2}$, in particular for $\mathcal{N} = 3, 4$, and then for the product of two or more discrete symmetries.

Under an abelian $Z_{\mathcal{N}}$ symmetry, a field transforms as

$$\phi \rightarrow \omega^q \phi, \quad \omega = \exp(2\pi i/\mathcal{N}), \quad (2.1)$$

where $q = 0, 1, 2, \dots, \mathcal{N} - 1$ are the $Z_{\mathcal{N}}$ charges. Depending on how the fields transform under a given charge, we can divide them into classes. Class 0 corresponding to $q = 0$ (implying that the field is neutral under the given symmetry) is by default common to all $Z_{\mathcal{N}}$ symmetries. For example, consider the case of Z_2 , where $q = 0, 1$ and the field transforms as $\phi \rightarrow \phi$ (even) and $\phi \rightarrow -\phi$ (odd) respectively. Thus, there are only two classes for $Z_2 : \{0, 1\}$. The lightest Z_2 even particle will be from the SM and the odd, the DM candidate.

In the case of Z_3 , we can assign charges $q = 0, 1, 2$ to the fields. However, the charge $q = 2$ corresponding to ω^2 is same as assigning $q = -1$ since $\omega^2 = \omega^*$, implying that the particle with charge $q = 2$ is nothing but the antiparticle of the particle with $q = 1$. Therefore, in the case of Z_3 , we also have only two classes $\{0, 1\}$ and there is only one DM candidate, the lightest particle transforming under a non-trivial Z_3 charge [18].

For Z_4 , $\omega^q = 1, i, -1, -i$ for $q = 0, 1, 2, 3$ respectively. Here, note that $q = 3$ is equivalent to $q = -1$, and they would thus be placed in the same class. So, the classes for Z_4 are $\{0, 1, 2\}$. The lightest particle of class 1 ($i, -i$), let's say ϕ_1 is always stable and a DM candidate. The lightest particle of class 2 (-1), say ϕ_2 can also be stable, provided $m_{\phi_2} < 2m_{\phi_1}$. Thus, Z_4 guarantees the stability of at least one DM candidate, but can also allow for a second DM candidate in case its decays are kinematically forbidden to the particles belonging to class 1 [19]. For DM models with Z_3 and Z_4 symmetries, see refs. [9, 20–22].

In general, if $\mathcal{N}$ is a prime number, the lightest particle transforming with a non-trivial $Z_{\mathcal{N}}$ charge is always stable, and a potential DM candidate. One can also have heavier particles with $Z_{\mathcal{N}}$ charges as DM if their decay modes are kinematically forbidden. If $\mathcal{N}$ is a composite number, such that $\mathcal{N} = p_1^{s_1} p_2^{s_2} \dots p_k^{s_k}$, where p_i is a prime number ($p_i \neq p_j$ for

$i \neq j$) and s_i is natural, $Z_{\mathcal{N}}$ can be decomposed as a product of smaller groups [23, 24]

$$Z_{\mathcal{N}} \simeq Z_{p_1}^{s_1} \times Z_{p_2}^{s_2} \times \dots \times Z_{p_k}^{s_k}, \quad (2.2)$$

and there can be at most $s_1 + s_2 + \dots + s_k$ stable particles depending on the kinematic conditions. Therefore, in order to realize multi-component DM with just one imposed discrete symmetry $Z_{\mathcal{N}}$, we require $\mathcal{N} \geq 4$. For studies of multi-component scalar DM in larger $Z_{\mathcal{N}}$ groups, see ref. [25].

Now we turn our discussion to the product of two or more of the same discrete symmetries. The simplest and most commonly used symmetry is $Z_2 \times Z_2'$ [7, 8, 11, 26–28]. Such a choice of symmetry allows the existence of at least two stable particles and at most three, subject to kinematics. For the case of an n Z_2 product: $Z_2 \times Z_2' \times \dots \times Z_2^{(n)'}$, there are at least n stable DM particles and at most $2^n - 1$ stable particles can be DM candidates. It is easy to generalise this for the case of Z_3 and Z_4

$$Z_{\mathcal{N}} \times Z_{\mathcal{N}}' \times \dots \times Z_{\mathcal{N}}^{(n)'} \Rightarrow \begin{cases} n & \text{at least} \\ c^n - 1 & \text{at most,} \end{cases} \quad (2.3)$$

where c denotes the number of classes corresponding to the discrete symmetry $Z_{\mathcal{N}}$. For $\mathcal{N} = 3$, since there are only two classes, the number of DM candidates for $Z_3 \times Z_3'$ is equal to that of $Z_2 \times Z_2'$ (see ref. [29] for a model with a $Z_3 \times Z_3'$ symmetry). On the other hand, in the case of Z_4 which has three classes, $Z_4 \times Z_4'$ allows for at max $3^2 - 1 = 8$ DM candidates. The ‘ -1 ’ term guarantees stability of the lightest particle in the SM.

We can also consider products of different discrete symmetries, for example $Z_2 \times Z_3$ (see ref. [30]), and generalise eq. (2.3) as

$$Z_{\mathcal{N}_1}^{n_1} \times Z_{\mathcal{N}_2}^{n_2} \times \dots \times Z_{\mathcal{N}_k}^{n_k} \Rightarrow \begin{cases} n_1 + n_2 + \dots + n_k & \text{at least} \\ \prod_{i=1}^k c_i^{n_i} - 1 & \text{at most,} \end{cases} \quad (2.4)$$

where $\mathcal{N}_i$ corresponds to the discrete symmetry used and n_i denotes its multiplicity in the product.

In the following, we will also be interested in the case of multi-component DM with states that can be asymmetric. In this case, the most natural framework is the stabilisation of the DM components due to continuous global symmetries, which may be accidental. The possibilities we will implicitly consider are:

i) a single global continuous $U(1)$ symmetry:

$$U(1), \quad (2.5)$$

with the DM χ_i having charge q_i such that no linear terms can be written in χ_i at least at the renormalizable level. There could be linear terms at the non-renormalizable level, possibly suppressed by Planck-scale physics and therefore sufficiently suppressed [31].

ii) The direct product of several $U(1)$ s, such that each particle only transforms under one of them and is also the lightest particle charged under it,

$$U_1(1) \times U_i(1) \times \cdots \times U_N(1), \quad (2.6)$$

with DM i transforming with the global charge $\chi_i \sim (0, \dots, i = 1, \dots, 0)$. For multi-component DM models stabilised by the product of global symmetries, see refs. [6, 7].

3 Asymmetries in the dark sector

In general, DM may have both symmetric and asymmetric components. One of the simplest options is that a dark asymmetry can be generated from an asymmetry in the visible sector [2]. The idea is that an asymmetry in the SM sector is transferred to the DM sector (for instance, via a generalised $B - L$). Standard Model (DM) operators $\mathcal{O}_{\text{SM}} (\mathcal{O}_i)$ with $Q_{B-L} (-Q_{B-L})$ charges generate a mixed operator,

$$\mathcal{L}_{B-L=0} = \sum_i \frac{1}{\Lambda_i^n} \mathcal{O}_i \mathcal{O}_{\text{SM}}, \quad (3.1)$$

which is active at high energies and transfers any initial $B - L$ asymmetry in the SM to the DM sector. For definiteness, let us consider the case of Dirac DM, χ_i . We define $n_i^{+(-)}$ as the number densities of the χ_i particle (anti-particle), with $n_i^+ > n_i^-$ without loss of generality, and similarly for baryons. The baryon number density, $B \equiv n_b^+ - n_b^-$, can be expressed in terms of $B - L$ as $B \equiv b (B - L)$, where $b = 30/97 (28/79) = 0.31 (0.35)$ for $T < T_{\text{EW}} (T > T_{\text{EW}})$ [32]. The DM number density is $X_i \equiv n_i^+ - n_i^-$, which can also be expressed in terms of $B - L$ as $X_i \equiv x_i (B - L)$. It is useful to express the number densities in terms of the yields $Y_i \equiv n_i/s$, where s is the entropy density given by $s(T) = (2\pi^2/45) h_{\text{eff}}(T) T^3$ and $h_{\text{eff}}(T)$ are the effective degrees of freedom related to entropy. We also define $\eta_i = Y_i^+ - Y_i^-$, $\eta_B = Y_b^+ - Y_b^-$, and their ratio

$$\epsilon_i \equiv \frac{x_i}{b} \equiv \frac{\eta_i}{\eta_B}. \quad (3.2)$$

Depending on whether the interaction rate generated by the transfer operator, eq. (3.1), is active below or above the EW scale, one obtains [32, 33]:

- $T_{\text{int}} < T_{\text{EW}}$: $x_i = -11/111 Q_{B-L}$, $\epsilon_i = -0.32 Q_{B-L}$.
- $T_{\text{int}} > T_{\text{EW}}$: $x_i = -11/79 Q_{B-L}$, $\epsilon_i = -0.39 Q_{B-L}$.

If there are strong dark interactions (large annihilations), the symmetric component may be totally erased. In that case, in order to reproduce the DM abundance, we require

$$\frac{\Omega_{\text{DM}}}{\Omega_{\text{B}}} \simeq 5 = \frac{\sum_i \epsilon_i m_i}{m_b}. \quad (3.3)$$

If this is the case, there is a prediction for the combination of DM masses. It only depends on the energy at which the operator is active, and on how the asymmetry is transferred, i.e., on the $B - L$ charge. For example, in the case of 2DM there are different possibilities:

1. Identical DM particles with same charges, e.g. $B - L \rightarrow X_1, X_2 = X_1$: the same operator couples to both DM particles. For instance, if the operator is active above (below) the EW scale, one obtains

$$m_1 + m_2 \simeq -\frac{13(16)}{Q_{B-L}} \text{ GeV} . \quad (3.4)$$

In the case of N components, we get

$$\sum_{i=1}^N m_i = -\frac{13(16)}{Q_{B-L}} \text{ GeV} . \quad (3.5)$$

Therefore, the larger the number of asymmetric states, the lower their masses. Indeed, if all masses are similar, the presence of N -asymmetric components would translate into an N -suppressed mass-scale compared to the 1-DM scenario,

$$m^{\text{N-DM}} \simeq \frac{m^{\text{1-DM}}}{N} . \quad (3.6)$$

Notice that this can be understood as an upper limit on the mass scale in order for the DM not to be overabundant. Of course, as for 1DM, heavier DM is allowed if the DM asymmetry is unrelated or suppressed compared to the baryonic one.

2. Different DM particles with different charges, e.g., $B - L \rightarrow X_1, X_2 \neq X_1$: different operators couple to both DM particles. For example, if operator one (two) with $B - L$ charge Q_1 (Q_2) is active above (below) the EW scale, one gets

$$\frac{m_1 Q_1}{6.5 \text{ GeV}} + \frac{m_2 Q_2}{8 \text{ GeV}} \simeq 2 . \quad (3.7)$$

In the case of $N = (N_1 + N_2)$ DM components, where N_1 (N_2) components are coupled to operators active above (below) the EW scale, the above equation generalizes to

$$\sum_{i=1}^{N_1} \frac{m_i Q_i}{(13 \text{ GeV}/N)} + \sum_{j=1}^{N_2} \frac{m_j Q_j}{(8 \text{ GeV}/N)} = N . \quad (3.8)$$

Note how eq. (3.4) (eq. (3.5)) is a particular case of eq. (3.7) (eq. (3.8)) for identical charges above (below) the EW scale.

3. Different DM particles with different charges, $B - L \rightarrow X_1 \rightarrow X_2$. Only one of them, χ_1 , is coupled to the SM; the asymmetry is first transferred to one DM state and then there is a further sharing of the asymmetry with another DM state. For example, this can be generated by the operators

$$\mathcal{L}_{B-L=0} = \frac{1}{M_1^3} \chi_1^3 L H + \frac{1}{M_2} \chi_1^2 \phi_2^2 , \quad (3.9)$$

with ϕ being a scalar DM particle. This implies $x_2 = 2 x_1$ and therefore¹

$$\frac{m_1}{6.5 \text{ GeV}} + \frac{m_2}{3.25 \text{ GeV}} = 2 . \quad (3.10)$$

¹The generalisation to N -component DM is very model-dependent and should be further investigated.

In the next section we consider the general case where both asymmetric and symmetric components may be present. If the asymmetric component dominates, the previous mass relations are very good approximations.

4 Multi-component framework

In the following, we set up the framework for multi-component DM with both symmetric and asymmetric components. Let us consider an N -component DM scenario stabilised by a certain choice of continuous $U(1)$ symmetries, as discussed in section 2, with $m_1 > m_2 > \dots > m_N > m_{\text{SM}}$. Within this setup we may have one or more of the following reactions:

1. Annihilations $\chi_i \bar{\chi}_i \rightarrow \text{SM} \bar{\text{SM}}$. We use the notation $\sigma_{\chi_i \bar{\chi}_i \rightarrow \text{SM} \text{SM}} \equiv \sigma_{\text{ann}, i}$.
2. Coannihilations $\chi_i \bar{\chi}_j \rightarrow \text{SM} \bar{\text{SM}}$, where χ_j is a state close in mass to χ_i . We can neglect coannihilations if $\Delta_{ij} \equiv (m_i - m_j)/m_j \gtrsim 0.05$.
3. Conversions $\chi_i \bar{\chi}_i \rightarrow \chi_j \bar{\chi}_j$, for $m_i > m_j$, which change the individual DM number densities but not the total one.² We use the notation $\sigma_{\chi_i \bar{\chi}_i \rightarrow \chi_j \bar{\chi}_j} \equiv \sigma_{\text{conv}, ij}$.
4. Semi-annihilations $\chi_i \chi_j \rightarrow \chi_k \text{SM}$, for $N > 2$. These processes depend on the symmetry imposed. In this case, DM needs to carry some SM charge.

In the following, we consider only the case of annihilations and conversions. We are interested in conversions of the form $\chi_i \bar{\chi}_i \leftrightarrow \chi_j \bar{\chi}_j$, which are not forbidden by any symmetry that allows co/annihilations and are therefore typically always present. This is the case if both χ_i and χ_j are charged under Z_2 and Z_4 (class 2) symmetries respectively. Other conversions such as $\chi_i \chi_i \leftrightarrow \chi_j \chi_j$, for example, are phase-space suppressed and allowed only if χ_i is charged under Z_3 , and χ_j under Z_2 . Similarly, conversions of the form $\chi_i \chi_i \chi_j \rightarrow \bar{\chi}_i \bar{\chi}_j \bar{\chi}_j$ are also phase-space suppressed.

Furthermore, we always have annihilations that are significant and comparable to conversions. This avoids significant reheating in the dark sector if DM states become non-relativistic and their temperatures increase exponentially [39]. The non-negligible coupling to the SM also ensures that they are thermalised with the SM plasma at freeze-out. Therefore, considering only $2 \rightarrow 2$ processes, assuming kinetic equilibrium with the SM of all species and neglecting the quantum statistics for the species involved, the evolution equation for the number densities of the particles read

$$\begin{aligned} \frac{dn_i^\pm}{dt} + 3Hn_i^\pm = & -\langle \sigma_{\text{ann}, i} v \rangle \left(n_i^+ n_i^- - \bar{n}_i^+ \bar{n}_i^- \right) - \sum_{i < j} \langle \sigma_{\text{conv}, ij} v \rangle \left(n_i^+ n_i^- - \frac{\bar{n}_i^+ \bar{n}_i^-}{\bar{n}_j^+ \bar{n}_j^-} n_j^+ n_j^- \right) \\ & + \sum_{i > j} \langle \sigma_{\text{conv}, ij} v \rangle \left(n_j^+ n_j^- - \frac{\bar{n}_j^+ \bar{n}_j^-}{\bar{n}_i^+ \bar{n}_i^-} n_i^+ n_i^- \right), \end{aligned} \quad (4.1)$$

²Note that the forbidden channel processes such as $\chi_j \bar{\chi}_j \rightarrow \chi_i \bar{\chi}_i$, i.e. annihilation of lighter DM to heavier DM are operative at high temperatures and may be significant for small splittings, $\Delta_{ij} < 1$ [34–38]. The Boltzmann equations already take this effect into account.

where $\bar{n}_i^\pm = n_{\text{eq}} e^{\pm \xi}$ denote equilibrium number densities of the species, for which we use Maxwell-Boltzmann distributions, $n_{\text{eq}} = g(mT/2\pi)^{3/2} e^{-m/T}$, with $\mu = \xi T$ the chemical potential and g the number of internal degrees of freedom. Throughout this work, we denote by $\langle \sigma_{\text{ann}, i} v \rangle$ and $\langle \sigma_{\text{conv}, ij} v \rangle$ the thermally-averaged annihilation and conversion cross sections, respectively. The Hubble parameter is given in the radiation-dominated epoch by $H(T) = (8\pi^3/90)^{1/2} g_{\text{eff}}^{1/2}(T) T^2/M_{\text{Pl}}$, with Planck's mass M_{Pl} equal to 1.2×10^{19} GeV, and $g_{\text{eff}}(T)$ the effective relativistic degrees of freedom.³ Furthermore, the temperature evolution of the Universe can be expressed by scaling it with the mass of the heaviest component (m_1 in our case) using $x = m_1/T$, so that eq. (4.1) can be written in terms of the yields as

$$\begin{aligned} \frac{dY_i^\pm}{dx} = & -\frac{s}{Hx} \left[\langle \sigma_{\text{ann}, i} v \rangle \left(Y_i^+ Y_i^- - \bar{Y}_i^+ \bar{Y}_i^- \right) - \sum_{i < j} \langle \sigma_{\text{conv}, ij} v \rangle \left(Y_i^+ Y_i^- - \frac{\bar{Y}_i^+ \bar{Y}_i^-}{\bar{Y}_j^+ \bar{Y}_j^-} Y_j^+ Y_j^- \right) \right. \\ & \left. + \sum_{i > j} \langle \sigma_{\text{conv}, ij} v \rangle \left(Y_j^+ Y_j^- - \frac{\bar{Y}_j^+ \bar{Y}_j^-}{\bar{Y}_i^+ \bar{Y}_i^-} Y_i^+ Y_i^- \right) \right]. \end{aligned} \quad (4.2)$$

For ease of analysis and to highlight the qualitative aspects of the setup, we first consider a two-component DM system and discuss the cases where both components are symmetric or asymmetric. Later on, we extend the analysis to more complex sectors.

4.1 Symmetric components

For 2DM with symmetric components, eq. (4.2) can be written as

$$\begin{aligned} \frac{dY_1}{dx} = & -\frac{s}{Hx} \left[\langle \sigma_{\text{ann}, 1} v \rangle \left(Y_1^2 - \bar{Y}_1^2 \right) + \langle \sigma_{\text{conv}, 12} v \rangle \left(Y_1^2 - \frac{Y_2^2}{\bar{Y}_2^2} \bar{Y}_1^2 \right) \right], \\ \frac{dY_2}{dx} = & -\frac{s}{Hx} \left[\langle \sigma_{\text{ann}, 2} v \rangle \left(Y_2^2 - \bar{Y}_2^2 \right) - \langle \sigma_{\text{conv}, 12} v \rangle \left(Y_1^2 - \frac{Y_2^2}{\bar{Y}_2^2} \bar{Y}_1^2 \right) \right]. \end{aligned} \quad (4.3)$$

If conversions are neglected, the Boltzmann equations in eqs. (4.3) decouple, and in order to reproduce the observed relic abundance, we require (for $m_i \geq \mathcal{O}(\text{GeV})$ and assuming s -wave annihilations),

$$\Omega h^2 \sim \frac{1}{\langle \sigma_{\text{ann}, 1} v \rangle} + \frac{1}{\langle \sigma_{\text{ann}, 2} v \rangle} \equiv \frac{1}{\langle \sigma v \rangle_{\text{eff}}} \simeq \frac{1}{2.2 \cdot 10^{-26} \text{ cm}^3/\text{s}} = 1.37 \text{ pb}^{-1}. \quad (4.4)$$

Therefore, the abundance is dominated by the particle with the *smallest* annihilation cross section, i.e., it is like 1DM, unless $\langle \sigma_{\text{ann}, 1} v \rangle \simeq \langle \sigma_{\text{ann}, 2} v \rangle$. If both particles have similar annihilation cross sections, the yield is dominated by the lightest particle, and the contributions to the relic abundance from both states are of similar size.

On the other hand, if conversions are significant, both components may end up with similar abundances, i.e., DM may be multi-component. For instance, assuming the hierarchy $\langle \sigma_{\text{ann}, 2} v \rangle \gg \langle \sigma_{\text{conv}} v \rangle \gg \langle \sigma_{\text{ann}, 1} v \rangle$, the abundance of χ_1 (χ_2) gets reduced (increased) with respect to the case without conversions [6–8, 11, 28]. This behaviour is illustrated in figure 1,

³We take the Hubble parameter to be dominated by the SM degrees of freedom by assuming that there are no significant numbers of relativistic degrees of freedom in the dark sector.

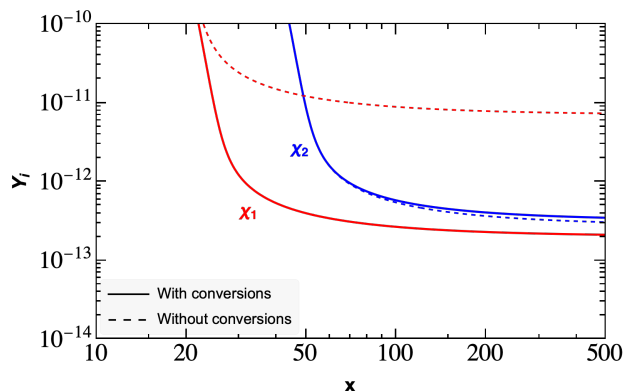


Figure 1. Evolution of yields $Y_{1,2}$ (solid curves) as a function of $x = m_1/T$ with $\langle\sigma_{\text{ann},2}v\rangle = 6 \times 10^{-9} \text{ GeV}^{-2}$, $\langle\sigma_{\text{ann},1}v\rangle = 0.1 \times 10^{-9} \text{ GeV}^{-2}$, $\langle\sigma_{\text{conv},12}v\rangle = 4 \times 10^{-9} \text{ GeV}^{-2}$ for $m_1 = 1000 \text{ GeV}$ and $m_2 = 500 \text{ GeV}$. The dashed curves correspond to the case without conversions i.e. $\langle\sigma_{\text{conv},12}v\rangle = 0$. Notice that the freeze-out of χ_2 happens at $x = m_1/T \simeq 40 \Rightarrow m_2/T \simeq 20$. A similar result was obtained in ref. [8].

where we show the evolution of abundances for χ_1 and χ_2 for the assumed hierarchy of cross sections in presence and absence of conversions. With conversions, eq. (4.4) is roughly modified as

$$\Omega h^2 \sim \frac{1}{\langle\sigma_{\text{ann},1}v\rangle + \langle\sigma_{\text{conv},12}v\rangle} + \frac{1}{\langle\sigma_{\text{ann},2}v\rangle} \simeq \frac{1}{\langle\sigma v\rangle_{\text{eff}}}. \quad (4.5)$$

It is clear from this that conversions reduce the total relic abundance, and so, in comparison to the single-DM-component case, smaller annihilations into SM states are necessary. This in turn allows to have smaller couplings between the visible and the dark sector, and therefore, to relax existing constraints.

We can extend the above discussion to N components. Let us consider the case where there are N DM components that have similar cross-sections for annihilations and conversions, e.g. $\langle\sigma_{\text{ann},i}v\rangle = \langle\sigma_{\text{conv},ij}\rangle = \langle\sigma v\rangle$ for all i, j . This simplifies the analysis so that the only parameters of interest are the individual DM masses m_i and $\langle\sigma v\rangle$. For example, let us consider the case of $N = 10$ and take $m_i - m_{i+1} = 100 \text{ GeV}$, with $m_1 = 1000 \text{ GeV}$.

In figure 2, we plot the individual asymptotic yields (left panel) and total abundance (right panel), as well as the fits to power laws of DM mass. Since the cross section is equal for all components, the yield goes as $\sim m^{-1}$ for the case without conversions and $\sim m^{-2}$ for the case with conversions. The inverse dependence on mass is well known from the analytic expression for the yield in the standard freeze-out case without conversions, as all the Boltzmann equations are decoupled. The additional mass suppression in the case of conversions can be understood as follows: the heavier the mass of the DM component, the larger the number of channels for conversions. For example, χ_1 , being the heaviest component, can undergo conversions to all the other 9 components whereas for χ_9 , the only allowed conversion is $\chi_9\chi_9 \rightarrow \chi_{10}\chi_{10}$. This implies that the individual abundances are almost independent of mass in the case without conversions, whereas in presence of conversions, they are inversely proportional to mass, as $\Omega_i \propto m_i Y_i$. For the lightest component, the yields are identical in both cases since there are no channels for its conversion

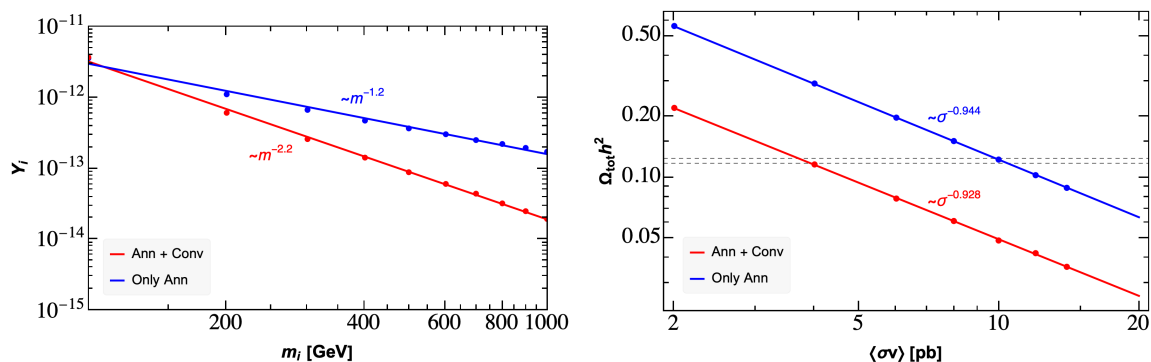


Figure 2. Power-law fits for the yields (left panel) and total abundance (right panel) in the symmetric case for 10 DM particles with equal annihilation and conversion cross section, $\langle\sigma v\rangle = 2 \text{ pb}$. *Left panel:* individual yields versus DM mass; we choose $\langle\sigma v\rangle = 2 \text{ pb}$. *Right panel:* total abundance vs cross section; the grey dashed lines represent the 3σ range for the observed DM relic abundance, $\Omega_{\text{DM}} h^2 = 0.12 \pm 0.0012$ [1].

and the excess due to its production from all the other components gets annihilated to SM. In figure 2 (right panel), we can see that the total abundance is inversely proportional to the cross section in both the cases. Furthermore, the presence of conversions allows us to reproduce the correct abundance for smaller values of the annihilation cross-section (roughly a factor of 2), and thus in general of the DM couplings to SM. This is a generic feature, which as argued before translates in a relaxation of the experimental bounds.

Finally, let us mention that for cold thermal relics, there exists a bound on how heavy the DM particle can be from unitarity in order not to overclose the universe: $m_\chi \lesssim 110 \text{ TeV}$ [40, 41]. In the case of multi-component DM, this translates into a bound on the masses of the DM components for s-wave annihilations,

$$\left(\sum_i^N m_i^2 \right)^{1/2} \lesssim 110 \text{ TeV} ,$$

when only annihilations are considered. When conversions are present, the bound reads⁴

$$\left(\sum_i^N \frac{m_i^2}{(N-i+1)} \right)^{1/2} \lesssim 110 \text{ TeV} .$$

Therefore, as for the case of asymmetric DM, the existence of a complex dark sector points to lower DM masses. Again, if all masses are similar, the presence of N -symmetric components translates into an $\sqrt{N}$ -suppressed mass-scale compared to the 1-DM scenario,

$$m_{\text{sym}}^{N\text{-DM}} \lesssim \frac{m_{\text{sym}}^{1\text{-DM}}}{\sqrt{N}} \lesssim \frac{110 \text{ TeV}}{\sqrt{N}} , \quad (4.6)$$

where in the last step we assume s -wave annihilations and conversions. Notice how, now, the suppression with N is milder as compared to the asymmetric case, cf. eq. (3.6). The above bounds may be relaxed to 190 TeV (220 TeV) for p -wave (for a combination of s - and p -wave) annihilations and conversions [41].

⁴Here, we assume s -wave approximation for both conversions and annihilations.

4.2 Asymmetric components

4.2.1 Analytical expressions

Now we study how the picture changes if there are initial DM asymmetries. We follow closely the work of ref. [5] and extend it to the case of 2DM with initial asymmetries $\eta_i = Y_i^+ - Y_i^-$, $i = 1, 2$. Of course, the asymmetries remain constant and equal to the initial value throughout the evolution. Following their notation, we can define the asymmetric ratios $r_i \equiv Y_i^-/Y_i^+$, which are in the range $0 \leq r_i \leq 1$, where the upper (lower) limit is for totally (a)symmetric DM. In this way, the co-moving number densities $Y_i^\pm$ can be completely expressed in terms of η_i and r_i as

$$Y_i^+ = \eta_i/(1 - r_i), \text{ and } Y_i^- = r_i Y_i^+. \quad (4.7)$$

Given that the chemical potentials of particles and their antiparticles are equal and opposite, we define $\bar{r}_i \equiv e^{-2\xi(x)}$ with $\xi = \sinh^{-1}(\eta_i/(2\bar{Y}_i))$. Hence, eq. (4.2) can be written for the asymmetric ratios r_i in the suggestive forms

$$\begin{aligned} \frac{dr_1}{dx} &= -\frac{s\eta_1}{Hx} \left[\langle \sigma_{\text{ann},1} v \rangle \left(r_1 - \bar{r}_1 \zeta^2(r_1) \right) + \langle \sigma_{\text{conv},12} v \rangle \left(r_1 - \frac{r_2 \bar{r}_1}{\bar{r}_2} \frac{\zeta^2(r_1)}{\zeta^2(r_2)} \right) \right], \\ \frac{dr_2}{dx} &= -\frac{s\eta_2}{Hx} \left[\langle \sigma_{\text{ann},2} v \rangle \left(r_2 - \bar{r}_2 \zeta^2(r_2) \right) - \langle \sigma_{\text{conv},12} v \rangle \frac{\eta_1^2 (1 - r_2)^2}{\eta_2^2 (1 - r_1)^2} \left(r_1 - \frac{r_2 \bar{r}_1}{\bar{r}_2} \frac{\zeta^2(r_1)}{\zeta^2(r_2)} \right) \right], \end{aligned} \quad (4.8)$$

where $\zeta(r_i) \equiv (1 - r_i)/(1 - \bar{r}_i)$. Notice that if the initial asymmetries η_i are zero, the ratios r_i would remain constant and equal to one, like in the standard freeze-out case.

It is interesting to derive some analytical solutions to better understand the behavior of the system. Considering the evolution of r_i ($i = 1, 2$), we have two different regimes:

- For $x < x_{f_i}$, there is thermal equilibrium and $r_i(x) \simeq \bar{r}_i(x)$.
- For $x > x_{f_i}$, $r_i(x) \gg \bar{r}_i(x)$ and also $\bar{r}_1(x)/\bar{r}_2(x) \ll 1$,

where x_{f_i} denotes the freeze-out temperature for the species. We can expand the annihilation and conversion cross sections in the relative velocity and express the thermally averaged cross sections as⁵ $\langle \sigma_{\text{ann},i} v \rangle \equiv \sigma_i (T/m_i)^{k_i}$ and $\langle \sigma_{\text{conv}} v \rangle \equiv \sigma_c (T/m_1)^{k_c}$, where one can consider different combinations of s -wave ($k_{i/c} = 0$) and p -wave ($k_{i/c} = 1$) cross sections.⁶ We can write

$$\begin{aligned} \frac{s}{Hx} \langle \sigma_{\text{ann},i} v \rangle &= \lambda_{a,i} g_*^{1/2} x^{-k_i-2} \left(\frac{m_1}{m_i} \right)^{k_i}, \\ \frac{s}{Hx} \langle \sigma_{\text{conv}} v \rangle &= \lambda_{a,1} f_c g_*^{1/2} x^{-k_c-2}, \end{aligned} \quad (4.9)$$

⁵Note that we expand the conversion cross section in terms of the velocity of the heavier component.

⁶In the numerical analysis we consider an s -wave annihilation cross section. Similar results can be obtained for higher partial waves.

where $\lambda_{a,i} = (\pi/45)^{1/2} M_{\text{Pl}} m_1 \sigma_i$, $f_c \equiv \sigma_c/\sigma_1$, and

$$g_*^{1/2}(x) = \frac{h_{\text{eff}}}{g_{\text{eff}}^{1/2}} \left(1 - \frac{1}{4} \frac{x}{g_{\text{eff}}} \frac{dg_{\text{eff}}}{dx} \right). \quad (4.10)$$

In the following, for simplicity, we consider the case: $k_1 = k_2 = k_c \equiv k$, $\eta_1 = \eta_2 \equiv \eta$, $\lambda_{a,2} \simeq \lambda_{a,1} \equiv \lambda_a$ and approximate $\zeta(r_i) \sim 1$. Therefore, with the above approximations, we obtain

$$\begin{aligned} \frac{dr_1}{dx} &\simeq -\lambda_a \eta g_*^{1/2} x^{-k-2} (1 + f_c) r_1, \\ \frac{dr_2}{dx} &\simeq -\lambda_a \eta g_*^{1/2} x^{-k-2} (r_2 - f_c r_1). \end{aligned} \quad (4.11)$$

Given that $m_2 < m_1$, for similar cross sections $r_2 \gg r_1$, and hence we can neglect $f_c r_1$ in the second line of eq. (4.11) and solve the decoupled system

$$\begin{aligned} r_1(x) &\simeq \bar{r}_{1,f} e^{-\lambda_a (1+f_c) \eta \Phi_1(x, m_1)}, \\ r_2(x) &\simeq \bar{r}_{2,f} e^{-\lambda_a \eta \Phi_2(x, m_2)}, \end{aligned} \quad (4.12)$$

where $\bar{r}_{i,f} \equiv \bar{r}_i(x_{f_i})$ and

$$\Phi_i(x, m_i) \equiv \int_{x_{f_i}}^x dx' x'^{-k-2} g_*^{1/2}. \quad (4.13)$$

Therefore, the DM asymmetry ratios decrease exponentially with the annihilation cross-section (i.e., DM becomes more asymmetric as annihilations are increased), in a similar way as for 1DM scenarios [5]. Furthermore, in presence of significant conversions, the asymmetry ratio for the heavier component also decreases exponentially with the conversion cross section. This is analogous to the symmetric case (see eq. (4.5)) where the yield for the heavier components was inversely proportional to the sum of annihilation and conversion cross sections whereas that for the lighter one it was only inversely proportional to its annihilation cross section. In this case, however, the crucial difference is that there is an exponential dependence.

4.2.2 Numerical analysis

In figure 3, we plot the evolution of the asymmetric ratios and components after solving eq. (4.8) and using eq. (4.7). Note that $Y_1^+ \sim Y_2^+ \sim \eta_B$, due to our choice of $\eta_1 = \eta_2 = \eta_B$ and large annihilation cross-sections, so that the symmetric component is significantly depleted. In general, one obtains $Y_i^+ \sim \eta_i$. From the left plot in figure 3, it can be seen that the heavier component, χ_1 , is more asymmetric than the lighter one, χ_2 , without conversions. Including conversions significantly reduces the symmetric component of the heaviest state. This can be seen in the right plot where $Y_i^+ - Y_i^- \sim \eta_B$ ($i = 1, 2$) but $Y_1^- \ll Y_2^-$. Hence, we observe that while the heavier component is mostly asymmetric, the lighter one is mostly symmetric. Notice that the analytic solutions derived in section 4.2.1 are contrasted with the numerical solutions in figure 3 (dotted lines) and are found to be in agreement.

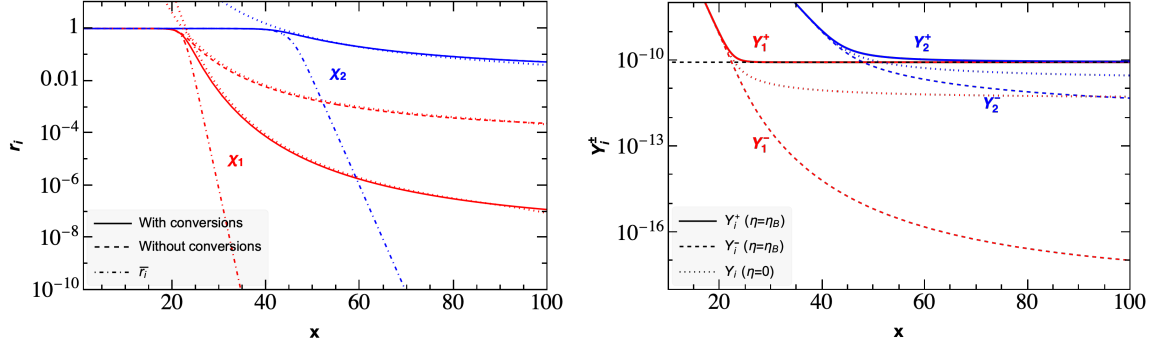


Figure 3. Evolution of asymmetric ratios and yields for χ_1 (red) and χ_2 (blue) in the asymmetric case. We choose $m_1 = 20$ GeV, $m_2 = 10$ GeV, $\langle\sigma_{\text{ann},1} v\rangle = \langle\sigma_{\text{ann},2} v\rangle = \langle\sigma_{\text{conv},12} v\rangle = 2$ pb and $\eta_{1,2} = \eta_B \sim 0.88 \times 10^{-10}$. *Left panel:* numerical solution for asymmetric ratios r_i in presence (solid) and absence (dashed) of conversions. The dot-dashed curves represent the equilibrium distributions $\bar{r}_i(x)$. The dotted red and blue curves correspond to the analytic solution and coincides well with the numerical solution. *Right panel:* the solid (dashed) curves correspond to Y_i^+ (Y_i^-), and the dotted curves correspond to the symmetric case i.e. $\eta_{1,2} = 0$. The black dashed line corresponds to the value of η_B .

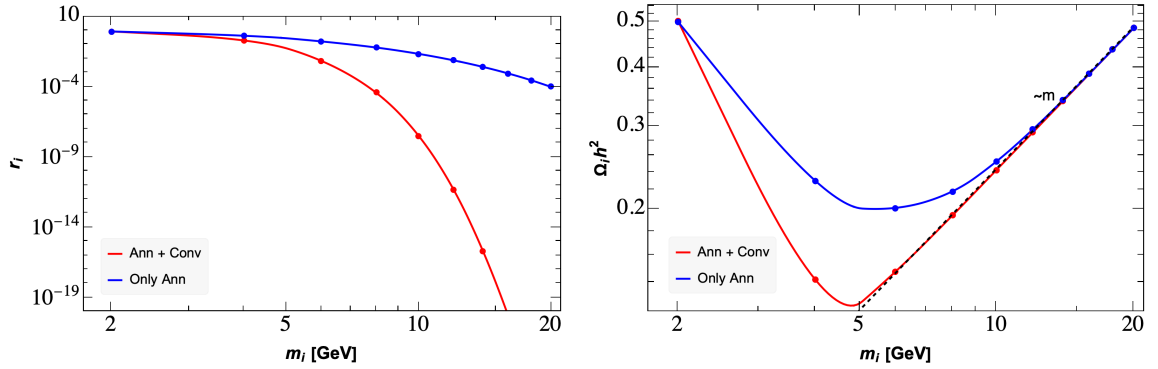


Figure 4. *Left panel:* evolution of asymmetric ratios for 10 DM particles with $m_i - m_{i+1} = 2$ GeV, $m_1 = 20$ GeV, $\langle\sigma v\rangle = 2$ pb and $\eta_i = \eta_B$, in presence (red) and absence (blue) of conversions. *Right panel:* the individual abundance of DM components versus mass in the asymmetric case. The black dashed curve depicts the direct proportionality to mass as particles become more asymmetric i.e. $r_i \rightarrow 0$.

Now we extend the analysis to $N = 10$ components to show the effects of conversions and annihilations, in a similar way to the analysis of symmetric components done above, keeping the value of annihilation and conversion cross sections equal for all components and fixing the initial asymmetries equal to the baryon asymmetry η_B . In figure 4, we show the individual asymmetric ratios (left panel) and abundances (right panel).

Here, unlike the symmetric case, one cannot fit the asymmetry ratios as power laws in the DM mass. For the case without conversions, we find $r_i \sim a e^{-b m_i}$ is a good fit whereas when conversions are included, we can fit the individual asymmetric ratios as the product of a power law and an exponential: $r_i \sim a' m_i^{b'} e^{-c m_i}$, where a', b', c depend on the value of cross sections. Notice that the DM mass enters subtly via $x_{f_i} = m_i/T_i$ in the expression

for Φ . We find that for similar values of the cross sections, the asymmetry ratios for the lighter species are larger than that of the heavier ones. This implies that, for a given initial asymmetry η , the heavier species are more asymmetric than the lighter ones, and they become even more asymmetric in the presence of conversions, thus they can be considered to be fully asymmetric.

From the right plot of figure 4, it can be seen that the individual abundances coincide for the heavier particles with and without conversions and scale with the DM mass m_i . Since the heavier species can be considered to be fully asymmetric, their abundances are independent of conversion and annihilation cross sections. This is valid if the cross-sections are larger than a certain value. For example, in this case, annihilation cross-sections larger than $\mathcal{O}(0.1)$ pb lead to $r < 0.1$.

As we can see, the presence of conversions reduces the total abundance also in the partially-asymmetric case. For initial asymmetries close to the baryonic one, i.e., $\eta_i \simeq \eta_B$, we find that only in the case of purely asymmetric DM states with masses above ~ 10 GeV, is the effect of conversions irrelevant. Furthermore, the direct proportionality of abundance with mass is a good approximation for even lower masses in the presence of conversions, which tend to make the species more asymmetric, as illustrated by the black dashed curve. We also see that the abundance is at or below current Planck values for DM masses around 5 GeV, as expected in asymmetric DM models (see section 3) due to the initial asymmetry being equal to the baryonic one.

4.3 Total relic abundance

We have seen that it is natural that if, DM has different states, the individual symmetric and asymmetric components can contribute significantly to the relic abundance for reasonable values of the cross sections. We follow closely the notation in ref. [5].

For the case with initial asymmetries, the total DM abundance will consist of both symmetric and asymmetric components and the DM is partially asymmetric, with an energy density given by

$$\rho_{\text{DM}} \equiv \sum_i \rho_i = s \sum_i m_i \eta_i \left(1 + 2 \frac{r_{\infty,i}}{1 - r_{\infty,i}} \right), \quad (4.14)$$

where $r_{\infty,i}$ depends on $\Phi_i(\infty, m_i)$. The first term corresponds to the asymmetric components (present for $r_{\infty,i} \rightarrow 0$), and the second one to the symmetric ones (which scale as the inverse power of the sum of annihilation and conversion cross sections for $r_{\infty,i} \rightarrow 1$). Reproducing the relic abundance yields a constraint between the DM masses, the asymmetries and the annihilation and conversion cross sections through this relation. Notice that in some cases the initial asymmetries may be similar, $\eta_i \simeq \eta$, i.e., if they come from the same operator, or if conversions/annihilation are large, and we can factorize η from the first term.

In the completely asymmetric case, i.e. $r_{\infty,i} = 0$, the total DM abundance just depends on the masses and initial asymmetries,

$$\rho_{\text{ADM}} = s \sum_i \eta_i m_i. \quad (4.15)$$

In the symmetric limit $r_{\infty,i} \rightarrow 1$, the expansion of the second term gives

$$\rho_{\text{SDM}} \simeq s \sqrt{\frac{180}{\pi M_{\text{Pl}}^2}} \sum_i \frac{1}{\sigma_i \Phi_i(\infty, m_i)}, \quad (4.16)$$

which is independent of the initial asymmetries and very mildly dependent on the DM mass via $\Phi_i(\infty, m_i)$. In presence of conversions, σ_i has to be summed over the conversion cross sections as well, such that total cross-section becomes

$$\sigma_{i,t} = \langle \sigma_{\text{ann},i} v \rangle + \sum_{j>i} \langle \sigma_{\text{conv},ij} v \rangle. \quad (4.17)$$

Using eqs. (3.3) and (4.16), one can also re-express the fractional asymmetry for heavier components in presence of conversions as [5]

$$r_{i,\infty} = e^{-\lambda_{i,t} \eta_i \Phi_i} = \exp \left[-2 \left(\frac{\sigma_{i,t}}{\sigma_{\text{WIMP}}} \right) \left(\frac{\Phi_{i,t}}{\Phi_{\text{WIMP}}} \right) \frac{1 - r_{i,\infty}}{1 + r_{i,\infty}} \right], \quad (4.18)$$

where $\lambda_{i,t} = (\pi/45)^{1/2} M_{\text{Pl}} m_1 (\sigma_i + \sigma_c)$, and σ_{WIMP} and Φ_{WIMP} correspond to a symmetric DM particle of the same mass, m_i . It should be noted that the freeze-out temperature (x_{f_i}) entering the expression for Φ in eq. (4.13) is modified when conversions are included, hence we use $\Phi_{i,t}$.

Naively, the relic abundance in the partially-asymmetric case (eq. (4.14)) might be expected to be the sum of eqs. (4.15) and (4.16); however, that's not the case as the symmetric component also depends on η [5]. Only for scenarios where the DM components are completely symmetric or asymmetric, the total abundance can be given by their sum

$$\frac{\rho_{\text{DM}}}{s} \simeq \sum_{i=\text{ADM}} \eta_i m_i + \left(\frac{180}{\pi M_{\text{Pl}}^2} \right)^{1/2} \sum_{j=\text{SDM}} \frac{1}{\sigma_j \Phi(\infty, m_j)}, \quad (4.19)$$

for $i \neq j$.

In figure 5, we plot the dependence of the relic abundance on masses⁷ and cross sections to show the differences between the symmetric, asymmetric and partially-asymmetric cases. It can be seen that the curve corresponding to partially asymmetric is not equivalent to the sum of the other two, and starts tracking the asymmetric curve for larger masses. The mass at which it starts tracking the asymmetric case depends on the value of σ_i and η_i .

5 Effective parameterisation

In the analysis for symmetric and asymmetric components performed above, we have considered all cross sections to be independent of DM mass. However, these are model dependent and in particular, also scale with the masses of the particles involved. For instance, considering only s -wave annihilations we can parameterise the DM annihilation cross section at freeze-out as

$$\langle \sigma_{\text{ann},i} v \rangle \simeq \lambda_i^2 \frac{m_i^2}{(m_i^2 + \Lambda^2)^2}, \quad (5.1)$$

⁷The plot is valid for 1DM case or for multi-component DM without conversions.

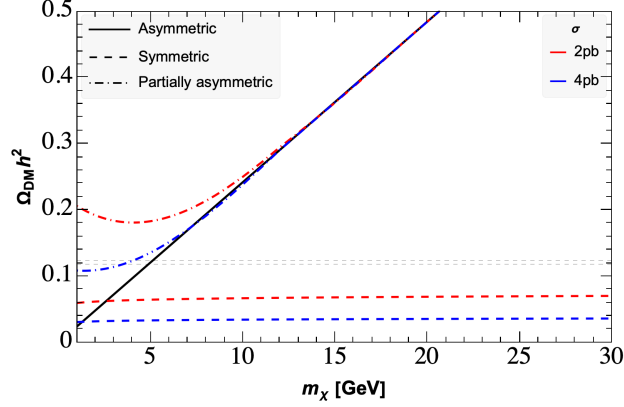


Figure 5. Relic abundance as the function of DM mass for asymmetric (solid), symmetric (dashed) and partially-asymmetric (dot-dashed) cases. The red and blue curves correspond to two different choices of $\sigma_i = 2$ pb, 4 pb respectively. The asymmetric case is linearly dependent on mass and η , which is chosen to be equal to η_B . The symmetric case is almost independent of mass and depends only on the values of σ_i . The partially-asymmetric case becomes independent of σ_i for large mass values, e.g., it starts exhibiting a completely asymmetric behaviour.

where m_i is the DM mass, λ_i the coupling characterising the DM-SM interaction and Λ is the mass scale of the heavy mediator involved. We consider two limiting cases for our analysis,

$$\langle \sigma_{\text{ann}, i} v \rangle \sim \begin{cases} a_i/m_i^2 & \text{for } \Lambda \ll m_i \\ b_i m_i^2/\Lambda^4 & \text{for } \Lambda \gg m_i \end{cases}, \quad (5.2)$$

corresponding to light and heavy mediator respectively, with a_i and b_i being model-dependent numerical parameters. In a similar manner, we parameterise the conversion cross section as

$$\langle \sigma_{\text{conv}, ij} v \rangle \simeq \lambda_{ij}^2 \frac{m_i^2}{(m_i^2 + \Lambda^2)^2} \sqrt{1 - \frac{m_j^2}{m_i^2}}, \quad (5.3)$$

where λ_{ij} is the interaction strength between the DM components with masses m_i and m_j . As before, we consider the limiting cases

$$\langle \sigma_{\text{conv}, ij} v \rangle \sim \begin{cases} c_{ij} \frac{(m_i^2 - m_j^2)^{1/2}}{m_i^3} & \text{for } \Lambda \ll m_i \\ d_{ij} \frac{m_i (m_i^2 - m_j^2)^{1/2}}{\Lambda^4} & \text{for } \Lambda \gg m_i. \end{cases} \quad (5.4)$$

Hence, in order to describe the dynamics of N -component DM, we require N DM masses, N DM-SM couplings and $N(N-1)/2$ DM-DM couplings i.e. $N(N+3)/2$ parameters in total. In the following, we analyse the results in a model independent way, treating a, b, c, d as free parameters and fixing them to be equal for all i, j . We then provide an example model in appendix A.

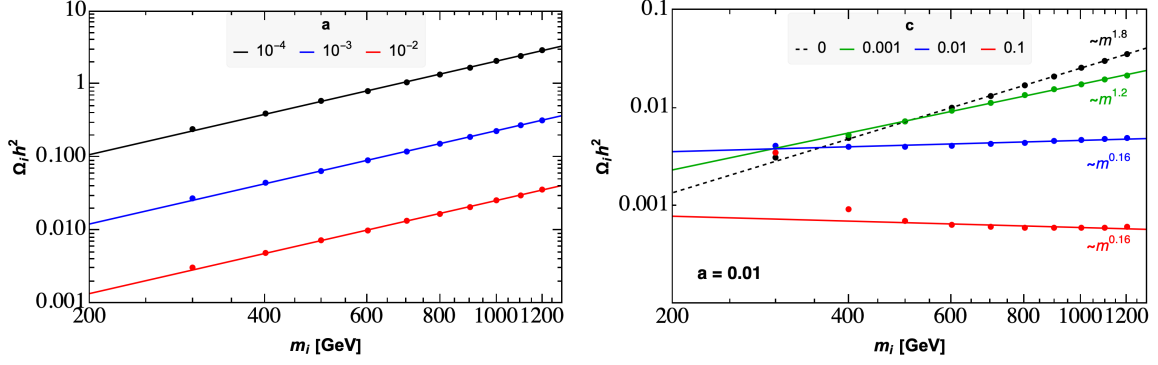


Figure 6. Individual abundances with DM mass in the light mediator case, without (left panel) and with (right panel) conversions, for $N = 10$.

5.1 Symmetric components

5.1.1 Light mediators

First, let us consider the case of light mediators, where the cross-sections are independent of the mediator mass, see first line of eqs. (5.2) and (5.4), and depend on the new parameters a_i and c_{ij} , taking $N = 10$ and varying the DM masses from 300 GeV to 1200 GeV. In the absence of conversions ($c_{ij} = 0$), all Boltzmann equations are decoupled and we roughly obtain the individual yields (abundances) scaling as $\sim m_i$ ($\sim m_i^2$). This is expected since Y_i scales as $\sim m^{-1}$ when we treat all cross sections to be similar and independent of mass.

The asymptotic abundances are plotted as a function of mass in the left plot of figure 6, for different choices of a . In this case, the abundances follow a power-law of the form

$$\Omega_i h^2 \sim \frac{\kappa}{a} m_i^{1.8}, \quad (5.5)$$

where κ is to be determined from the fits.

Moreover, we find that the scenario becomes interesting in presence of conversions with an additional dependence on the parameter c (we take all $c_{ij} \equiv c$). The total conversion cross sections of a given DM particle into the lighter states $\langle \sigma_{\text{conv}, i} v \rangle = \sum_{j>i} \langle \sigma_{\text{conv}, ij} v \rangle$ are of similar size. This is due to the fact that while heavier DM particles have more DM states to convert into, their conversion cross sections are suppressed by their mass, whereas lighter states have less DM states to convert into, but their cross sections are much larger. These effects partially compensate each other. As a result of this, the power law dependence on mass for abundances varies depending on the strength of conversions with respect to annihilations, unlike the case of mass-independent cross sections, where we saw $\Omega_i h^2 \propto m_i^{-1}$ in presence of conversions. We illustrate this in the right plot of figure 6. As can be seen, for $c \gtrsim a$, the abundances do not change much with the mass of the DM component. The abundances in this case can be expressed in the form

$$\Omega_i h^2 \sim \frac{\kappa'}{\left(\frac{a}{m_i^2} + \delta c \right)}, \quad (5.6)$$

where κ' and δ are numbers determined from the fit. If the conversions are weak, i.e. $c \ll a$, the first term in the denominator of eq. (5.6) dominates and $\Omega_i h^2 \propto m_i^2$. On

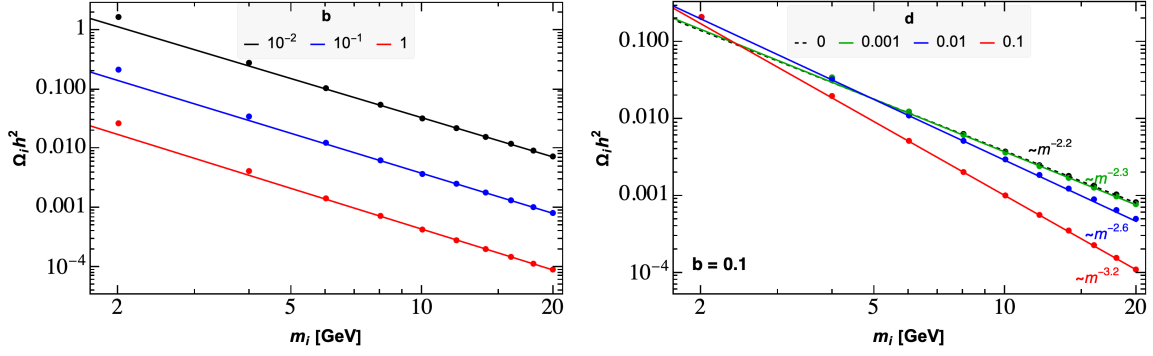


Figure 7. Individual abundances with DM mass in *the heavy mediator case*, without (*left panel*) and with (*right panel*) conversions, for $N = 10$, with $\Lambda \sim 100$ GeV.

the other hand, for significant conversions i.e. $c \sim a$ or $c \gg a$, the second term in the denominator dominates and we find $\Omega_i h^2$ does not depend on m_i . Therefore with conversions, $\Omega_i h^2 \sim \kappa' m^\beta$ with $0 \lesssim \beta \lesssim 2$.

Note that for extremely large conversions, the fits are determined by ignoring the lightest and next-to-lightest DM particles as they do not follow the same mass-suppression as heavier ones. However, for moderate to weak conversions, we find that it is possible to perform a fit for all the DM components.

5.1.2 Heavy mediators

In the case of heavy mediators and/or light DM, see second line of eqs. (5.2) and (5.4), where annihilations are proportional to b_i and conversions to d_{ij} (we take all equal), the cross sections are proportional to the squares of DM masses and the individual yields and abundances are thus expected to scale as $\sim m_i^{-3}$ and $\sim m_i^{-2}$, respectively.

The case of no conversions is shown in the left plot of figure 7 for different choices of the parameter b . We have taken the heaviest DM component mass to be 20 GeV and the lightest 2 GeV. We also fix the mediator scale to $\Lambda = 100$ GeV. The asymptotic relic abundances follow a power law of the following form

$$\Omega_i \sim \frac{\kappa}{b} m_i^{-2.2}, \quad (5.7)$$

where κ is determined from the numerical fits.

In presence of conversions, the dependence on mass will vary depending on the strength of conversions with respect to annihilations, similar to the light mediators case. However, due to the direct proportionality to mass, the total conversion cross sections of each species are not similar unlike in the previous case, and scale as $\sim m_i^3$. Therefore, we can express the abundances as

$$\Omega_i h^2 \sim \frac{\kappa' \Lambda^4}{(b m_i^2 + \delta d m_i^3)}, \quad (5.8)$$

where κ', δ are determined from the numerical fits. When conversions dominate, i.e. $d \gtrsim b$, we get $\Omega_i h^2 \propto m_i^{-3}$, whereas for weak conversions, we obtain roughly $\Omega_i h^2 \propto m_i^{-2}$. Hence, for the heavy mediators case, we have $\Omega_i h^2 \sim \kappa' m^\gamma$ with $-3 \lesssim \gamma \lesssim -2$, see the right plot of figure 7.

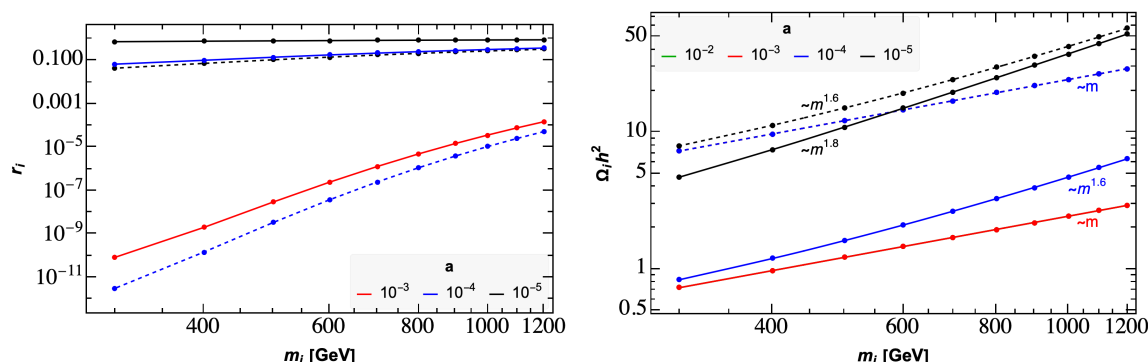


Figure 8. Asymmetry ratio (*left panel*) and abundance (*right panel*) versus DM mass for different choices of a in the *light mediator case*, without conversions, for $N = 10$. *Left:* notice that all the dashed curves ($\epsilon = 1$) are suppressed compared to the solid ($\epsilon = 0.1$) ones. The red dashed curve is highly suppressed and does not appear in the plot. *Right:* The green solid curve corresponding to $a = 10^{-2}$ coincides exactly with the red solid curve, whereas the dashed green and red curves coincide exactly with the blue dashed curve.

5.2 Asymmetric components

As discussed above, in the case of completely asymmetric dark matter, the relic abundance is independent of the annihilation and conversion cross sections, whereas there is a slight dependence on these cross sections in the case of partially-asymmetric dark matter.

From the analytic expressions derived above, we know that the asymmetry ratio is exponentially suppressed by the annihilation cross section and the initial asymmetry η_i . In the following, we discuss the effect this has on the DM relic abundance for the light and heavy mediator cases.

5.2.1 Light mediators

Let us first consider the case without conversions. When the cross sections are inversely proportional to the square of the DM masses, the lighter species will have a larger cross section and their asymmetry ratio gets exponentially suppressed compared to the heavier ones. The trends vary depending on the value of ϵ in eq. (3.2) and a in eq. (5.2), which enter in eq. (4.12). In figure 8, we show the fractional asymmetries versus mass for different choices of η (e.g. ϵ , see eq. (3.2)).

In the left plot of figure 8, it can be seen that the solid curves ($\epsilon = 0.1$) are larger than the dashed ($\epsilon = 1$) ones for the same value of a , due to the suppression by η . At the same time, curves corresponding to larger values of a have $r \ll 1$ and can be treated as completely asymmetric, while the curves corresponding to tiny values of a are close to one, and are mostly symmetric. This is reflected in the relic abundance of the species, where we see that for larger values of a , the curves coincide and are a linear function of m (i.e., the solid red and the dashed blue curves), whereas for smaller values of a , the abundances scale as $\sim m_i^2$, similar to what we observed for the symmetric components.

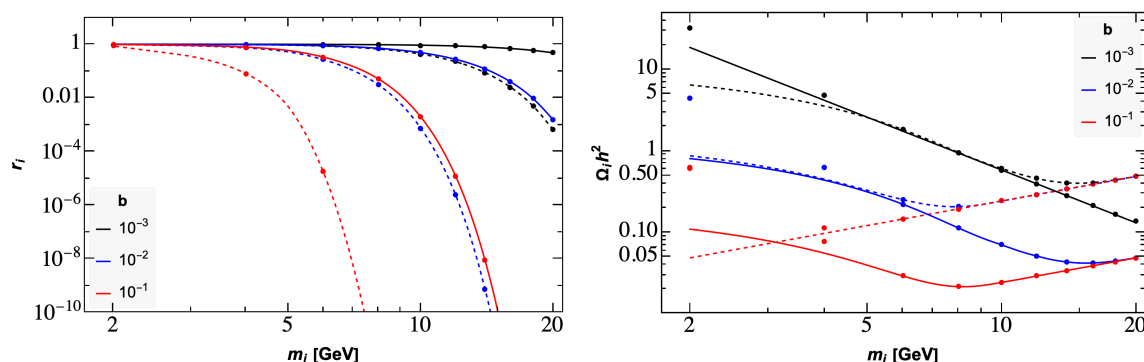


Figure 9. Asymmetry ratio (*left panel*) and abundance (*right panel*) versus DM mass for different choices of b in the heavy mediator case, without conversions, for $N = 10$, with $\epsilon = 1$ (dashed) and $\epsilon = 0.1$ (solid).

5.2.2 Heavy mediators

In the case of heavy mediators, as heavier components have larger cross sections, they are expected to be highly asymmetric in nature compared to the lighter ones. In figure 9, we show the asymmetry ratios and abundances for different choices of b (annihilations) and ϵ . We fix the mediator mass scale at 100 GeV and vary the DM masses in the range 2 – 20 GeV. Comparing the asymmetry ratios with the light mediator case, cf. figure 8, we see a reversed behaviour and now heavier particles in the set-up tend to be more asymmetric than the lighter ones. The suppression due to ϵ is evident in the left plot. Depending on the strength of annihilations, the abundance of heavier particles that are completely asymmetric is proportional to their mass whereas the abundance of the lighter ones roughly scales as $\sim m^{-2}$, similar to what we had seen in the symmetric case. Notice also how, the smaller the value of ϵ , the larger the value of the DM mass where the abundance changes behaviour from mainly being symmetric-dominated to being asymmetric-dominated.

5.2.3 Effect of conversions

Here we discuss the effect of including conversions for both cases above. Since conversions deplete the symmetric components, the heavier particles which can convert to the lighter ones become more asymmetric, compared to the case of no conversions. However, the heavier states may be more symmetric (asymmetric) than the lighter ones in the case of light (heavy) mediators (c.f the left panels in figures 8 and 9). Therefore, in presence of strong conversions, the abundance of the heavier species scales proportional to their mass and the initial asymmetry. For weak conversions, the behavior is similar to the case of only annihilations. Depending on the strength of conversions versus annihilations, one can categorise the multi-component asymmetric DM set-up as shown in table 1. The only case where the DM can be treated as partially asymmetric is when both annihilations and conversions are weak; all other combinations lead to a completely asymmetric behavior.

Annihilation/Conversion	<i>Weak</i>	<i>Strong</i>
<i>Weak</i>	Partially asymmetric	Asymmetric
<i>Strong</i>	Asymmetric	Asymmetric

Table 1. Behavior of multi-component DM (except for the lightest component), depending on interaction strength. ‘Asymmetric’ denotes that the DM is completely asymmetric.

	Cross-section	Symmetric		Partially-Asymmetric
Model	σ	Ann.	Ann.+Conv.	Ann./ Ann.+Conv.
Mass-independent	$\neq f(m_i)$	$\sim m^0$	$\sim m^{-1}$	$\sim \eta m^{[0,1]*}$
Heavy mediators	$\sim m_i^2/\Lambda^4$	$\sim m^{-2}$	$\sim m^{-[3,2]}$	$\sim \eta m^{[-2,1]*}$
Light mediators	$\sim 1/m_i^2$	$\sim m^2$	$\sim m^{[0,2]}$	$\sim \eta m^{[1,2]}$

Table 2. Scaling of the relic abundance with the DM mass for the heavier components in a multi-component DM with symmetric or asymmetric components, with and without conversions, for different types of dependence of the cross sections on the DM mass. The left/right exponents indicate the extreme limits that are obtained for strong/weak conversions (and annihilations) in the symmetric (partially-asymmetric) case. In the cases with an *, left/right exponents indicate weak/strong interactions.

5.3 Mass-dependence: power-law scaling

For the parameterisations of annihilation and conversion cross-sections used we observe that the multi-component models exhibit a power-law behavior. The scaling of the relic abundance with DM mass is summarised in table 2. The power laws are a good fit for the heavier components in the set-up and do not apply to the lightest component.⁸

We show the case of treating all cross sections as mass-independent ($\sigma \neq f(m_i)$), as well as the cases of light and heavy mediators. In the mass-independent symmetric case without conversions, all of the particles contribute with a similar abundance to the total one. Similarly, when the cross sections are inversely proportional to the DM mass, e.g. for the light mediator case, strong conversions may lead to several stable components with similar abundances in the symmetric case, e.g. the relic abundance becomes independent of the DM mass. On the other hand, in models with heavy mediators, we expect the lighter components of the dark sector to dominate the relic abundance in presence of significant conversions.

For the partially-asymmetric case, if the annihilations and conversions are strong, the heavier components are highly asymmetric and the abundance depends linearly on mass, whereas if both the processes are weak, we recover the dependence found in the symmetric case with weak conversions. Therefore, depending on the values of interaction cross-sections, it may be possible to have several partially-asymmetric particles with similar abundances in the heavy mediator case or in the mass-independent case with weak annihilations and conversions.

⁸In order for the expressions to be dimensionally correct, the scalings go as $(m_i/m_N)^k$, where k is the exponent shown in the table and m_N the lightest mass. We do not show this explicitly in the table.

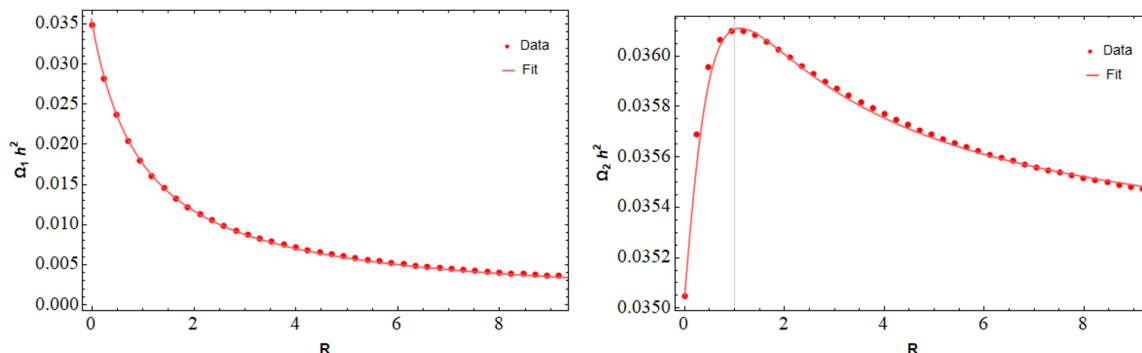


Figure 10. Abundance of the heaviest (*left panel*) and the lightest component (*right panel*) versus $R = \langle \sigma_{\text{conv}, 12} v \rangle / \langle \sigma_{\text{ann}, 1} v \rangle$ for the case of mass-independent cross-sections with $N = 2$, $\langle \sigma_{\text{ann}, 1} v \rangle = \langle \sigma_{\text{ann}, 2} v \rangle = 5.85 \times 10^{-9} \text{ GeV}^{-2}$, $m_1 = 40 \text{ GeV}$ and $m_2 = 20 \text{ GeV}$. The data is fitted with eq. (5.9), with $\alpha = 9.38 \times 10^{-5}$, $\beta = 1.65$ and $\gamma = 0.77$. The peak close to $R \simeq 1$.

5.4 Abundance of the lightest state

The power laws provided above are a good fit to describe the behavior of heavier components that undergo conversions. However, this is not true for the lightest component that is being produced by all the heavier components. In this section, we comment on an interesting effect that the lightest component undergoes in presence of conversions, which to our best of knowledge has not been discussed elsewhere.

It turns out to be useful to define $R = \langle \sigma_{\text{conv}, 12} v \rangle / \langle \sigma_{\text{ann}, 1} v \rangle$ as the ratio of the conversion and the annihilation thermally-averaged cross-sections. In figure 10, we plot the abundance of the heaviest and the lightest component versus R , for $N = 2$, in the mass-independent case. In general, one would expect the abundance of the lightest (heaviest) component to be directly (inversely) proportional to the strength of conversions (i.e. from annihilations of the heavy states into it). While the abundance of heaviest does decrease with conversions, for the lightest component, we find that the abundance does not scale directly with the conversion rate. We observe the presence of a peak in the abundance when $R \simeq 1$, after which the abundance drops but is always (slightly) larger than the case without conversions ($R = 0$). We observe that the peak appears at $R \simeq 1$, a phenomenon we term as ‘annihilation-conversion resonance’. The presence of this peak is inherent to conversions and can be attributed to the interplay of freeze-out dynamics for all the components involved. A word of caution, however, is in order: for $R \gg 1$, the computation of the relic abundance becomes more complicated, with the possible presence of significant reheating in the dark sector [39].

A suitable fit to the profile of the peak of the lighter component in figure 10 (right panel) is given by

$$\Omega_l h^2(R) = \Omega_l h^2(0) + \frac{\alpha R}{R^\beta + \gamma}, \quad (5.9)$$

where α and γ are constants to be determined from the fit, and $\Omega_l h^2(R)$ [$\Omega_l h^2(0)$] is the abundance of the lightest component in the case of conversions [no conversions]. For the heavier components (figure 10, left panel), the fit goes as $\alpha'/(1+R)$, where α' is a constant.

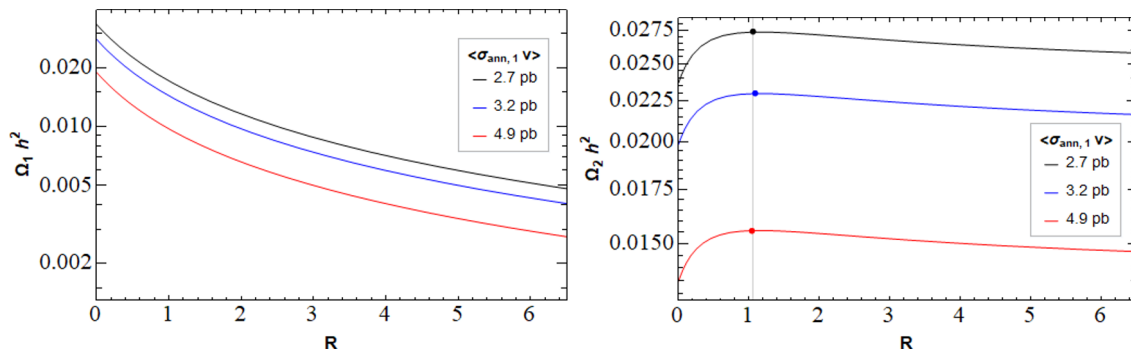


Figure 11. Abundance of the heaviest (*left panel*) and the lightest component (*right panel*) for different values of annihilation cross-sections in the *light mediator* case for $N = 2$, $m_1 = 1200$ GeV and $m_2 = 1000$ GeV.

Such peaks in the lightest component abundance can be seen in all the scenarios discussed above, whether the DM is symmetric/partially asymmetric or heavier/lighter than the mediator involved, for $N \geq 2$, as long as the mass splitting between the DM components fulfills $\Delta_{ij} \geq 0.05$. Note that the mass splitting Δ_{ij} of both DM species plays a central role in the above discussion as a consequence of eq. (5.4), since for small Δ_{ij} no peak in the abundance is observed. In figure 11, we show the case of mass-dependent cross-sections, using the parameterization of eqs. (5.2) and (5.4). We consider the light mediator case with different values of a (annihilations) and c (conversions).

6 Phenomenological implications

In general, the phenomenological implications are model dependent, and mostly driven by the interaction strength with the SM (annihilations), and to a lesser extent by conversions. Assuming similar values for the different components, the following generic features are expected.

In multi-component DM scenarios, if there are no conversions, typically the component with the smallest annihilation cross section dominates the abundance. Given that indirect detection (ID) signals are proportional to $n^2 \langle \sigma_{\text{ann},i} v \rangle \propto 1/(m^2 \langle \sigma_{\text{ann},i} v \rangle)$, for equal mass-independent annihilation cross sections, we expect the lightest component to dominate. In presence of conversions, the abundance of the heavier components gets reduced and thus the ID signals should be suppressed compared to the case with no conversions. For example, in 2DM case, the ID rates for the heavier component χ_1 and the lighter component χ_2 are

$$R_{\text{sym},1}^{\text{ID}} = \frac{1}{m_1^2} \frac{\langle \sigma_{\text{ann},1} v \rangle}{(\langle \sigma_{\text{ann},1} v \rangle + \langle \sigma_{\text{conv},12} v \rangle)^2},$$

$$R_{\text{sym},2}^{\text{ID}} = \frac{1}{m_2^2} \frac{1}{\langle \sigma_{\text{ann},2} v \rangle}. \quad (6.1)$$

The scaling of indirect detection rates with DM mass for the different scenarios are shown in table 3. We observe that in the light mediator case with only annihilations, or in presence of weak conversions, the ID rates of the different DM components are of similar order.

	Cross-section	Indirect Detection	
Model	σ	Only Ann.	Ann.+Conv.
Mass-independent	$\neq f(m_i)$	$\sim m^{-2}$	$\sim m^{-4}$
Heavy mediators	$\sim m_i^2/\Lambda^4$	$\sim m^{-4}$	$\sim m^{-[6,4]}$
Light mediators	$\sim 1/m_i^2$	$\sim m^0$	$\sim m^{[-4,0]}$

Table 3. Indirect detection rates proportional to $n^2\langle\sigma v\rangle$ for heavier components in a multi-component DM with symmetric components, for different scenarios. The exponents (left/right) indicate the extreme limits that are obtained depending on the strength of conversions (strong/weak).

In the partially-asymmetric case with 1DM, there is an exponential reduction of the asymmetry ratio (and therefore on the indirect signals) with the annihilation cross section compared to the vanilla symmetric case, as pointed in ref. [5]. In the case of partially-asymmetric multi-component DM, we find an exponential dependence of the asymmetry ratios on the conversion cross section for the heavier species, making them more asymmetric, whereas the lighter ones may remain mostly symmetric. Therefore, with strong conversions, the heavier species become even more asymmetric, e.g. $r_i \ll 1$, and their indirect detection signals are highly suppressed compared to a WIMP candidate of the same mass [5],

$$\frac{R_{\text{asym}}^{\text{ID}}}{R_{\text{sym}}^{\text{ID}}} = \frac{\sigma_{\text{ann},h}}{\sigma_{\text{sym}}} \times r_{h,\infty} \times \left(\frac{2}{1+r_{h,\infty}} \right)^2 \ll 1. \quad (6.2)$$

The ID rates for lighter components are also suppressed but not as much as the heavier ones

$$\frac{R_{\text{asym}}^{\text{ID}}}{R_{\text{sym}}^{\text{ID}}} = \frac{\sigma_{\text{ann},l}}{\sigma_{\text{sym}}} \times r_{l,\infty} \times \left(\frac{2}{1+r_{l,\infty}} \right)^2 \lesssim 1. \quad (6.3)$$

On the other hand, if the conversions are weaker and the annihilation rates are similar to the thermal one, the symmetric part of the components may be similar, as already discussed above.

One may wonder whether strong conversions may lead to a locally-enhanced population of the lighter states compared to the heavier ones in high DM-density environments, such as the galactic centre or dwarf galaxies, enhancing the previous rates. Schematically, assuming that the local energy density ρ_{local} in the considered body scales as the global one Ω_{global} , as is the case for cold DM [42, 43], we may have

$$\frac{\rho_{\text{local}}^{\text{light}}}{\rho_{\text{local}}} > \frac{\Omega_{\text{global}}^{\text{light}}}{\Omega_{\text{global}}}, \quad \frac{\rho_{\text{local}}^{\text{heavy}}}{\rho_{\text{local}}} < \frac{\Omega_{\text{global}}^{\text{heavy}}}{\Omega_{\text{global}}}. \quad (6.4)$$

If the lighter states become more populated, their annihilation channels, say into SM states α , will dominate. This would be interesting phenomenologically, given the stronger ID limits on light DM. For example, in the case of 2DM with masses $m_{1,2}$ ($m_1 > m_2$) and number densities $n_{1,2}$, conversions will lead to an increase in the number density of the lighter component. One can estimate

$$\Delta n_2 \simeq 2 n_{1,i}^2 \langle \sigma_{\text{conv},12} v \rangle t, \quad (6.5)$$

where t is the relevant time scale from the formation of the local body, and $n_{1,i}$ is the initial local number density of the heavier component at the time where the body is formed. The change in the lighter species may be estimated as

$$\Delta n_2 \simeq 10^{-9} \text{ cm}^{-3} \left[\left(\frac{\mathcal{P}}{0.67} \right) \left(\frac{\rho_{\text{local}}}{3 \text{ GeV/cm}^3} \right) \left(\frac{10 \text{ GeV}}{m_1} \right) \right]^2 \left(\frac{\langle \sigma_{\text{conv},12} v \rangle}{10^{-8} \text{ GeV}^{-2}} \right) \left(\frac{t}{8 \text{ Gy}} \right), \quad (6.6)$$

where $\mathcal{P} \equiv \rho_{\text{local}}^1 / \rho_{\text{local}} = \Omega_{\text{global}}^1 / \Omega_{\text{global}}$ is the initial fraction of the energy density in the heavy state. Therefore, the current local number densities of the states become

$$n_{1,f} \simeq n_{1,i} - \Delta n_2, \quad n_{2,f} \simeq n_{2,i} + \Delta n_2. \quad (6.7)$$

Though in this case the change is negligible, it could be significant for local bodies with much higher dark matter densities, for lighter DM, and for larger conversions. A detailed analysis of the evolution of those systems would be needed in that case.

Multi-component DM could also give signatures in direct detection, specially kinks in the energy spectra of averaged rates [44, 45], and partial cancellations and a non-sinusoidal behavior in time-dependent rates [46] (being able to reproduce DAMA Phase 2 results). The implications on direct detection rate for reproducing the relic abundance in thermal-freeze-out and asymmetric scenarios have been discussed in ref. [47]. In general, the presence of large conversions allows to reproduce the relic abundance with smaller DM-SM couplings, hence relaxing the strong bounds coming from DD null-results.

Astrophysical signals from capture in the sun and other celestial bodies (e.g. neutron stars) can constrain the set-up. For example, the asymmetric component may yield constraints from creating black-holes, etc. The symmetric part may yield the typical signals, for instance annihilation of DM components into neutrinos, and/or anomalous heating of the celestial body. The conversions $\chi_1 \chi_1 \rightarrow \chi_2 \chi_2$ with $m_1 \simeq m_2$ (inelastic scattering) can yield a larger population of the lowest mass state in the Sun (assuming thermalisation is reached) [48]. The opposite reaction $\chi_2 \chi_2 \rightarrow \chi_1 \chi_1$ is kinematically forbidden in typical bodies for non-negligible mass splittings.

Similarly, the conversions $\chi_1 \chi_1 \rightarrow \chi_2 \chi_2$ with $m_1 \simeq m_2$ (inelastic scattering) can explain small-scale structure problems in dwarfs if large enough (also bounds from clusters) [49] and again yield a larger population of the lowest mass state. The opposite reaction $\chi_2 \chi_2 \rightarrow \chi_1 \chi_1$ is typically kinematically forbidden also in clusters for non-negligible mass splittings.

At colliders, even if strong conversions make the DM to be mainly composed of the lighter state, the heavier states, if kinematically accessible, may also be produced. For LHC studies of multi-component scenarios see for example refs. [50, 51].

7 Conclusions

Guided by the visible sector, one may argue that the most natural dark sector is multi-component and partially-asymmetric. In this work, we have studied the relic abundance of several stable particles in the dark sector with symmetric and asymmetric components, with special emphasis on the role that conversions have in the final relic abundance. We

have found that the total abundance gets reduced by conversions in all cases, generically helping to *hide* the dark sector from experiment.

In many scenarios, we have found that several DM components may have similar abundances, contributing on equal footing to the relic abundance. These are the cases of: *i*) symmetric or partially-asymmetric DM with mass-independent cross sections and small conversions, or light-mediators with strong conversions; *ii*) heavy-mediator with weak conversions. We expect the dependence on the DM mass shown in table 2 to be valid for a large class of models, and to be rather universal as long as the masses and cross sections are at the weak scale.

In the symmetric case, for mass-independent cross sections, the abundance of heavier species is suppressed due to significant conversions and they are mostly asymmetric, while the lightest ones dominate the energy density. Furthermore, for large conversions and low masses, the local density of lighter states may be enhanced with respect to the global one in very high-density environments. These effects specially affect ID signals, with lighter states being more populated and yielding the largest SM final state products, while heavier ones remain mostly asymmetric and less populated.

We conclude that, in multi-component scenarios, it is natural to have symmetric and asymmetric components in presence of conversions. These can help to explain the absence of any positive signal up to now, and furthermore provide distinctive experimental signatures in current and future direct and indirect detection experiments, as well as at colliders, where cosmologically subdominant components may be produced.

Acknowledgments

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A Model example

Here we provide an example model to realize multi-component DM and demonstrate the behavior of the cross section with DM mass, in line with the discussion in section 5. One of the simplest models where this can be seen is the scalar singlet DM model, where the DM couples to the standard model via Higgs portal [52, 53]. Most of the parameter space of the real scalar singlet model is ruled out, leading to DM masses either in the funnel region of Higgs resonance or $m_{\text{DM}} > 1$ TeV [54, 55]. In the following, we use it as an illustration and extend it to N components. The parameter space for the two component model has been studied in detail in ref. [11].

Consider the standard model extended by N scalar singlets ϕ_i , all of them stabilised by $Z_2 \times Z'_2 \dots \times Z'^N_2$ symmetry and coupling to the SM via a Higgs portal coupling λ_i . The

part of the Lagrangian concerning DM interactions is given by

$$\mathcal{L}_{\text{DM-SM}} = \sum_{i=1}^N \left(\frac{1}{2} (\partial_\mu \phi_i)^2 - \frac{1}{2} m_{0_i}^2 \phi_i^2 - \frac{1}{4!} \lambda_{s_i} \phi_i^4 - \frac{1}{2} \lambda_i \phi_i^2 H^\dagger H - \frac{1}{4} \sum_{j>i} \lambda_{ij} \phi_i^2 \phi_j^2 \right), \quad (\text{A.1})$$

where λ_{s_i} is the quartic self coupling, m_{0_i} is the DM bare mass, λ_{ij} is the DM-DM coupling and H is the SM Higgs doublet, which acquires a vacuum expectation value after electroweak symmetry breaking and is parameterised by $H = (0, (v + h)/\sqrt{2})^T$, where $v = 246$ GeV, while $\langle \phi_i \rangle = 0$. Hence, the mass of the DM candidates are $m_i^2 = m_{0_i}^2 + \frac{\lambda_i v^2}{2}$ and we choose $m_i > m_j$ for $j > i$. The Higgs portal couplings (λ_i) opens annihilation channels for DM states, $\phi_i \phi_i \rightarrow ff, WW, ZZ, hh$. Furthermore, they also lead to conversions $\phi_i \phi_i \rightarrow \phi_j \phi_j$ along with the direct interactions λ_{ij} , which dominate for larger DM masses.

In the large DM mass regime (light mediator case), $m_i \gg m_h$, the annihilation cross section (s -channel) for the model can be roughly approximated as

$$\langle \sigma_{\text{ann}, i} v \rangle \approx \frac{\lambda_i^2}{16\pi m_i^2} \quad \text{for } m_i \gg m_h. \quad (\text{A.2})$$

According to our adopted convention, this implies $a_i \equiv \lambda_i^2/16\pi$.

For conversions, the interaction rate is given by

$$\langle \sigma_{\text{conv}, ij} v \rangle = \frac{\sqrt{s - 4m_j^2}}{8\pi s \sqrt{s}} \left[\frac{v^4 \lambda_i^2 \lambda_j^2}{(s - m_h^2)^2 + m_h^2 \Gamma_h^2} + \frac{2(s - m_h^2) v^2 \lambda_i \lambda_j \lambda_{ij}}{(s - m_h^2)^2 + m_h^2 \Gamma_h^2} + \lambda_{ij}^2 \right]. \quad (\text{A.3})$$

Taking the threshold value $s = 4m_i^2$, in the heavy mass regime, the direct interaction term λ_{ij} dominates over $\lambda_{i,j}$, and we can approximate

$$\langle \sigma_{\text{conv}, ij} v \rangle \approx \frac{\lambda_{ij}^2 (m_i^2 - m_j^2)^{1/2}}{32\pi m_i^3}. \quad (\text{A.4})$$

Therefore, $c_{ij} \equiv \lambda_{ij}^2/32\pi$.

In the low DM mass regime (heavy mediator case), the annihilation and conversion cross sections can be approximated by

$$\begin{aligned} \langle \sigma_{\text{ann}, i} v \rangle &\approx \frac{\lambda_i^2}{8\pi} \frac{m_i^2}{m_h^4} \quad \text{for } m_i \ll m_h, \\ \langle \sigma_{\text{conv}, ij} v \rangle &\approx \frac{\lambda_j^2}{8\pi m_h^4} m_i (m_i^2 - m_j^2)^{1/2}. \end{aligned} \quad (\text{A.5})$$

The approximation for the conversion cross section above holds in the case where the DM mass is given by $m_i^2 \sim \lambda_i v^2$. Thus, $b_i \equiv \lambda_i^2/8\pi$ and $d_{ij} \equiv \lambda_j^2/8\pi$ in this case.

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Lowering the scale of fermion triplet leptogenesis with two Higgs doublets

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In this paper, we consider the possibility of generating the observed baryon asymmetry of the Universe via leptogenesis in the context of a triplet fermion-mediated type-III seesaw model of neutrino mass. With a hierarchical spectrum of the additional fermions, the lower bound on the lightest triplet mass is $\sim 10^{10}$ GeV for successful leptogenesis, a couple of orders higher than that of the type-I case. We investigate the possibility of lowering this bound in the framework of two-Higgs-doublet models. We find that the bounds can be lowered down to 10^7 GeV for a hierarchical spectrum. If we include the flavor effects, then a further lowering by one order of magnitude is possible. We also discuss if such lowering can be compatible with the naturalness bounds on the triplet mass.

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I. INTRODUCTION

The Standard Model (SM) is immensely successful in explaining majority of the properties of the fundamental particles and their interactions. The SM was further bolstered by discovery of the Higgs boson by the ATLAS and CMS experiments of the Large Hadron Collider (LHC) [1,2]. However, there are some experimental results that cannot be accounted for in the framework of the SM. Two of the most relevant of these for our purposes are as follows: (i) The SM fails to explain the origin of the baryon asymmetry of the Universe, $Y_{\Delta B} = (8.75 \pm 0.23) \times 10^{-11}$, determined from cosmic microwave background radiation and relative abundances of light elements [3]. (ii) On the other hand, several neutrino oscillation experiments have shown that neutrinos have tiny but nonzero masses, unlike what is predicted by the SM. One of the most appealing scenarios to explain the tiny neutrino masses needed to explain the results from oscillation experiments is the seesaw mechanism, which is generated by the dimension-5 Weinberg operator [4]. The ultraviolet completion of this can be accomplished in three ways, giving rise to three different versions of the model. These are achieved by introduction of heavy right-handed neutrinos

(type I), scalar triplets (type II), or fermion triplets (type III). Since the Weinberg operator violates lepton number by two units, seesaw implies that neutrinos are Majorana particles.

In the type-I seesaw model, the out-of-equilibrium decays of the heavy Majorana neutrinos can give rise to a lepton asymmetry (leptogenesis) which can be subsequently reprocessed into baryon asymmetry via sphaleron processes [5]. For thermal leptogenesis with a hierarchical spectrum of heavy neutrinos, generation of required CP asymmetry puts a lower bound on the mass of the lightest right-handed neutrino. This is the famous Davidson-Ibarra bound which gives $M_{\text{lightest}} > 10^8\text{--}10^9$ GeV [6]. For the case of the type-III seesaw model, in which the lepton asymmetry is generated from the CP -violating decays of fermion triplets, the lower bound on the CP asymmetry combined with the fact that the $SU(2)$ triplet fermions have additional gauge interactions increases the above bound to $M_{\text{lightest}} > 3 \times 10^{10}$ GeV [7].

These bounds are in tension with the constraints coming from naturalness, implying that the presence of these heavy particles can lead to large quadratic corrections to the Higgs mass, enhancing the hierarchy problem. If we assume naturalness to exist, it puts an upper limit on the mass scale of the seesaw mediators. The naturalness constraints in the context of the minimal type-III seesaw model were studied in Ref. [8]. For instance, assuming that the quadratic corrections to Higgs mass are $< 1 \text{ TeV}^2$, they obtained $M_{\text{lightest}} \leq 3 \times 10^5\text{--}1.84 \times 10^7$ GeV, depending on the parameters of the model. Current searches for such fermion triplets at the ATLAS [9] and CMS [10] detectors at the 13 TeV LHC have set the lower limits on their mass at 790 and 880 GeV at 95% C.L., respectively.

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In the literature, attempts have been made to reduce the bound on the lightest seesaw mediator mass by adding extra scalar fields [11–14]. In particular, Ref. [12] considered a two-Higgs-doublet model (2HDM) in the context of the type-I seesaw, and they showed that the scale of leptogenesis can be lowered and made consistent with the naturalness bound if the second Higgs doublet gets a very small vacuum expectation value (VEV). In this paper, we investigate if a similar phenomenon can happen for the type-III seesaw model. In particular, we examine the role of gauge scatterings and an additional doublet in lowering the leptogenesis scale. In addition, we also include the flavor effects and see the implications of these in lowering the leptogenesis scale further.

The paper is structured as follows. In Sec. II, we discuss the two-Higgs-doublet model with triplet fermions. In Sec. III, we briefly discuss type-III leptogenesis in the usual case with just one Higgs doublet and present the numerical results for the 2HDM case in Sec. IV. We summarize our findings in Sec. V.

II. MODEL

We extend the Standard Model by three hierarchical fermionic $SU(2)$ triplets Σ_i ($i = 1, 2, 3$) with zero hypercharge ($M_{\Sigma_3} \gg M_{\Sigma_2} \gg M_{\Sigma_1}$) and the scalar sector by an additional Higgs doublet with hypercharge one. Thus, our model contains two $SU(2)$ doublets Φ_i ($i = 1, 2$). Since the general 2HDM has tree-level Higgs-mediated flavor-changing neutral currents (FCNCs) (contradicting the experimental bounds on FCNCs), a discrete Z_2 symmetry is imposed on the scalar potential as well as to the fermions to avoid such currents [15]. These FCNCs can also be eliminated if the Z_2 symmetry is softly broken by demanding that each right-handed fermion couples to only one of the doublets [16,17]. Here, we assume that Φ_1 (Φ_2) couples only to u_R^i (Σ_i).

The CP -conserving, softly broken Z_2 symmetric potential in the 2HDM model can be written as

$$\begin{aligned} V(\Phi_1, \Phi_2) = & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) \\ & + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 \\ & + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) \\ & + \frac{\lambda_5}{2} [(\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2], \end{aligned} \quad (1)$$

where all the parameters can be chosen real without loss of generality. The m_{12}^2 term with operator dimension two above softly breaks the Z_2 symmetry.¹

¹This implies that the Z_2 symmetry is respected at short distances, i.e., at distances shorter than the cutoff $1/M$ to all orders in the perturbative series. Thus, in the limit $M^2 \rightarrow \infty$ or, equivalently, at high energies, the Z_2 symmetry is restored [18].

After electroweak symmetry breaking, the scalar fields can be decomposed as

$$\Phi_i = \begin{pmatrix} \phi_i^+ \\ (v_i + \rho_i + i\eta_i)/\sqrt{2} \end{pmatrix}, \quad (2)$$

where v_i are the VEVs given by

$$\begin{aligned} v_1 & \approx \sqrt{\frac{-2m_{11}^2}{\lambda_1}}, \\ v_2 & \approx \frac{1}{1 + \frac{v_1^2}{2m_{22}^2} \lambda^*} \frac{m_{12}^2}{m_{22}^2} v_1, \end{aligned} \quad (3)$$

with $\lambda^* = \lambda_3 + \lambda_4 + \lambda_5$, $m_{11}^2 < 0$, and $v_1^2 + v_2^2 = v^2 \approx (246 \text{ GeV})^2$.

From Eq. (3), we can see that the VEV v_2 can be very small if $m_{12}^2/m_{22}^2 \ll 1$. This is natural in the 't Hooft sense, since, as $m_{12}^2 \rightarrow 0$, a $U(1)$ or Z_2 symmetry is reinstated for $\lambda_5 = 0$ or $\lambda_5 \neq 0$, respectively. Furthermore, the minimization conditions of the scalar potential lead to an upper bound on m_{22}^2 [12]:

$$m_{22}^2 \lesssim \frac{1}{2} m_h^2 \tan^2 \beta \quad (4)$$

in the limit $m_{22}^2 \gg \lambda_2 v_2^2$, provided that m_{11}^2 remains negative. Here, $m_h = 125 \text{ GeV}$ is the Higgs mass, and $\tan \beta = v_1/v_2$.

The part of the Lagrangian concerning neutrino masses and leptogenesis is given by

$$\begin{aligned} \mathcal{L}_\Sigma = & \text{Tr}[\bar{\Sigma}_i i \not{D} \Sigma_i] - \tilde{\Phi}_2^\dagger \bar{\Sigma}_i \sqrt{2} Y_{ai} L_a \\ & - \frac{1}{2} \text{Tr}[\bar{\Sigma}_i M_{\Sigma_i} \Sigma_i] + \text{H.c.}, \end{aligned} \quad (5)$$

where $L_\alpha = (\nu, l)^T$ ($\alpha = e, \mu, \tau$) and $\Phi_2 = (\phi_2^+, \phi_2^0)^T$ are the lepton and Higgs doublet, respectively, with $\tilde{\Phi} \equiv i\sigma_2 \Phi^*$ and $D_\mu \equiv \partial_\mu - i g \tau^A W_\mu^A/2$. The triplets can be represented as

$$\Sigma = \begin{pmatrix} \Sigma^0/\sqrt{2} & \Sigma^+ \\ \Sigma^- & -\Sigma^0/\sqrt{2} \end{pmatrix}, \quad (6)$$

with $\Sigma^\pm = (\Sigma^1 \mp i\Sigma^2)/\sqrt{2}$.

For $m_{12}^2 > 0$ and $\lambda_5 = 0$, the $U(1)$ symmetry is softly broken, and one can write the neutrino mass matrix in terms of the usual type-III seesaw formula with v replaced by v_2 :

$$\mathcal{M}_\nu^{\text{tree}} = -\frac{v_2^2}{2} Y^T \frac{1}{M_\Sigma} Y. \quad (7)$$

The model also has the possibility of generating neutrino masses radiatively if $m_{12} = 0$, $\lambda_5 \neq 0$. In this case, $v_2 = 0$

and the Z_2 symmetry remains unbroken; therefore, neutrinos cannot have masses at tree level. This case corresponds to the type-III scotogenic model, and the neutrino masses are given by [19]

$$\mathcal{M}_{\alpha\beta}^{\text{loop}} = \frac{Y_{ai}Y_{\beta i}\lambda_5 v^2}{32\pi^2 M_{\Sigma_i}} \left[\frac{r}{1-r} \left(\frac{1}{1-r} \right) \ln r \right], \quad (8)$$

where $r = M_{\Sigma_i}^2/M_\phi^2$ and $M_\phi^2 = m_{22}^2 + (\lambda_3 + \lambda_4)v^2$.

For further analysis, we consider $\lambda_5 = 0$, so that neutrino masses are generated only at the tree level. Leptogenesis in the context of a hybrid model of scotogenic type-I and -III seesaw has been discussed in Ref. [20].

III. LEPTOGENESIS

The mass term for the Majorana triplets in Eq. (5) breaks the lepton number, and, in general, the Yukawa coupling matrix $\mathbf{Y}$ is complex and contains new sources of CP violation; therefore, if the decay of the fermion triplet $\Sigma \rightarrow l\Phi_2$ is out of equilibrium, then all the Sakharov conditions [21] are satisfied for the dynamic generation of a lepton asymmetry. This is similar to type-I leptogenesis; however, qualitative differences arise due to the coupling of fermion triplets to SM gauge bosons, as we discuss later.

The lepton asymmetry produced by the decay of a heavy triplet can be written as [22]

$$Y_L \equiv \varepsilon_\Sigma \eta Y_\Sigma(T \gg M_\Sigma) = \frac{135\zeta(3)}{4\pi^4 g_*} \varepsilon_\Sigma \eta, \quad (9)$$

where ε_Σ is the average lepton asymmetry ΔL produced per decay of Σ (CP asymmetry parameter), $Y_a \equiv n_a/s$ is the number density of species a per unit entropy, n_Σ is the total number of triplets, and g_* is the number of active degrees of freedom at $T = M_1$. The efficiency factor $0 < \eta \leq 1$ incorporates the effects of washouts and scatterings and is determined from the numerical solution of the Boltzmann equations for the lepton asymmetry and triplet population. The baryon asymmetry produced from this lepton asymmetry is

$$Y_B = 3CY_{B-L} = -0.0041\varepsilon_\Sigma \eta, \quad (10)$$

where $C = -8/23$ is the sphaleron conversion factor for three flavors and two Higgs doublets [23] and the factor of 3 comes from the three $SU(2)$ degrees of freedom of the triplets.

A. CP asymmetry

For a hierarchical spectrum of heavy seesaw states ($M_{\Sigma_1} < M_{\Sigma_2} < M_{\Sigma_3}$), only the decays of the lightest state are important for leptogenesis, as the asymmetry generated by the heavier states can be washed out by Σ_1 -mediated interactions until they drop out of equilibrium. Considering

the decay of the lightest triplet Σ_1 , CP asymmetry can be calculated as [7]

$$\varepsilon_{\Sigma_1} = \sum_{j=2,3} \frac{3}{2} \frac{M_{\Sigma_1}}{M_{\Sigma_j}} \frac{\Gamma_j}{M_{\Sigma_j}} I_j \frac{V_j - 2S_j}{3}, \quad (11)$$

where

$$I_j = \frac{\text{Im}[(Y^\dagger Y)_{1j}^2]}{|Y^\dagger Y|_{11}|Y^\dagger Y|_{jj}} \quad (12)$$

and S_j and V_j are the loop factors associated with the self-energy and vertex corrections, respectively. Γ_j corresponds to the decay width of triplet components, which is similar to that of a right-handed neutrino and can be expressed as

$$\Gamma_j = \frac{1}{8\pi} M_{\Sigma_j} (Y^\dagger Y)_{jj} (1 - y_j)^2, \quad (13)$$

where the factor $y_j = m_{22}^2/M_{\Sigma_j}^2$ takes the mass of the second Higgs doublet into account. In the type-I seesaw, there is an upper bound on CP asymmetry [6]:

$$\varepsilon_N \leq \frac{3}{8\pi} \frac{M_N}{v^2} (m_{\nu_3} - m_{\nu_1}), \quad (14)$$

which can be translated into a lower bound on the mass of right-handed neutrinos $M_N \gtrsim 6 \times 10^8$ GeV, assuming maximal efficiency, i.e., $\eta = 1$. For triplets, the same expression holds with a factor of $1/3$ and v replaced by v_2 in our model, and we get

$$\varepsilon_\Sigma \leq \frac{1}{8\pi} \frac{M_\Sigma}{v_2^2} (m_{\nu_3} - m_{\nu_1}). \quad (15)$$

Note the inverse dependence on v_2 , which allows us to relax the bounds on maximum allowed CP asymmetry for smaller values of VEVs as shown in Fig. 1. However, unlike the case of the type-I seesaw, we cannot translate this into a lower bound on the triplet mass, as gauge scatterings can suppress the efficiency considerably and we cannot assume $\eta = 1$. Hence, in order to compute the lower bound for the triplet mass, one must solve the coupled Boltzmann equations for the evolution of number density of triplets and lepton asymmetry.

B. Boltzmann equations

Assuming flavor-blind leptogenesis and considering only gauge scatterings mediated by SM gauge bosons, decays (and inverse decays), and $\Delta L = 2$ scatterings, the Boltzmann equations can be written as²

²Since we are concerned with only the decay of the lightest triplet, we suppress the index 1.

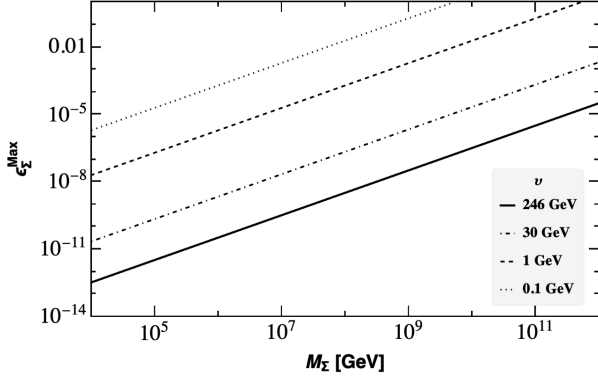


FIG. 1. The maximum bound on CP asymmetry as a function of triplet mass M_{Σ} for different values of v_2 . The solid black line represents the standard case with one Higgs doublet and $v = 246$ GeV.

$$sH z \frac{dY_{\Sigma}}{dz} = -\left(\frac{Y_{\Sigma}}{Y_{\Sigma}^{\text{eq}}} - 1\right)\gamma_D - 2\left(\frac{Y_{\Sigma}^2}{Y_{\Sigma}^{2\text{eq}}} - 1\right)\gamma_A, \quad (16)$$

$$sH z \frac{dY_{B-L}}{dz} = -\gamma_D \epsilon_{\Sigma} \left(\frac{Y_{\Sigma}}{Y_{\Sigma}^{\text{eq}}} - 1\right) - \frac{Y_{B-L}}{Y_l^{\text{eq}}} \left(\frac{\gamma_D}{2} + 2\gamma_{\Sigma}^{\text{sub}}\right), \quad (17)$$

where $Y_a = n_a/s$ is the yield, $z \equiv M_{\Sigma}/T$ is the evolution variable, $H = 1.66\sqrt{g_*}T^2/M_{\text{Pl}}$ is the Hubble expansion rate of the Universe, and $s = 0.44g_*T^3$ is the entropy density. The suffix eq denotes the equilibrium values, and γ 's are the reaction densities for various processes: γ_D is the decay or inverse decay reaction rate, γ_A is the gauge scattering reaction rate for the processes $\Sigma\bar{\Sigma}' \leftrightarrow GG', f\bar{f}, H\bar{H}$, and $\gamma_{\Sigma}^{\text{sub}}$ comes from $\Delta L = 2$ scatterings $LL \leftrightarrow \Phi_2^*\Phi_2^*, L\Phi_2 \leftrightarrow \bar{L}\Phi_2^*$ with the on-shell part subtracted.

The decay and gauge scattering densities are given by [7,24]

$$\gamma_D(z) = \frac{1}{8\pi^3} \frac{M_{\Sigma}^5}{v_2^2} \frac{K_1(z)}{z} \tilde{m}, \quad (18)$$

$$\gamma_A(z) = \frac{M_{\Sigma}^4}{64\pi^4} \int_4^{\infty} dx \frac{\sqrt{x} K_1(z\sqrt{x}) \hat{\sigma}_A(x)}{z}, \quad (19)$$

where $x = s/M_{\Sigma}^2$, s is the center of mass energy squared, and $K_1(z)$ is the modified Bessel function of the first kind. We define

$$\tilde{m} = \frac{(Y^\dagger Y)_{11}}{2M_{\Sigma}} v_2^2 = 8\pi\Gamma_{\Sigma} \frac{v_2^2}{M_{\Sigma}^2} \quad (20)$$

as the effective neutrino mass and use the following expression for the reduced cross section $\hat{\sigma}_A(x)$ [7]:

$$\hat{\sigma}_A(x) = \frac{6g^4}{\pi} \left(1 + \frac{2}{x}\right) r + \frac{2g^4}{\pi} \left[-\left(4 + \frac{17}{x}\right) r + 3\left(1 + \frac{4}{x} - \frac{4}{x^2}\right) \log\left(\frac{1+r}{1-r}\right) \right], \quad (21)$$

with $r = \sqrt{1 - 4/x}$. Note that, since the gauge scatterings do not violate L , they affect only the evolution of Σ population. In the absence of γ_A , the Boltzmann equations take the exact same form as that for right-handed neutrinos, with all the interaction terms multiplied by 3 in Eq. (17). It can be seen that the decay and gauge scattering rates depend only on $\tilde{m}$ and M_{Σ} , which we will take as free parameters for our analysis.³

C. Gauge versus Yukawa interactions

Now, we discuss the effect of gauge reactions which are a function of M_{Σ} only. At high temperatures $z \ll 1$, the gauge reactions are in equilibrium and tend to thermalize the triplet population; thus, no asymmetry can be generated. However, at low temperatures $z \gg 1$, the gauge reaction rate is doubly Boltzmann suppressed and an asymmetry can be generated when these gauge reactions decouple, which implies

$$\frac{\Gamma_A}{H} = \frac{\gamma_A}{n_{\Sigma}^{\text{eq}} H} \lesssim 1, \quad (22)$$

provided that the inverse decays $L\Phi_2 \rightarrow \Sigma$ are also decoupled at that stage. If the inverse decays are still active, the asymmetry will be generated at a later stage once they go out of equilibrium, i.e., $\gamma_D/n_l^{\text{eq}} H \lesssim 1$. To understand the dynamics better, we plot the thermalization rates for various values of M_{Σ} , $\tilde{m}$, and v_2 in Fig. 2.

As can be seen from Fig. 2, the gauge thermalization rate $\Gamma_A/H = \gamma_A/n_{\Sigma}^{\text{eq}} H$ reaches its maximum value at $z = 1$, i.e., $T \sim M_{\Sigma}$, and the lower the triplet mass, the more the triplet is thermalized. As long as $\gamma_A/n_{\Sigma}^{\text{eq}} H > 1$, Y_{Σ} tracks its equilibrium distribution Y_{Σ}^{eq} , and the gauge interactions dominate over the Yukawa ones for $\gamma_D < 4\gamma_A$; i.e., triplets scatter before they decay. The asymmetry is, thus, suppressed by a factor of $\gamma_D/4\gamma_A$ in this period until either $\gamma_A/n_{\Sigma}^{\text{eq}} H < 1$ or $4\gamma_A < \gamma_D$, where the latter condition is used to define $z = z_A$ [22].

For $z > z_A$, the decays dominate over the gauge interactions and the triplets decay before scattering. Hence, there is no suppression from gauge scatterings, even if they are in equilibrium. The only suppression comes from the number of triplets that survive gauge annihilations at $z = z_A$, which is a Boltzmann-suppressed quantity. For the SM VEV (shown in black in Fig. 2), the value of z_A lies around 4 for $M_{\Sigma} = 10^{12}$ GeV and 21 for $M_{\Sigma} = 10^4$ GeV. Notice that the larger the value of z , the lower the number of

³Measuring neutrino masses does not fix $\tilde{m}$ and M_{Σ} , and, hence, they can be treated as free parameters. See Ref. [25].

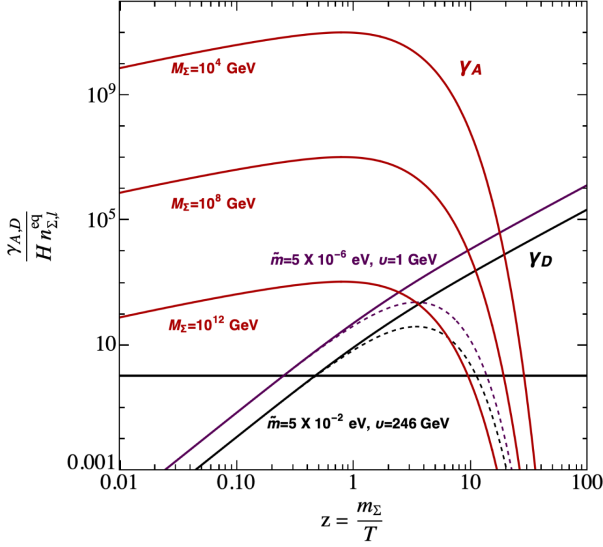


FIG. 2. Evolution of gauge and Yukawa reaction densities relevant for Boltzmann equations. The decay (inverse decay) thermalization rate $\gamma_D/H n_{\Sigma,l}^{\text{eq}}$ denoted by solid (dashed) curves depends only on $\tilde{m}$ for a particular value of VEV [black (purple) lines], whereas the gauge thermalization rates $\gamma_A/H n_{\Sigma,l}^{\text{eq}}$ denoted by red curves depend only on M_{Σ} . The horizontal black line indicates when the reactions go out equilibrium.

triplets left at that point to generate the lepton asymmetry. Therefore, successful low-scale type-III leptogenesis becomes extremely difficult.

The washout from inverse decays, characterized by the thermalization rate $\gamma_D/n_l^{\text{eq}}H$, is depicted by the dashed lines in Fig. 2. The inverse decays are suppressed at $z > 1$, so, if $\gamma_D/n_l^{\text{eq}}H < 1$ at $z = z_A$, the inverse decays are already decoupled and their contribution to washout can be neglected, and the major contribution to washout comes from gauge reactions. This usually holds true for smaller values of $\tilde{m}$, where $\gamma_D/n_l^{\text{eq}}H$ never reaches unity. Larger values of $\tilde{m}$ can keep inverse decays in equilibrium at $z \geq z_A$, where their contribution to the washout starts dominating over that of gauge reactions and the efficiency is suppressed mainly by inverse decays.

Therefore, we can define a minimum value of $\tilde{m} \equiv \tilde{m}_A$ [22] for a given value of M_{Σ} , below (above) which the gauge interactions are more dominant (suppressed) than the Yukawa interactions, allowing us to define the lower $\tilde{m}$ “gauge regime,” where gauge reactions are dominant, and the larger $\tilde{m}$ “Yukawa regime,” where gauge reactions can be neglected; see Fig. 3. In the former, efficiency has a strong dependence on M_{Σ} , and in the latter, efficiency depends entirely on $\tilde{m}$, just like type-I leptogenesis. Maximum efficiency is seen at the juncture of the two regimes. $\tilde{m}_A$ can be computed by requiring

$$\left. \frac{\gamma_D}{n_l^{\text{eq}}H} \right|_{z=z_A} = 1. \quad (23)$$

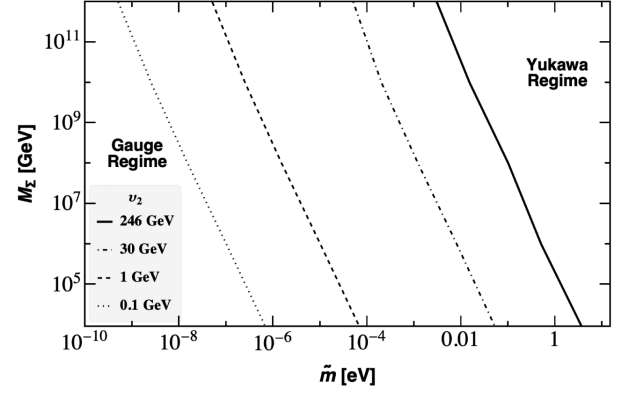


FIG. 3. In the lower (upper) $\tilde{m}$ region, i.e., to the left (right) of each line corresponding to different VEVs, the gauge (Yukawa) interactions are dominant. Lowering the VEV shifts the line delimiting the two regimes toward smaller $\tilde{m}$ values. Hence, with smaller VEVs, unless $\tilde{m}$ is very small, one always remains in the Yukawa regime.

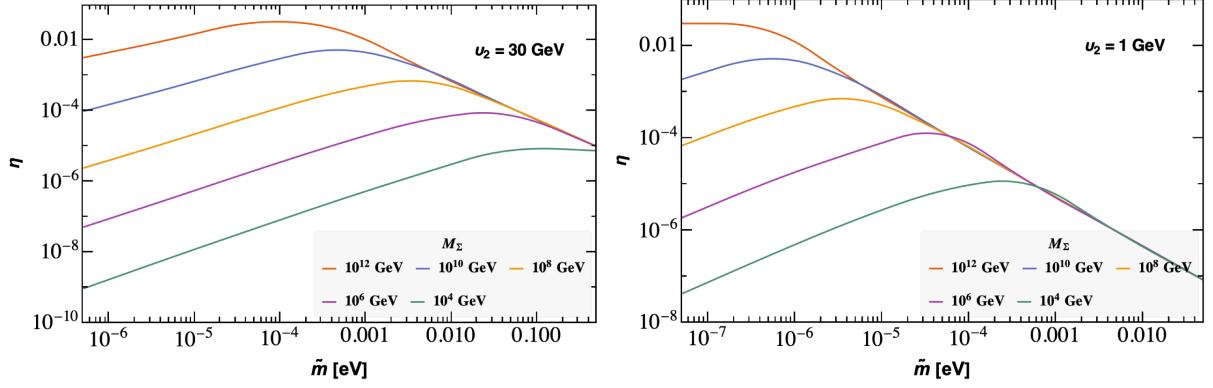
IV. RESULTS AND DISCUSSION

We now discuss the implications for triplet leptogenesis in the presence of another Higgs doublet acquiring a tiny VEV. Note that, for smaller VEVs, even smaller values of $\tilde{m}$ give a higher decay or inverse decay rate than that of the standard VEV with higher $\tilde{m}$; see the purple lines corresponding to $v_2 = 1$ GeV and $\tilde{m} = 0.0005$ eV in Fig. 2. Hence, as one decreases the VEV, the Yukawa interactions start dominating over gauge interaction at smaller values of z for a given value of $\tilde{m}$ and M_{Σ} . While lower values of z_A imply there is more triplet population left to produce the required asymmetry, the inverse decay rate is also enhanced and they can stay in equilibrium for $z > z_A$, suppressing the asymmetry generation. A direct consequence of lowering the VEV is the shifting of $\tilde{m}$ scale separating the two regimes, since the Yukawa interactions are inversely proportional to v_2^2 , as can be seen in Fig. 3.

A. Lower bound on triplet mass

As discussed above, in order put a lower bound on M_{Σ} , we need to determine the efficiency η by numerically solving the Boltzmann equations, Eqs. (16) and (17). Considering only gauge scatterings and $\mathcal{O}(Y^2)$ interactions, the efficiency factor corresponding to different values of $\tilde{m}$, M_{Σ} , and v_2 is shown in Fig. 4. The efficiency reaches a maximum at $\tilde{m} = \tilde{m}_A$ for different values of M_{Σ} . It can be seen that, for higher values of $\tilde{m}$, efficiency is independent of M_{Σ} . We also find that lowering the VEV does not affect the maximum value of efficiency obtained for a particular M_{Σ} significantly; however, maximum efficiency is obtained for lower values of $\tilde{m}$ for smaller values of v_2 .

In conjunction with the maximum bound on CP asymmetry [Eq. (15)], which is relaxed due to lower VEVs, we plot the maximum of $\eta \epsilon_{\Sigma}$ that one can obtain for a particular

FIG. 4. η as a function of $\tilde{m}$ and M_Σ for different values of v_2 .

value of M_Σ and v_2 in Fig. 5. We also show the value of $\eta\epsilon_\Sigma$ required to produce the observed baryon asymmetry from Eq. (10).

Figure 5 is the main result of this work, allowing us to read the lower bound on M_Σ for successful leptogenesis:

$$M_\Sigma > \begin{cases} 3 \times 10^9 \text{ GeV} & \text{for } v_2 = 30 \text{ GeV}, \\ 3 \times 10^7 \text{ GeV} & \text{for } v_2 = 1 \text{ GeV}, \\ 1 \times 10^7 \text{ GeV} & \text{for } v_2 = 0.5 \text{ GeV}. \end{cases} \quad (24)$$

B. Effect of $\Delta L = 2$ scatterings

Now we consider the effect of $\Delta L = 2$ scatterings mediated by right-handed neutrinos [$\mathcal{O}(Y^4)$], proportional to $M_\Sigma^2 \tilde{m}^2 / v^4$, where $\tilde{m}^2 = \sum m_i^2$ is the neutrino mass scale $\gtrsim 0.05$ eV. They contribute to the washout of lepton asymmetry and are significant in the strong washout regime. The thermally averaged $\Delta L = 2$ scattering for $T \lesssim M/3$ in type-I leptogenesis has been approximated in Ref. [26]. For type III, as interactions rates are multiplied by 3, we have

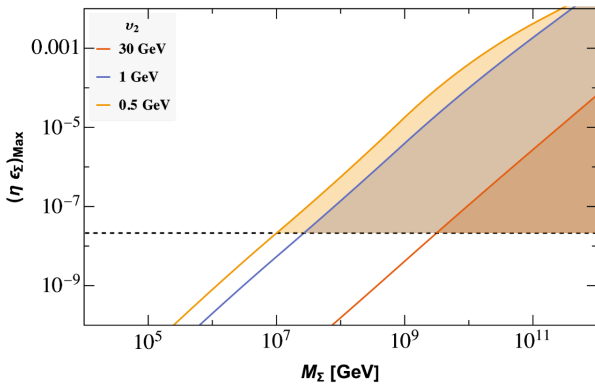


FIG. 5. The maximum value of $\eta\epsilon_\Sigma$ one can obtain for different values of M_Σ and v_2 . The dashed black line indicates the value of $\eta\epsilon_\Sigma = (2.13 \pm 0.06) \times 10^{-8}$ required to produce the observed baryon asymmetry of the Universe in type-III leptogenesis with two Higgs doublets.

$$\frac{\Gamma_{\Delta L=2}}{H} \approx \frac{3T}{2.2 \times 10^{13} \text{ GeV}} \left(\frac{246 \text{ GeV}}{v_2} \right)^4 \left(\frac{\tilde{m}}{0.05 \text{ eV}} \right)^2. \quad (25)$$

At $T \lesssim M_\Sigma/3$, if these scatterings are in equilibrium, they can potentially wash out the asymmetry; however, this depends on the value of $\tilde{m}$. With a hierarchical neutrino mass spectrum and $\tilde{m} \lesssim 10^{-1}$ eV, these scatterings do not affect the final results unless $M_\Sigma > 10^{14}$ GeV. For $\tilde{m} \sim 1$ eV, these scatterings affect the results for $M_\Sigma > 10^{12}$ GeV.

In our model with smaller VEVs, these scatterings become more important than the standard case, as they are inversely proportional to v_2^4 . The corresponding $\tilde{m}$ scale for a particular value of M_Σ above which these scatterings can potentially wash out the generated asymmetry shifts by $(v_2/v)^4$. Since we observe that the maximum efficiency for a given triplet mass is seen toward low $\tilde{m}$ values where these scatterings are not very significant, the bounds in Eq. (24) remain valid. Though, in order to avoid too strong washout, these scatterings must go out of equilibrium at the sphaleron decoupling temperature around $T \sim 100$ GeV, which leads to a lower bound on the VEV of Φ_2 : $v_2 \gtrsim 0.47$ from Eq. (25).

C. Effect of flavor

The flavor blind analysis is justified for decaying right-handed neutrino masses above 10^{12} GeV, as the flavor effects become important at temperatures below this value [27,28]. In the window 10^9 – 10^{12} GeV, the τ Yukawa interactions come into thermal equilibrium and the $B-L$ asymmetry is distributed into L_τ and L_1 ($L_e + L_\mu$). Consequently, the Boltzmann equation for $B-L$ asymmetry [Eq. (17)] splits into Y_{L_τ} and Y_{L_1} , making it a two-flavor problem. For temperatures below 10^9 GeV, the μ Yukawa interactions also come into thermal equilibrium, and one has to consider three independent Boltzmann equations for $Y_{L_{e,\mu,\tau}}$. The flavor-dependent Boltzmann equation for lepton asymmetry can be written as [24]

$$\frac{dY_{\Delta_\alpha}}{dz} = -\frac{\gamma_D}{sHz} \left[\varepsilon_\Sigma^{l_\alpha} \left(\frac{Y_\Sigma}{Y_\Sigma^{\text{eq}}} - 1 \right) + \frac{K_\alpha}{2Y_l^{\text{eq}}} \sum_{\beta=e,\mu,\tau} C_{\alpha\beta}^l Y_{\Delta_\beta} \right], \quad (26)$$

where $\Delta_\alpha = Y_{B/3-L_\alpha}$ and $Y_{L_\alpha} = 2Y_{l_\alpha} + Y_{e_\alpha}$. The flavor projector K_α is defined as

$$K_\alpha = \frac{Y_{\alpha 1}^* Y_{\alpha 1}}{(Y^\dagger Y)_{11}} = \frac{\tilde{m}_\alpha}{\tilde{m}}, \quad (27)$$

and $\sum_\alpha K_\alpha = 1$. The coefficients $C_{\alpha\beta}^l$ couple the differential equations for different flavors and are determined from the reactions that are in equilibrium for a given temperature regime. The total asymmetry is then given by $Y_{B-L} = \sum_\alpha Y_{\Delta_\alpha}$.

Depending on the regime in which asymmetry is generated, the flavor effects allow us to obtain a less suppressed efficiency. Since the suppression in the gauge regime comes from flavor-blind gauge scatterings, flavor effects play a negligible role there. However, they can become important and sizable in the Yukawa regime where inverse decays are important. For example, if

$$\varepsilon_\Sigma^{l_\beta} \gg \varepsilon_\Sigma^{l_\alpha} \quad \text{and} \quad \Gamma(\Sigma \rightarrow L_\alpha H) \gg \Gamma(\Sigma \rightarrow L_\beta H), \quad (28)$$

i.e., the decay of the lightest triplet has a larger CP asymmetry and a smaller decay width in a particular flavor channel, then the suppression from inverse decays in that channel becomes negligible and we obtain a larger asymmetry than the unflavored case.

For simplicity and to illustrate the effect of flavor in both the regimes, we first work in the two-flavor framework (L_τ, L_1) and follow closely the work in Ref. [24]. The flavor coupling matrix is given by [29]

$$C^l = \frac{1}{460} \begin{pmatrix} 196 & -24 \\ -9 & 156 \end{pmatrix}. \quad (29)$$

Assuming $K_1 = 0.99$, $K_\tau = 0.01$ and $\varepsilon_\Sigma^{l_\tau} = 1.1 \times \varepsilon_\Sigma$, $\varepsilon_\Sigma^{l_1} = -0.1 \times \varepsilon_\Sigma$, we show the results for three different values of M_Σ , $v_2 = 1$ GeV in Fig. 6.

For a given value of M_Σ and $\tilde{m}$, it can be seen that in the gauge regime the efficiency is roughly of the same order, whereas, for larger values of $\tilde{m}$, the efficiency can be enhanced by roughly 1–2 orders of magnitude. In conjunction with the bound on CP asymmetry, this allows one to reduce the bound on the triplet mass for successful leptogenesis in Eq. (24) by an order of magnitude. Therefore, for a hierarchical spectrum of decaying fermion triplets, we can lower the scale to roughly $\mathcal{O}(10^6 \text{ GeV})$ for $v_2 \sim 1$ GeV once the flavor effects are included.

We have also checked the results for lower temperature regimes, where the three-flavor decomposition becomes

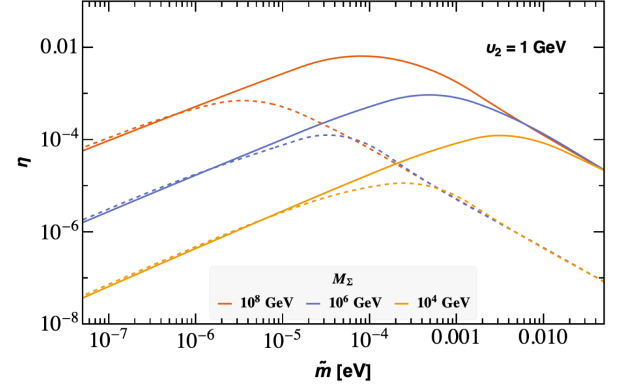


FIG. 6. η as a function of $\tilde{m}$ for the flavor-dependent (solid curves) and flavor-independent (dashed curves) cases assuming $K_1 = 0.99$, $K_\tau = 0.01$ and $\varepsilon_\Sigma^{l_\tau} = 1.1 \times \varepsilon_\Sigma$, $\varepsilon_\Sigma^{l_1} = -0.1 \times \varepsilon_\Sigma$.

important, and found that efficiency enhancement is of the same order as that of the two-flavor case; hence, the bound $M_\Sigma > 10^6 \text{ GeV}$ remains valid for lower temperatures as well, provided that Eq. (28) is satisfied.

D. Naturalness

Having discussed the bound on triplet mass for successful leptogenesis, here we focus our attention toward naturalness, similar to the studies done for type-I and type-III seesaws in the literature [8,12,30], for completeness. The high scale of the seesaw mediators and the presence of heavy particles in the model can provide large corrections to the Higgs mass, thus inducing the naturalness problem [31]. Naturalness demands that these corrections to the Higgs mass are less than $\mathcal{O}(\text{TeV}^2)$ and leads to a conservative naturalness bound on the masses of the heavy particles.

The naturalness constraints in the present setup involve the influence of the mass of the second Higgs doublet (m_{22}) on the first (m_{11}) and the influence of triplet masses (M_Σ) on that of the second Higgs doublet (m_{22}). For the former, the authors of Ref. [12] find that a TeV-scale m_{22} is completely natural. For the latter, the corrections to m_{22}^2 from M_Σ and the naturalness bound can be computed as

$$\delta m_{22}^2 \approx \frac{3}{4\pi^2} \text{Tr}[\mathbf{Y}^\dagger \mathcal{D}_M^2 \mathbf{Y}] < \Lambda^2, \quad (30)$$

where Λ denotes the scale of corrections and $\mathcal{D}_M = \text{diag}(M_{\Sigma_1}, M_{\Sigma_2}, M_{\Sigma_3})$. Similar to the correction to the Higgs mass, if we take $\Lambda = 1 \text{ TeV}$, the bound on M_Σ becomes [30]

$$M_\Sigma \lesssim 3.5 \times 10^7 \text{ GeV} \left(\frac{v_2}{246} \right)^{2/3}. \quad (31)$$

However, since the heavy scalars are expected to be at the TeV scale, one can consider slightly relaxed bounds by

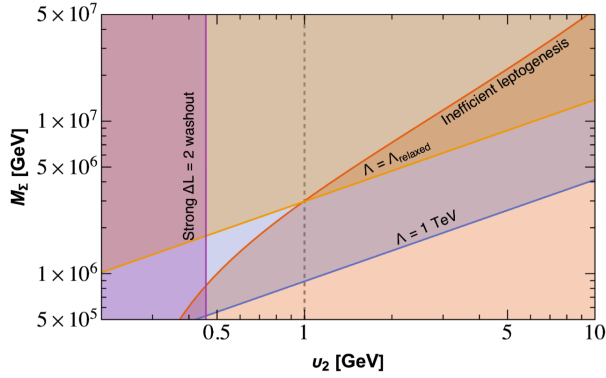


FIG. 7. Bounds on the triplet mass and the VEV of the second Higgs doublet from naturalness (blue and orange curves) and leptogenesis (red and purple curves). The gray dashed line indicates the value of v_2 below which the leptogenesis and relaxed naturalness bounds become compatible.

taking the correction to m_{22}^2 larger than $\mathcal{O}(\text{TeV}^2)$. Following the approach in Ref. [12], we take this relaxed bound $\Lambda_{\text{relaxed}} = \min(10 \text{ TeV}, 10\sqrt{m_h^2 \tan \beta/2})$, where we have used the consistency condition of Eq. (4). In Fig. 7, we show these bounds in conjunction with the bounds from efficient leptogenesis including the flavor effects (a similar figure is provided for the type-I seesaw in Ref. [12]).

It can be seen that demanding corrections to m_{22}^2 of the order TeV is not consistent with the bounds from leptogenesis despite the lowering of the leptogenesis scale via the small VEV of the second Higgs doublet. However, if we consider the relaxed naturalness bound, a small region of parameter space opens up that is compatible with all the bounds (as indicated by the light blue region in the center) and can be read as $8.5 \times 10^5 \text{ GeV} < M_\Sigma < 3 \times 10^6 \text{ GeV}$. Furthermore, in order for these bounds to be compatible, v_2 should be less than 1 GeV (indicated by the dashed gray line).

V. SUMMARY

We considered the possibility of lowering the scale of type-III leptogenesis in the context of a two-Higgs-doublet model. Because of the gauge scatterings of fermion triplets, the asymmetry produced by their decays is highly suppressed, and successful leptogenesis is possible only for

$M_\Sigma > 3 \times 10^{10} \text{ GeV}$ (10^9 GeV) in single-flavor (three-flavor) approximation [7]. We have shown that if all the heavy fermion triplets couple only to the second Higgs doublet, which obtains a small VEV, the absolute bound on the mass of the decaying triplet can be lowered to 10^7 GeV , roughly 3 orders of magnitude below the current bounds for successful leptogenesis. We also include the flavor effects and find that there can be a further reduction to 10^6 GeV .

We have studied the interplay of gauge and Yukawa scatterings for different VEVs and find the latter get more dominant than the former as one decreases the VEV. This results in shifting the scale delimiting the gauge and Yukawa regimes, at the junction of which one obtains maximum efficiency. Hence, the efficiency is enhanced for smaller values of $\tilde{m}$ compared to the standard case with just the SM Higgs doublet. This suggests that, for lower VEVs, the effective contribution to the lightest active neutrino needs to be smaller. Higher values of $\tilde{m}$ lead to significant suppression from inverse decays, and the dynamics become similar to the case of type-I leptogenesis. The lowering of the scale is mostly due to the relaxation of Davidson-Ibarra bound on the maximum CP asymmetry for smaller VEVs. Furthermore, constraints from strong washout due to $\Delta L = 2$ scatterings and compatibility with naturalness impose stringent constraints on the VEV of the second Higgs doublet to be $0.47 \text{ GeV} \lesssim v_2 \lesssim 1 \text{ GeV}$.

The analysis above highlights that unlike the case of type-I leptogenesis, where the scale can be lowered to TeV scales (and below) in the presence of a second Higgs doublet, there exists only a small region of parameter space around $M_\Sigma \sim 10^6 \text{ GeV}$ that is compatible with both leptogenesis and naturalness.

ACKNOWLEDGMENTS

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Part III: Resumen

RESUMEN DE LA TESIS

El modelo estándar de la física de partículas (SM), con tres generaciones de quarks y leptones, constituye la teoría más completa para explicar las interacciones entre las partículas elementales y tres de las fuerzas fundamentales la fuerte, la débil y la electromagnética, y ha tenido un éxito notable en la descripción de los resultados experimentales. Sin embargo, está lejos de ser una teoría completa, ya que existe evidencia convincente de la existencia de masas de los neutrinos, del exceso de materia bariónica sobre la materia antibariónica, y de la abundancia de materia oscura:

- **Masas de los neutrinos:** Los neutrinos, los únicos fermiones neutros en el modelo estándar, se consideran sin masa. En las últimas décadas, varios experimentos con neutrinos solares, atmosféricos y de reactores han confirmado el fenómeno de las oscilaciones de neutrinos, lo cual indica una mezcla entre los autoestados de sabor y de masa [57–60]. Estas mediciones sugieren que las masas de los neutrinos son seis órdenes de magnitud menores que la masa del electrón, lo que convierte la observación de masas de neutrinos no solo en una indicación directa de física BSM, sino que también nos lleva a preguntarnos: ¿Por qué las masas de los neutrinos son tan pequeñas?
- **Asimetría bariónica del universo:** Prácticamente todo lo que vemos e interactuamos está compuesto de materia, mientras que la antimateria sigue siendo esquiva. Esta asimetría entre materia y antimateria en el universo ha sido medida con gran precisión, siendo del orden de 10^{-10} , a partir de la Nucleosíntesis del Big Bang y de mediciones del Fondo Cósmico de Microondas [143, 355, 356]. Para generar una asimetría bariónica neta (bariogénesis), deben cumplirse las tres condiciones de Sakharov: violación del número bariónico, violación de CP y decaimientos fuera del equilibrio en el universo temprano [148]. Aunque estas tres condiciones están presentes dentro del SM, no pueden explicar por sí solas la gran asimetría observada.

- **Materia oscura:** Existe una evidencia sólida, en forma de observaciones de curvas de rotación galácticas [215,216], cúmulos de galaxias y lentes gravitacionales [223–226,357,358], que indica que la mayor parte de la materia en el universo observable es no luminosa y no bariónica. Esta forma de materia, denominada materia oscura (DM), contribuye con aproximadamente un 25% de la densidad total de energía del universo, es decir, alrededor de cinco veces más que la materia visible. La identidad de la materia oscura aún sigue siendo un misterio. Los neutrinos activos del SM, las únicas partículas del SM que podrían ser candidatos viables a DM, no pueden explicar la densidad relicta observada de DM debido a las restricciones experimentales sobre su masa.

Para poder explicar estas observaciones, es necesario extender el SM. Tales extensiones se conocen como teorías Más Allá del Modelo Estándar (BSM), cuyo objetivo es resolver uno o varios de los problemas no resueltos del SM. En el pasado, se han propuesto muchas soluciones para abordar conjuntamente una o más de estas cuestiones dentro de un marco común. A continuación, se hace una breve revisión de algunos de los escenarios BSM más comunes.

- **Mecanismo seesaw:** La explicación más elegante para las pequeñas masas de los neutrinos es a través del mecanismo seesaw, vinculado al operador de dimensión 5 $LLHH$, conocido como el operador de Weinberg [55], donde L y H son, respectivamente, el doblete leptónico y el doblete de Higgs del SM. En el escenario más sencillo, conocido como seesaw de tipo I, se añaden al SM neutrinos pesados de mano derecha tipo Majorana N_i con masa M_i [93,94]. Tras la ruptura espontánea de la simetría electrodébil, el operador de Weinberg genera masas diminutas para los neutrinos: $m_\nu \propto v_H^2/M_i$. Sin embargo, dado el valor del vacío del Higgs (VEV), masas de neutrinos del orden sub-eV pueden generarse para $M \sim \mathcal{O}(10^{15})$ GeV, asumiendo acoplamientos de $\mathcal{O}(1)$. Esta escala tan alta resulta difícil de explorar en experimentos terrestres.
- **Leptogénesis:** Entre los diversos mecanismos propuestos para la bariogénesis, la leptogénesis ha suscitado gran interés y popularidad debido a su conexión natural con el mecanismo seesaw [161]. En el mecanismo seesaw de tipo I, el número leptónico se viola debido a la naturaleza Majorana de los RHNs. Los decaimientos con violación de CP de estos RHNs pueden generar una asimetría leptónica neta que posteriormente es transformada en una asimetría bariónica mediante los esfalerones electrodébiles. La leptogénesis térmica ha sido ampliamente discutida en la literatura, y una revisión al respecto puede encontrarse en [162]. Dado que estos mecanismos están vinculados a la escala seesaw, también resultan difíciles de poner a prueba.
- **WIMPs:** Las partículas masivas que interactúan débilmente (Weakly Interacting Massive Particles) son uno de los candidatos a materia oscura más comúnmente considerados. Estas partículas a escala débil están inicialmente en equilibrio térmico con el baño térmico del SM y participan en interacciones a escala

débil con un valor aproximado de $\langle\sigma v\rangle \sim 10^{-9} \text{ GeV}^{-2}$. Cuando la tasa de interacción disminuye por debajo de la tasa de Hubble del universo, su distribución se desvía de la distribución térmica y las interacciones se congelan (freeze-out). El rendimiento congelado puede reproducir la abundancia relicta observada. Sin embargo, hasta ahora ninguno de los experimentos de búsqueda directa o indirecta de materia oscura ha proporcionado evidencia convincente de la existencia de WIMPs.

10.1 Objetivos

Observamos que las extensiones mínimas propuestas son o bien muy difíciles de poner a prueba o no han proporcionado ninguna señal en las búsquedas experimentales. Por ello, los objetivos de esta tesis se centran en la fenomenología de modelos BSM que puedan ser testeados en un futuro cercano, que evadan las restricciones existentes y que generen predicciones comprobables a través de observables astrofísicos, cosmológicos o en colisionadores.

Debido a los resultados nulos obtenidos en diversos experimentos en la búsqueda de nueva física, resulta razonable inferir que la nueva física podría estar débilmente acoplada al Modelo Estándar o residir en escalas de energía elevadas. Una rápida revisión de los títulos de los artículos incluidos sugiere un enfoque en extensiones BSM que van más allá de marcos mínimos. Este enfoque contrasta con el principio de la navaja de Occam. No obstante, la mayoría de los trabajos presentados aquí abordan estas cuestiones dentro de marcos unificados. Aunque estos marcos no son estrictamente mínimos, resultan interesantes al establecer conexiones entre la física de neutrinos y el sector oscuro, manteniendo a la vez la elegancia de la extensión mínima sobre la cual fueron construidos.

Los objetivos clave de este trabajo se resumen entonces de la siguiente manera:

- Estudiar modelos populares de nueva física que expliquen los problemas de las masas de neutrinos, la asimetría bariónica y la materia oscura mediante la extensión del SM.
- Proponer nuevas extensiones del Modelo Estándar para explicar estos problemas, ya sea de forma separada o dentro de un marco común.
- Analizar y construir teorías efectivas de campo que generen la abundancia relicta correcta y las masas de neutrinos adecuadas.
- Estudiar la fenomenología de los modelos y mostrar el espacio de parámetros viable que pueda ser probado en experimentos directos e indirectos de neutrinos y materia oscura, tanto en curso como propuestos.
- Investigar las implicaciones para señales astrofísicas y cosmológicas, así como las búsquedas en colisionadores.

10.2 Metodología

En general, tales extensiones implican ampliar el contenido de partículas o simetrías del SM para incluir nuevas interacciones, nuevos acoplamientos y nuevos números cuánticos que conformen un modelo BSM. Sin embargo, esto no debe perjudicar el SM que conocemos y que ya ha sido probado. Con esta idea de mantener intactas las predicciones principales del SM, se tiene cierta libertad creativa para implementar nuevas ideas y marcos conceptuales con un enfoque de apertura hacia nuevas posibilidades. El objetivo final no es solo crear modelos que puedan resolver los desafíos planteados, sino también aquellos que puedan ser probados con nuestras capacidades experimentales actuales o, al menos, explorados en el futuro. Por tanto, la sinergia entre teóricos y experimentalistas es vital en esta búsqueda de nueva física. Identificar la(s) extensión(es) correcta(s) que nos acerquen a una comprensión más completa de la naturaleza sigue siendo una tarea desafiante, aunque el conjunto finito de resultados experimentales sirve como una guía para los constructores de modelos. A continuación, examinaremos los pasos usuales que se deben seguir para construir un modelo BSM.

Una receta para cocinar un modelo BSM

Habiendo establecido el contexto, es momento de discutir el enfoque y camino que siguen los constructores de modelos BSM al extender el SM. Veamos cómo cocinar un modelo BSM:

1. **[Ingrediente principal] Seleccionar un problema a resolver:** Hemos visto varias deficiencias del SM. El primer paso es enfocarse en uno o más de estos problemas, ya sea que provengan de fenómenos inexplicados o inconsistencias teóricas. Se puede elegir abordar múltiples problemas dentro de un mismo modelo. Esto es análogo a decidir qué platillo quieres cocinar y comer.
2. **[Cocina] Elegir el enfoque:** Esto sirve como principio guía para construir el modelo, similar a escoger una cocina — italiana, española, india, etc. El enfoque más común se basa en argumentos de simetría, dado que el SM mismo está construido sobre principios de simetría. Ampliar el grupo de simetría puede conducir a esquemas de unificación o clases enteras de modelos, especialmente dentro de marcos supersimétricos. Alternativamente, se puede adoptar una perspectiva de abajo hacia arriba usando el marco EFT. Otras ideas guía incluyen la *minimalidad* y la *naturalidad*, dependiendo de tu preferencia.
3. **[Ingredientes] Añadir nuevas partículas; extender las simetrías:** El siguiente paso es ampliar el contenido de partículas y el grupo de simetría del SM. Se pueden añadir nuevos fermiones o escalares; las simetrías pueden ser continuas (locales o globales) o discretas. Los bosones gauge serán determinados por si la simetría es gauged o no.
4. **[Cocción I] Construir el lagrangiano:** Una vez hechas las adiciones, es necesario

- escribir todos los términos renormalizables permitidos por las simetrías (tanto las antiguas como las nuevas). Si se usa el enfoque EFT, también habrá que incluir operadores de dimensión superior. Luego, hay que decidir si se desea que la simetría se rompa espontáneamente (es decir, nuevos bosones gauge, el nuevo campo escalar adquiere un VEV) o que permanezca intacta, dependiendo del problema que se intente resolver.
5. **[Degustar y ajustar] Comprobación de consistencia:** Con el modelo en mano, es esencial verificar la consistencia teórica. Esto incluye restricciones por perturbatividad, unitariedad, cancelación de anomalías, estabilidad del vacío y sensibilidad al ajuste fino (si la naturalidad es un principio guía). Si surgen problemas, puede ser necesario ajustar los ingredientes — quizá añadir más ‘sal’ — para corregirlos.
 6. **[Cocción II] Fenomenología:** En esta etapa, el modelo está sólo medio cocinado. Ahora hay que explorar sus implicaciones fenomenológicas. Esto implica estudiar nuevas interacciones (dispersión, decaimientos, etc.), calcular el espectro de masas y realizar cualquier diagonalización necesaria, además de calcular contribuciones de diagramas en lazo.
 7. **[Servir] Predicciones:** El paso final es ‘servir’ el modelo, primero restringiéndolo con los límites experimentales (sensibilidades actuales y futuras), por ejemplo: pruebas de precisión electrodébil, límites de astrofísica y cosmología, para mapear el espacio de parámetros permitido. Se deben hacer predicciones sobre las firmas únicas del modelo que puedan ser probadas con el montaje actual o en el futuro.

Este flujo general es comúnmente seguido en la construcción de modelos BSM y también se ha empleado en el desarrollo de los modelos presentados en la Parte II de esta tesis.

Herramientas computacionales

Para trabajar la fenomenología de los modelos presentados en este trabajo, hemos utilizado diversas herramientas de física de altas energías (HEP). Usamos SARAH [359] y FeynRules [360] para construir y analizar algunos de estos modelos. El análisis numérico, especialmente la resolución de las ecuaciones de Boltzmann, se realiza con la ayuda de Mathematica. En el caso de los modelos de materia oscura, la abundancia relicta se calcula con herramientas como micrOMEGAs [361].

Una vez que un modelo ha sido implementado en estas herramientas, obtenemos los vértices de interacción. Luego utilizamos este resultado para generar diagramas de Feynman para un proceso a través de FeynArts [362]. La sección eficaz relevante o la tasa de decaimiento para dicho proceso se calcula con la ayuda de FeynCalc [363]. Frecuentemente, los resultados de los cálculos se expresan en términos de integrales de Passarino-Veltman, las cuales se descomponen utilizando Package-X [364].

La sección eficaz de producción y las firmas en colisionadores para algunas de las nuevas partículas que aparecen en nuestras extensiones se analizan mediante MadGraph [350] y CalcHEP [365].

10.3 Resultados

10.3.1 Sector de neutrinos

En el Capítulo 2, discutimos los mecanismos de masa de neutrinos y los modelos estándar de seesaw. También abordamos las variantes a baja escala del seesaw tipo-I, junto con algunos modelos radiativos de masa de neutrinos. La posibilidad de testear la violación del número leptónico (LNV) en estos seesaw estándar de alta escala sigue siendo un problema.

La escala puede ser suprimida paramétricamente si existen nuevos escalares en la teoría, lo que a su vez podría conducir a nuevos operadores tipo Weinberg y modelos de seesaw [6, 7]. Además, estos seesaw también pueden explicar la asimetría bariónica a través de leptogénesis, como se discutió en el Capítulo 3. Se analizaron los casos estándar de leptogénesis, así como enfoques que van más allá del método canónico. El trabajo en Ref. [3] es un ejemplo de una aproximación que busca reducir la escala de la leptogénesis tipo-III.

Nuevos operadores tipo Weinberg y modelos seesaw

En las referencias [6, 7] investigamos nuevos operadores tipo Weinberg que involucran uno o dos multipletes escalares SU(2) adicionales, extendiéndose hasta la representación quintuplete. Estos operadores tienen el potencial de generar masas de neutrinos de Majorana tras la ruptura espontánea de la simetría electrodébil (EWSSB), conduciendo así a nuevos modelos seesaw que corresponden a completaciones UV a nivel árbol de estos operadores efectivos. Estos escenarios se muestran en la Tabla 10.1.

A medida que estos nuevos escalares adquieren valores de Expectación del Vacío (VEVs), se espera que permanezcan cercanos a la escala electrodébil. Además, el parámetro ρ restringe estos nuevos VEVs a ser del orden de $\sim O(1)$ GeV. Analizamos los potenciales escalares para evaluar una posible supresión natural de estos nuevos VEVs y discutimos sus implicaciones en relación con los pseudo-Goldstone bosones.

Adicionalmente, establecemos restricciones generales sobre los nuevos multipletes escalares a partir de correcciones de bucle a las pruebas de precisión electrodébiles (EWPTs) y mediante búsquedas directas de escalares con carga doble en colisionadores (discutido en la Sección 2.5.3), ver Fig. 10.1.

Asimismo, exploramos la posibilidad de explicar la medida del W-bosón reportada

Model	Scalar Multiplets	Mediators	Wilson Coefficients
A ₁	$\Phi_1 = 4_{-1/2}^S$	$\Sigma = 5_0^F$	$C_5^{(2)} = y_1 M_{\Sigma}^{-1} y_1^T$
A ₂	$\Phi_1 = 4_{-3/2}^S$	$\mathcal{F} = 3_{-1}^F$	$C_5^{(1)} = y_1 M_{\mathcal{F}}^{-1} y_H^T + y_H M_{\mathcal{F}}^{-1} y_1^T$
B ₁	$\Phi_1 = 4_{1/2}^S, \Phi_2 = 4_{-3/2}^S$	$\mathcal{F} = 5_{-1}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B ₂	$\Phi_1 = 3_0^S, \Phi_2 = 5_{-1}^S$	$\mathcal{F} = 4_{-1/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B ₃	$\Phi_1 = 5_{-2}^S, \Phi_2 = 5_1^S$	$\mathcal{F} = 4_{3/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$
B ₄	$\Phi_1 = 5_{-1}^S, \Phi_2 = 5_0^S$	$\mathcal{F} = 4_{1/2}^F$	$C_5^{(3)} = y_1 M_{\mathcal{F}}^{-1} y_2^T + y_2 M_{\mathcal{F}}^{-1} y_1^T$

Table 10.1: Lista de modelos genuinos obtenidos a partir de nuevos operadores tipo Weinberg con un nuevo multiplete escalar Φ_1 (A_i) y con dos nuevos multipletes escalares $\Phi_{1,2}$ (B_i) [7]. La completación UV correspondiente a fermiones de Majorana o fermiones vectoriales se muestra en la columna de mediadores. Los coeficientes de Wilson en función de la masa del mediador y los acoplamientos de Yukawa $y_1, 2, H$ con los escalares $\Phi_{1,2}, H$ se presentan en la última columna. Los multipletes se representan como $N_Y^{S,F}$, donde N denota la representación de $SU(2)$, S, F indican la naturaleza escalar o fermiónica del multiplete, y Y su hipercarga.

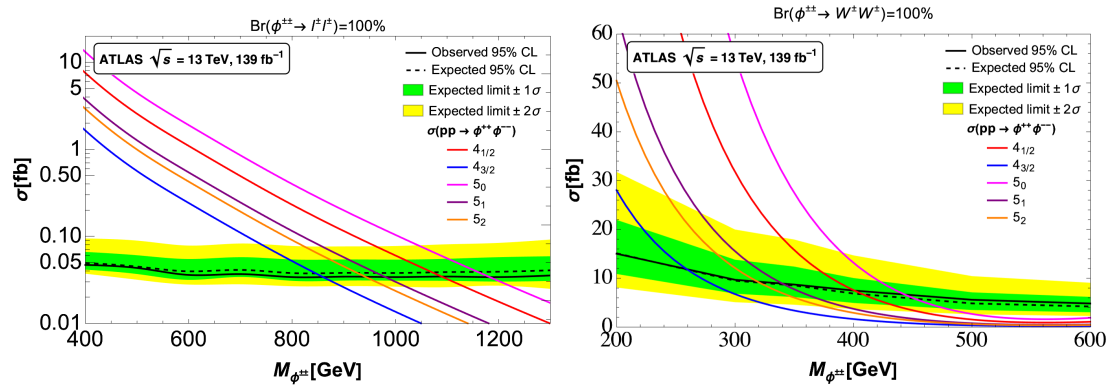


Figure 10.1: Límites superiores sobre la sección eficaz de producción de escalares doblemente cargados multiplicada por su razón de decaimiento en dileptones del mismo signo (izquierda) y en un par de bosones W (derecha). La región por encima del límite observado está excluida. Las bandas verde y amarilla representan las curvas de exclusión esperadas dentro de $1 - 2\sigma$. [136, 352].

por CDF [351].¹

Calculamos la contribución a las masas de neutrinos tanto a nivel árbol como a un lazo, y discutimos variantes a baja escala basadas en seesaws inversos o asumiendo jerarquías entre los dos acoplamientos de Yukawa diferentes. Dependiendo de estas jerarquías, es posible obtener una fenomenología rica proveniente de los multipletes con los acoplamientos de Yukawa mayores. También examinamos la generación de operadores de dimensión 6 y exploramos sus implicaciones, tales como la no unitariedad de la matriz de mezcla leptónica, interacciones no universales del bosón Z y violación de la conservación de sabor leptónico (LFV).

Además, en presencia de simetrías específicas, construimos modelos escotogénicos o modelos escotogénicos generalizados que no solo pueden producir candidatos viables

¹ Mediciones más recientes de la masa del W son consistentes con resultados anteriores; por lo tanto, la anomalía de CDF podría ser errónea. En cualquier caso, ya sea que la anomalía desaparezca o persista, abordarla no fue el objetivo principal de nuestro trabajo.

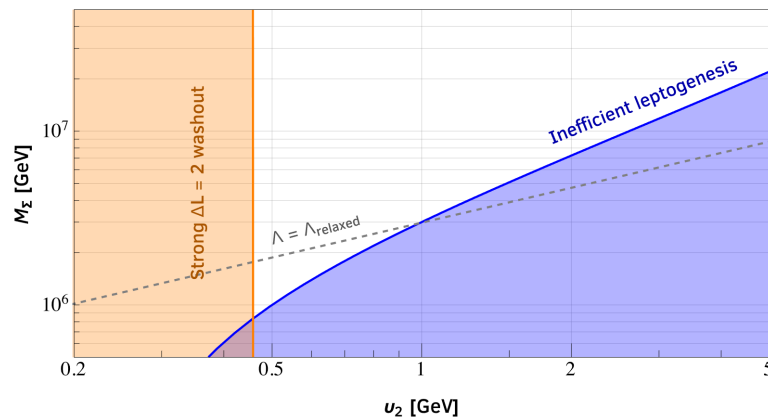


Figure 10.2: El espacio de parámetros viable para una leptogénesis exitosa en un modelo de balancín tipo III, donde los fermiones tripletes Σ se acoplan a un segundo doblete de Higgs Φ_2 con VEV v_2 . Véase la referencia [3] para más detalles.

a materia oscura (DM), sino también generar masas de neutrinos. Para satisfacer la abundancia relicta de DM, proporcionamos rangos de masa para los nuevos multipletes escalares y fermiónicos. Al ser modelos de seesaw, también podemos incorporar leptogénesis, obteniendo nuevamente un modelo común para los tres problemas. Además, la presencia de nuevos escalares podría conducir directamente a la bariogénesis electrodébil.

Leptogénesis de tripletes a escalas inferiores

Vimos en el Capítulo 3 que el mecanismo canónico de leptogénesis tipo-I requiere que las masas de los singletes fermiónicos M_N sean del orden de $O(10^9)$ GeV. Sin embargo, para los tripletes fermiónicos (seesaw tipo-III), este umbral aumenta a $M_\Sigma O(10^1)$ [168]. El requisito más alto para los tripletes se debe a la presencia de interacciones gauge, que compiten con las interacciones Yukawa, causando una disminución en su población.

Aunque varios trabajos demuestran la plausibilidad de reducir la escala de leptogénesis tipo-I hasta $O(\text{TeV})$ mediante la extensión del sector escalar, por ejemplo, en modelos de dos dobles de Higgs o modelos escotogénicos, esta posibilidad no había sido explorada para la leptogénesis tipo-III. En Ref. [3], investigamos si es posible bajar la escala de leptogénesis de tripletes fermiónicos cuando se acoplan a un segundo doble de Higgs Φ_2 , que asume un VEV mucho menor comparado con el del SM.

Demostramos que, efectivamente, es posible reducir la escala hasta $O(10^6)$ GeV, cuando se incluyen los efectos de sabor. Esto se atribuye a las interacciones Yukawa mejoradas con el segundo doble. También discutimos la compatibilidad con la naturalidad y las restricciones por el lavado fuerte. Los resultados se resumen en la Fig.10.2. La región sobre la línea discontinua corresponde a la restricción de naturalidad, donde la corrección a la masa de Φ_2 por M_Σ es mayor que $\Lambda_{\text{relaxed}} = 10\text{TeV}$. La región azul sombreada está excluida por leptogénesis ineficiente, mientras que el área naranja representa la región con posible lavado fuerte $\Delta L = 2$.

Este estudio, aunque realizado en el contexto del seesaw tipo-III, puede extenderse para explorar leptogénesis en los nuevos modelos de seesaw discutidos anteriormente debido a características comunes: *i*) los nuevos multipletes tienen interacciones gauge; *ii*) los nuevos VEVs son comparativamente más pequeños, lo que conduce a Yukawas más grandes.

10.3.2 Sector oscuro

En el capítulo 4, analizamos que la materia oscura constituye aproximadamente el 25% del presupuesto energético total del universo. Dada esta proximidad entre las densidades de materia oscura y bariónica, es plausible suponer que el sector oscuro podría ser bastante rico, al estar compuesto por múltiples partículas estables. Además, la abundancia de reliquias de materia oscura podría estar determinada por una asimetría inicial similar a la bariónica. Esta idea de tener materia oscura multicomponente y asimétrica motivó los trabajos en las referencias [2,4].

Componentes múltiples simétricos y asimétricos

En la referencia [4], consideramos un sector oscuro extendido compuesto por varios componentes, que abarca candidatos a DM simétricos, parcialmente asimétricos y completamente asimétricos. En presencia de múltiples componentes estables, se requiere una combinación de simetrías estabilizadoras que garantice al menos n componentes de DM estables. Con una sola simetría discreta $\mathbb{Z}_N$, un DM multicomponente requiere $N \geq 4$. Resulta más conveniente utilizar combinaciones de simetrías discretas y continuas para lograr un marco multicomponente. Demostramos que, para un producto de n simetrías discretas $\mathbb{Z}_N$, se permiten al menos n componentes de DM estables, pero pueden admitir componentes más estables dependiendo de la cinemática. Luego discutimos la correlación entre las asimetrías en los sectores visibles y oscuros, dependiendo de las cargas conservadas en cada sector y el operador de mayor dimensión que transfiere las asimetrías.

Además, como se mencionó, esta configuración, con más interacciones que el enfoque WIMP convencional, genera interesantes características fenomenológicas. Investigamos el papel de las conversiones entre partículas en el sector oscuro, mostrando cómo estas conversiones tienden a alterar la abundancia de los componentes más pesados y ligeros en comparación con escenarios con solo aniquilaciones. También exploramos por primera vez el papel de las conversiones en presencia de asimetrías iniciales en el sector oscuro. Nuestros hallazgos muestran que las conversiones tienden a hacer que los componentes más pesados de la materia oscura sean altamente asimétricos; por el contrario, los componentes más ligeros exhiben un comportamiento predominantemente simétrico.

También realizamos una parametrización efectiva considerando mediadores pesados/ligeros para determinar el escalamiento de la ley de potencia para el rendimiento de materia

Model	σ	Symmetric		Partially Symmetric
		Ann.	Ann. + Conv.	Ann. / Ann. + Conv.
Mass-independent	$\neq f(m_i)$	$\sim m^0$	$\sim m^{-1}$	$\sim \eta m^{[1,0]}$
Heavy mediators	$\sim m_i^2/\Lambda^4$	$\sim m^{-2}$	$\sim m^{-[3,2]}$	$\sim \eta m^{[1,-2]}$
Light mediators	$\sim 1/m_i^2$	$\sim m^2$	$\sim m^{[0,2]}$	$\sim \eta m^{[1,2]}$

Table 10.2: Escalamiento de la abundancia de reliquias con la masa de materia oscura para todos los componentes, excepto el más ligero, en una materia oscura multicomponente, donde β denota la asimetría del sector oscuro. Los exponentes izquierdo/derecho entre corchetes corresponden a conversiones fuertes/débiles (y aniquilaciones) en el caso simétrico (parcialmente asimétrico).

oscura en función de su masa. Este marco debería ser aplicable a una amplia gama de modelos, siempre que las masas/secciones eficaces se encuentren en la escala débil (véase la Tabla 10.2).

Se puede observar que, en varios casos, $\Omega \sim m^0$. En estos casos, los componentes de DM en presencia de conversiones pueden tener una contribución similar a la abundancia total de reliquias. El efecto de las conversiones en componentes simétricos y asimétricos muestra que las abundancias pueden reducirse, y el sector oscuro podría permanecer oculto al alcance experimental actual.

Un modelo de cogénesis

En la Ref. [2] se introduce un modelo de cogénesis (véase la Sección 4.4.3) para generar simultáneamente una asimetría en los sectores visible y oscuro. Este modelo amplía el mecanismo de leptogénesis térmica de dos sectores, ampliando el sector oscuro para abarcar dos componentes distintos de materia oscura: un escalar S y un fermión ψ . En varios escenarios dentro del modelo, uno o ambos componentes de materia oscura presentan asimetría.

La motivación fue desarrollar un mecanismo de congelamiento asimétrico. Para ello, propusimos una partícula de interacción débil que sirve como posible candidata a materia oscura asimétrica producida en las desintegraciones de un fermión χ a escala de TeV que es un WIMP asimétrico. Sin embargo, que los productos de desintegración también presenten una asimetría depende de la fuerza del acoplamiento de congelamiento. Esto se ilustra en la Fig. 10.3. Si los acoplamientos débiles son comparativamente más fuertes que la producción de congelación de las desintegraciones, que ocurre antes de que χ se haya congelado, entonces la población de DM resultante (ψ) es simétrica. Por otro lado, si el acoplamiento débil es comparativamente menor, entonces la producción dominante ocurre a partir de las desintegraciones tardías, una vez que χ se ha congelado asimétricamente. En este último caso, los productos de desintegración también son asimétricos.

Por lo tanto, las desintegraciones tardías de la partícula progenitora χ desempeñan un papel fundamental en la determinación de la naturaleza asintótica y la producción del componente de materia oscura. Estas desintegraciones tardías también pueden poblar

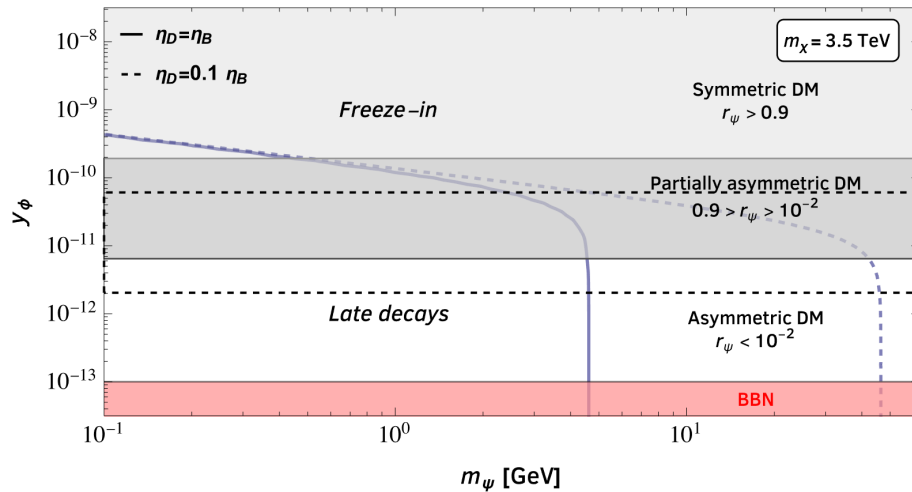


Figure 10.3: El espacio de parámetros para lograr la congelación asimétrica. La DM FIMP ψ se produce en las desintegraciones de χ mediante interacciones débiles de Yukawa y_ϕ . En la región sombreada en gris (gris oscuro) para valores mayores de y_ϕ , observamos que la DM ψ producida mediante la congelación es simétrica (parcialmente asimétrica), como lo indica la razón de asimetría r . En la región blanca para valores menores de y_ϕ , la ψ se produce predominantemente en las desintegraciones tardías de χ y resulta ser asimétrica. La región rosa representa las restricciones de BBN sobre el tiempo de vida de χ [2].

significativamente el componente simétrico en momentos posteriores, restaurando eficazmente las aniquilaciones y mejorando las señales en la detección indirecta. Cabe destacar que, en este escenario, la abundancia de materia oscura está determinada por la asimetría inicial, a pesar de ser asintóticamente simétrica por naturaleza. Por lo tanto, tenemos masas de materia oscura de $\mathcal{O}(1 - 10)$ GeV, consistentes con la hipótesis ADM.

Además, nuestro modelo presenta una interesante señal fenomenológica en forma de un flujo monocromático de neutrinos originado en las desintegraciones tardías del componente más pesado de la materia oscura, potencialmente detectable en telescopios de neutrinos. Si bien el modelo principal se basa en una extensión del balancín de tipo I y es difícil de investigar, también proporcionamos un modelo de balancín inverso que puede funcionar a escalas bajas y comprender el mecanismo con éxito.

Sectores claros y oscuros y Warm DM

En el capítulo 2, analizamos la tensión entre los límites de masa de los neutrinos obtenidos a partir de la cosmología y los experimentos de oscilación. Además, si se asumen estos límites cosmológicos como válidos, implicaría que sería difícil detectar las masas de los neutrinos en experimentos terrestres como KATRIN.

Dado que los límites de masa cosmológicos son sensibles únicamente a la densidad energética del neutrino y no directamente a su masa, estos límites pueden relajarse si existe un mecanismo para suprimir su densidad numérica. En la referencia [1], presentamos un mecanismo para relajar el límite de masa cosmológico mediante la introducción de un sector oscuro claro con varios fermiones sin masa (radiación oscura) y un candidato a neutrino estéril de masa molecular en el rango de masas $\mathcal{O}(10 - 100)$

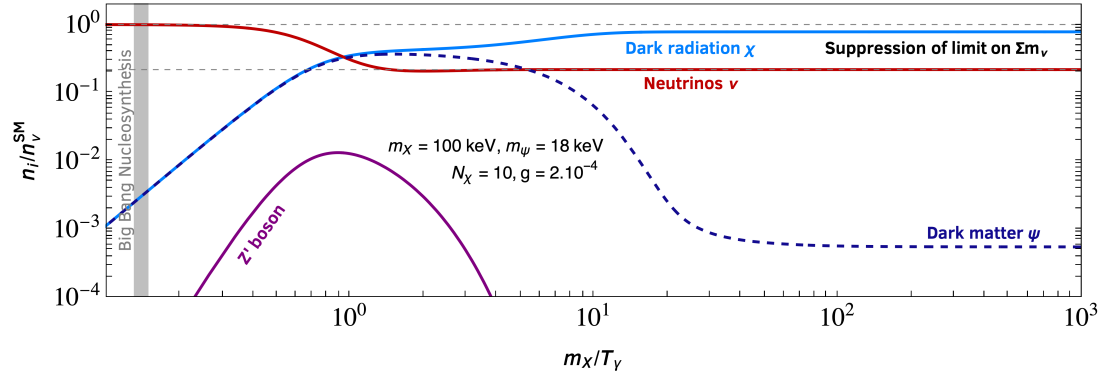


Figure 10.4: Representación esquemática del mecanismo para reducir la densidad numérica de neutrinos y relajar la masa ligada a ellos en la referencia [1]. La línea discontinua gris inferior corresponde a $n_\nu/n_\nu^{SM} \simeq 0.2$.

keV.

El sector oscuro alcanza el equilibrio térmico con los neutrinos activos tras el desacoplamiento de los neutrinos y antes de la recombinación en el rango de temperaturas $O(100 \text{ keV}) \gtrsim T \gtrsim O(10 \text{ eV})$ mediante una interacción de gauge $U(1)$ en el sector oscuro. Los estados del sector claro-oscuro, incluyendo el DM, la radiación oscura y el bosón de gauge oscuro Z' , se pueblan en la extensión de los neutrinos activos mediante estas nuevas interacciones de gauge. Cabe destacar que, dado que los neutrinos ya se han desacoplado, ya no se producen a partir de las partículas del SM. Por lo tanto, el mecanismo suprime la abundancia de neutrinos activos y, por lo tanto, concilia las masas considerables de neutrinos con la cosmología.

La representación esquemática del mecanismo obtenido mediante la resolución numérica de los EB relevantes se muestra en la Fig. 10.4. Tras el desacoplamiento, los neutrinos se equilibran con las partículas del sector oscuro: Z' , DM ψ y N_χ generaciones de fermiones sin masa χ . Las interacciones $\nu\nu \rightarrow Z'$ reducen la población de neutrinos, mientras que interacciones como $Z' \rightarrow \chi\chi$ y $Z' \rightarrow \psi\psi$ (según la jerarquía) pueblan los fermiones en el sector oscuro. Las N_χ generaciones de χ contribuyen como radiación oscura y compensan la disminución de N_{eff} debida a la reducción de la densidad numérica de neutrinos. La abundancia de reliquias DM se determina una vez que las aniquilaciones del neutrino estéril DM ψ como $\psi\psi \rightarrow \chi\chi$ se congelan en el sector oscuro (ver la distintiva curva de congelamiento).

Incorporamos el mecanismo de Ref. [353] en un modelo UV completo basado en el balancín tipo I mínimamente extendido [354]. Este modelo es similar a nuestro modelo en Ref. [2]; sin embargo, incluye un sector oscuro más rico y no presenta asimetría oscura. Exploramos la fenomenología del candidato WDM en detalle, incluyendo las restricciones derivadas de la estabilidad y la formación de la estructura.

El límite cosmológico de masa del neutrino puede relajarse por un factor de $\sim 5(8)$ para 10 (20) generaciones de fermiones oscuros sin masa, mientras que un número menor

de generaciones está restringido por N_{eff} . Una de las principales características de este modelo es un ligero aumento de N_{eff} en el CMB que en el BBN, así como la supresión del espectro de potencia de la materia a escalas menores. Además, el modelo es compatible con la leptogénesis convencional exitosa, aunque este no era el objetivo principal del trabajo.

10.3.3 Marcos unificados

Vimos que el balancín tipo I puede extenderse para incluir un candidato a DM. Estos marcos unificados se basan en la motivación de encontrar un origen común para los tres problemas clave de BSM en los que se centra esta tesis. Este origen común motivó los trabajos realizados en las referencias [8,9].²

Un argumento de simetría continua

En la referencia [9], proponemos un modelo que presenta un mediador escalar pesado con un número leptónico distinto de cero (L) que presenta interacciones que conservan y violan el número leptónico en los sectores visible y oscuro, respectivamente. El escalar se desintegra simultáneamente en los neutrinos dextrógiros pesados (cargados bajo $U(1)_L$) mediante interacciones sustanciales (lo que conduce al balancín de tipo I y la leptogénesis) y en los fermiones oscuros (sin carga bajo $U(1)_L$) mediante interacciones débiles, dando lugar así a una materia oscura similar a la FIMP. La pesadez y la pequeñez de los parámetros que caracterizan la nueva física se atribuyen naturalmente a la posterior conservación y violación de esta simetría del número leptónico, respectivamente.

Un argumento de simetría discreta

En la referencia [8], presentamos otro modelo basado en una simetría discreta $\mathbb{Z}_4$, en la que un escalar pesado conecta nuevamente el sector de neutrinos con un sector oscuro del universo con interacción débil. La simetría $\mathbb{Z}_4$ se rompe a un remanente $\mathbb{Z}_2$ cuando el escalar adquiere un VEV, generando así masas de neutrinos diminutas y explicando la asimetría bariónica observada mediante leptogénesis. Este modelo introduce nuevas contribuciones a la asimetría CP a partir de las dispersiones. La escala de ruptura de $\mathbb{Z}_4$ ocurre a escalas de energía significativamente mayores, lo que podría conducir a la formación de paredes de dominio (DW). Estos defectos podrían dominar la densidad energética del universo, resultando en una catástrofe cosmológica a menos que se vuelva inestable. En consecuencia, si la simetría discreta es solo aproximada, la colisión y aniquilación de estas ondas gravitacionales (DW) puede generar un fondo estocástico de ondas gravitacionales (OG). Dicha señal podría ser detectada por los detectores de OG actuales y futuros, lo que ofrece una ventana para explorar no solo modelos de balancín a gran escala, sino también para sondear el sector oscuro de interacción débil.

² Cabe destacar que los modelos de las referencias [1,2] también son marcos unificados, aunque se centraron principalmente en la fenomenología de la DM y, por lo tanto, se han analizado anteriormente.

10.3.4 Reflexiones

Al mirar retrospectivamente el Modelo Estándar, pasaron aproximadamente 115 años desde el primer descubrimiento del electrón en 1897 hasta el último descubrimiento del bosón de Higgs en 2012. Durante estos años, nuestra comprensión del universo se ha expandido significativamente, gracias a los extensos esfuerzos experimentales de las últimas décadas. Sin embargo, aún queda mucho por conocer. Ahora estamos entrando en la era de la física de neutrinos de precisión, con numerosos experimentos en curso o en desarrollo. En la próxima década, esperamos obtener respuestas a las preguntas abiertas en el sector neutrino. Aunque estos avances también nos ayudarán a explorar posibles conexiones entre los neutrinos y el sector oscuro, las búsquedas de materia oscura están volviéndose cada vez más sensibles. Con próximas misiones espaciales y estrategias novedosas para detectar materia oscura ligera y ultraligera, finalmente podríamos arrojar algo de luz sobre lo oscuro.

A pesar de la etapa difícil que enfrenta actualmente el campo, los marcos BSM bien motivados e interdisciplinarios ayudarán a allanar el camino hacia un nuevo Modelo Estándar. Aunque no sepamos cuánto tiempo tomará la aparición de este 'Nuevo SM', los esfuerzos continuos en construcción de modelos, las restricciones experimentales e incluso los resultados nulos seguirán acercándonos a desentrañar los secretos de nuestro universo.

Appendices

A

SEESAW EFFECTIVE LAGRANGIAN

Here we discuss the EFT framework to construct the effective Lagrangian for the Type-I seesaw as an example. Since, N_R 's are Majorana fermions, we can define $N \equiv N_R + N_R^c$ and express eq. (2.15) as

$$\mathcal{L}_N = \frac{1}{2} \overline{N}_i i \not{\partial} N_i - \frac{1}{2} M_i \overline{N}_i N_i - \left[Y_\nu \overline{L} \tilde{H} N_i + \text{H.c.} \right], \quad (\text{A.1})$$

where we have suppressed the flavour indices. The term above in the square brackets can be expanded as

$$Y_\nu \overline{L} \tilde{H} N_i + \text{H.c.} = \frac{1}{2} \left[\overline{N} (Y_\nu^T \tilde{H}^T L^c + Y_\nu^\dagger \tilde{H}^\dagger L) + (\overline{L} Y_\nu \tilde{H} + \overline{L}^c Y_\nu^* \tilde{H}^*) N \right], \quad (\text{A.2})$$

and we can write $\mathcal{L}_N$ in a compact form (omitting the N generation index i)

$$\mathcal{L}_N = \frac{1}{2} \overline{N} \mathcal{K} N - \frac{1}{2} \overline{N} \zeta - \frac{1}{2} \overline{\zeta} N, \quad (\text{A.3})$$

where $\mathcal{K} \equiv i \not{\partial} - M$ and $\zeta \equiv Y_\nu^T \tilde{H}^T L^c + Y_\nu^\dagger \tilde{H}^\dagger L$ with $\overline{\zeta} = \zeta^\dagger \gamma^0$. Now in order to obtain the effective Lagrangian, let us recall eq. (1.28) to write

$$e^{iS_{\text{eff}}} = e^{iS_{\text{SM}}} \int D N D \overline{N} e^{iS_N} = e^{iS_{\text{SM}}} \int D N D \overline{N} e^{i \int d^4 x \mathcal{L}_N}. \quad (\text{A.4})$$

The integral over $\mathcal{L}_N$ to integrate out the heavy fermion fields is typically Gaussian and we use the functional integration for fermion fields

$$\int D \psi D \overline{\psi} \exp \left[i \int \overline{\psi}(x) K(x, y) \psi(y) - i \int [\overline{J}(z) \psi(z) + \text{H.c.}] \right] \quad (\text{A.5})$$

$$= [\text{Det } K] \exp \left[-i \int d^4 x d^4 y \overline{J}(x) K^{-1}(x, y) J(y) \right], \quad (\text{A.6})$$

to express it as

$$\begin{aligned} & \int DND\bar{N} \exp \left[i \int d^4x \frac{1}{2} \left(\bar{N} \mathcal{K} N - \bar{N} \zeta - \bar{\zeta} N \right) \right] \\ &= [\text{Det } \mathcal{K}] \exp \left[-\frac{i}{2} \int \bar{\zeta} \mathcal{K}^{-1} \zeta d^4x \right], \end{aligned} \quad (\text{A.7})$$

where $[\text{Det } K]$ is a constant and field independent, hence its contribution can be ignored. Further, one can expand the operator $\mathcal{K}^{-1} = (i\not{\partial} - M)^{-1}$ in powers of M

$$\frac{1}{i\not{\partial} - M} = -\frac{1}{M} - \frac{i\not{\partial}}{M^2} + \dots \quad (\text{A.8})$$

The leading order term in this expansion gives us the lepton number violating dimension-5 Weinberg operator. Upon substituting the value of ζ in eq. (A.7), we get

$$\frac{\mathcal{L}_5}{\Lambda} = \frac{1}{2} Y_\nu \bar{L} \tilde{H} M^{-1} Y_\nu^T \tilde{H}^T L^c + \text{H.c.} = \frac{1}{2} (Y_\nu M^{-1} Y_\nu^T)_{\alpha\beta} \bar{L}_\alpha \tilde{H} \tilde{H}^T L_\beta^c + \text{H.c.}, \quad (\text{A.9})$$

recovering the Weinberg operator and the Wilson coefficient,

$$\kappa_{\alpha\beta} = (Y_\nu M^{-1} Y_\nu^T)_{\alpha\beta}. \quad (\text{A.10})$$

where we have reintroduced the flavour indices in the last step. Post EWSSB, this operator will generate Majorana masses for the light neutrinos $m_{a\beta} = \kappa_{\alpha\beta} v_H^2/2$. A similar approach can be used to derive the dimension-5 operators for the other seesaw cases. Note that one can also evaluate the tree level diagrams corresponding to s, t -channel exchange of heavy N_R , take the limit $M_R^2 \gg s, t$ and then *match* it to get the corresponding operator.

B

LOOP CONTRIBUTIONS TO CP ASYMMETRY

In this appendix, we explicitly derive the contributions from the one-loop diagrams to the CP asymmetry from RHN decays in type-I leptogenesis discussed in Sec. 3.3.1.

B.1 Vertex contribution to asymmetry

The first contribution to the CP asymmetry comes at loop level via the interference of the vertex diagram shown in and the tree level decay shown in Fig. 3.2, we show the one loop vertex correction diagram with mommentum flows in Fig. B.1. The amplitude of this interference term can be written as:

$$\begin{aligned}
 I_{\text{vertex}} &= \int \frac{d^4 q_1}{(2\pi)^4} \bar{u}_\alpha(q) [-iY_{\alpha j} P_R] \frac{i(\not{q}_3 + M_j)}{q_3^2 - M_j^2 + i\epsilon} [-iY_{\beta j} P_R] \frac{i(-\not{q}_1)}{q_1^2 + i\epsilon} [iY_{\beta i}^* P_L] u_i(p) \\
 &\quad \frac{i}{q_2^2 + i\epsilon} \times \underbrace{[\bar{u}_\alpha(q) [-iY_{\alpha i} P_R] u_i(p)]^\dagger}_{\mathcal{M}_{\text{tree}}} \\
 &= i\mathbf{Y} \int \frac{d^4 q_1}{(2\pi)^4} \frac{\bar{u}_\alpha(q) P_R (\not{q}_3 + M_j) P_R (\not{q}_1) P_L u_i(p) \bar{u}_i(p) P_L u_\alpha(q)}{(q_3^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)}, \tag{B.1}
 \end{aligned}$$

where $\mathbf{Y} = [Y_{\alpha i}^* Y_{\alpha j} Y_{\beta j} Y_{\beta i}^*]$ is the product of the involved couplings. Taking spin sum and trace of the numerator, we obtain

$$\begin{aligned}
 I_{\text{vertex}} &= i \frac{\mathbf{Y}}{2} \int \frac{d^4 q_1}{(2\pi)^4} \frac{\text{Tr}[P_R (\not{q}_3 + M_j) P_R (\not{q}_1) P_L (\not{q} + M_i) P_L (\not{q})]}{(q_3^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \\
 &= i\mathbf{Y} M_i M_j \int \frac{d^4 q_1}{(2\pi)^4} \frac{q_1 \cdot q}{(q_3^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)}. \tag{B.2}
 \end{aligned}$$

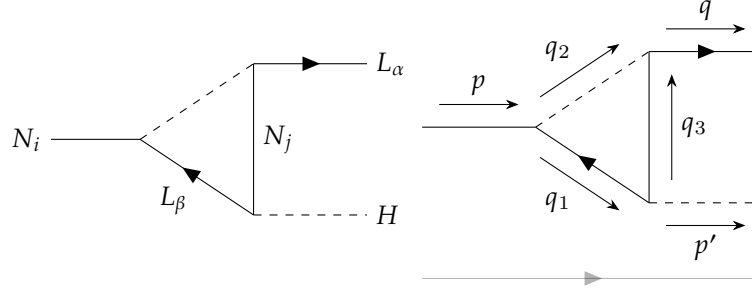


Figure B.1: One loop vertex correction for the process $N_j \rightarrow l\phi$ and the corresponding momentum flows. The grey line corresponds to the chosen orientation following the Feynman rules for Majorana fermions [366].

We can rewrite the above integral in the following form

$$I_{\text{vertex}} = -\mathbf{Y} \frac{M_i M_j}{16\pi^2} \left[\int \frac{-i d^4 q_1}{\pi^2 (2\pi)^4} \frac{q_1 \cdot q}{(q_3^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \right], \quad (\text{B.3})$$

where the integral in the square brackets can be solved using FeynCalc [363, 367] to reduce it to one loop functions (B, C) using Passarino-Veltman reduction [368]

$$I'_v = -\frac{M_i M_j}{16\pi^2} \left[\frac{1}{2} B_0(0, 0, M_j^2) + \frac{1}{2} B_0(M_i^2, 0, 0) + \frac{1}{2} M_j^2 C_0(0, 0, M_i^2, 0, M_j^2, 0) \right. \\ \left. + \frac{1}{2} M_i^2 C_0(0, 0, M_i^2, 0, M_j^2, 0) \right].$$

Since we need only the imaginary part of the interference term for evaluation of CP asymmetry, see eq. (3.11), we just need to extract the imaginary part of the B_0 and C_0 functions, which can be obtained from Package-X [364]:

$$\begin{aligned} \text{Im}[B_0(0, 0, M_m^2)] &= 0 \\ \text{Im}[B_0(M_k^2, 0, 0)] &= \pi \Theta(M_k^2) \\ \text{Im}[C_0(0, 0, M_k^2, 0, M_m^2, 0)] &= -\frac{\pi \Theta(M_k^2)}{M_k^2} \ln \left[1 + \frac{M_k^2}{M_m^2} \right]. \end{aligned} \quad (\text{B.4})$$

Using these, we get the imaginary part

$$\text{Im}(I'_v) = -\frac{M_i M_j}{32\pi^2} \left[\pi \Theta(M_i^2) - (M_j^2 + M_i^2) \frac{\pi \Theta(M_i^2)}{M_i^2} \ln \left(1 + \frac{M_i^2}{M_j^2} \right) \right]. \quad (\text{B.5})$$

Defining $z \equiv M_j^2/M_i^2$ and getting rid of the Θ functions, we get

$$\text{Im}(I'_v) = -\frac{M_i^2}{32\pi} \sqrt{z} \left(1 - (z + 1) \ln \left[\frac{1 + z}{z} \right] \right). \quad (\text{B.6})$$

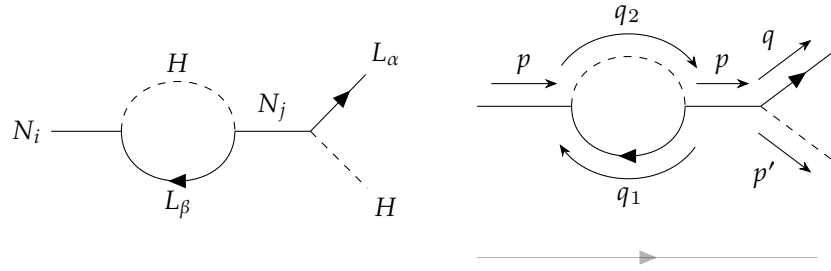


Figure B.2: Self-energy for the process $N_j \rightarrow l\phi$ and the corresponding momentum flows. The grey line corresponds to the chiral orientation.

We can now derive the vertex contribution to the CP asymmetry by using eq. (3.11)

$$\varepsilon_v = -\frac{4}{\Gamma_{\text{total}}} \sum_{i \neq j} \sum_{\beta} \text{Im}(\mathbf{Y}) \text{Im}(I'_v V_\varphi) \quad (\text{B.7})$$

where $\Gamma_{\text{total}} = \frac{(Y^\dagger Y)_{ii}}{8\pi} M_i$, and V_φ is the two body phase space factor corresponding to the two decay channels

$$V_\varphi = 2 \times \frac{|\vec{q}|}{8\pi E_{\text{cm}}^2} = \frac{1}{8\pi M_i}. \quad (\text{B.8})$$

Putting it all together, we have

$$\begin{aligned} \varepsilon_v &= 4 \times \frac{8\pi}{(Y^\dagger Y)_{ii} M_i} \sum_{i \neq j} \sum_{\beta} \text{Im}[Y_{\alpha i}^* Y_{\alpha j} Y_{\beta j} Y_{\beta i}^*] \times \frac{M_i^2}{32\pi} \sqrt{z} \left(1 - (z+1) \ln \left[\frac{1+z}{z} \right] \right) \frac{1}{8\pi M_i}, \\ &= \frac{1}{8\pi} \sum_{m \neq j} \frac{\text{Im}[Y_{\alpha i}^* Y_{\alpha j} (Y^\dagger Y)_{ij}]}{(Y^\dagger Y)_{ii}} f_v(z), \end{aligned} \quad (\text{B.9})$$

where we have defined the vertex contribution function f_v as

$$f_v(z) = \sqrt{z} \left(1 - (z+1) \ln \left[\frac{1+z}{z} \right] \right). \quad (\text{B.10})$$

B.2 Self energy contribution to asymmetry

It was pointed in [369] that the self-energy diagrams can also significantly contribute to the asymmetry, in addition to the vertex correction diagrams. Therefore, we consider the interference of the self energy diagram and the tree level decay shown in Fig. 3.2, the momentum flows are shown in Fig. B.2.

$$\begin{aligned}
I_{\text{self}} &= \int \frac{d^4 q_1}{(2\pi)^4} \bar{u}_\alpha(q) [-iY_{\alpha j} P_R] \frac{i(\not{p} + M_j)}{p^2 - M_j^2 + i\epsilon} [-iY_{\beta j} P_R] \frac{i(-\not{q}_1)}{q_1^2 + i\epsilon} [iY_{\beta i}^* P_L] u_i(p) \\
&\quad \frac{i}{q_2^2 + i\epsilon} \times \underbrace{[\bar{u}_\alpha(q) [-iY_{\alpha i} P_R] u_i(p)]^\dagger}_{\mathcal{M}_{\text{tree}}} \\
&= i\mathbf{Y} \int \frac{d^4 q_1}{(2\pi)^4} \frac{\bar{u}_\alpha(q) P_R (\not{p} + M_j) P_R (q_1) P_L u_i(p) \bar{u}_i(p) P_L u_\alpha(q)}{(p^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \\
&= i\frac{\mathbf{Y}}{2} \int \frac{d^4 q_1}{(2\pi)^4} \frac{\text{Tr}[P_R (\not{p} + M_j) P_R (q_1) P_L (\not{p} + M_i) P_L (q)]}{(p^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \\
&= i\mathbf{Y} M_i M_j \int \frac{d^4 q_1}{(2\pi)^4} \frac{q_1 \cdot q}{(p^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \tag{B.11}
\end{aligned}$$

where once again we have the same product of couplings: $\mathbf{Y} = [Y_{\alpha i}^* Y_{\alpha j} Y_{\beta j} Y_{\beta i}^*]$. Following the same approach, we rewrite the above integral and use FeynCalc to express it in terms of one loop functions.

$$I_{\text{self}} = -\mathbf{Y} \frac{M_i M_j}{16\pi^2} \left[\int \frac{-i d^4 q_1}{\pi^2 (2\pi)^4} \frac{q_1 \cdot q}{(p^2 - M_j^2 + i\epsilon)(q_2^2 + i\epsilon)(q_1^2 + i\epsilon)} \right] \tag{B.12}$$

$$I'_s = -\frac{M_i M_j}{16\pi^2} \left[\frac{M_i^2 B_0(M_i^2, 0, 0)}{4(M_i^2 - M_j^2)} \right]. \tag{B.13}$$

Considering only the imaginary part of B_0 function, we get:

$$\text{Im} I'_s = -\frac{M_i M_j}{64\pi^2} \left[\frac{M_i^2}{M_i^2 - M_j^2} \pi \Theta(M_i^2) \right] = -\frac{M_i^2}{64\pi} \sqrt{z} \left(\frac{1}{1-z} \right). \tag{B.14}$$

Thus, the self-energy contribution to the CP asymmetry can be written as

$$\varepsilon_s = -\frac{4}{\Gamma_{\text{total}}} \sum_{i \neq j} \sum_{\beta} \text{Im}(\mathbf{Y}) \text{Im}(I'_s V_\varphi). \tag{B.15}$$

The phase space in this case will be twice the one in eq. (B.8) as the loop can contain either $e^- \phi^+$ or $\nu \phi^0$, so $V_\varphi = 1/(4\pi M_j)$. Hence, the total contribution from the self-energy diagram is

$$\varepsilon_s = 4 \times \frac{8\pi}{(Y^\dagger Y)_{ii} M_i} \sum_{i \neq j} \sum_{\beta} \text{Im}[Y_{\alpha i}^* Y_{\alpha j} Y_{\beta j} Y_{\beta i}^*] \times \frac{M_i^2}{64\pi} \sqrt{z} \left(\frac{1}{1-z} \right) \frac{1}{4\pi M_i} \tag{B.16}$$

$$\varepsilon_s = \frac{1}{8\pi} \sum_{m \neq j} \frac{\text{Im}[(Y_\nu^*)_{aj} (Y_\nu)_{am} (Y_\nu^\dagger Y_\nu)_{jm}]}{(Y_\nu^\dagger Y_\nu)_{jj}} f_s(z), \tag{B.17}$$

where $f_s = \sqrt{z}[1/(1-z)]$.

BOLTZMANN EQUATIONS

In this appendix, we will first discuss the formalism for Boltzmann equations (BEs), and then derive the BE for thermal freeze-out.

C.1 Formalism

The Boltzmann equation can be written as

$$\hat{\mathbf{L}}[f] = \mathbf{C}[f], \quad (\text{C.1})$$

where $\mathbf{C}$ is the collision operator, $\hat{\mathbf{L}}$ is the Liouville operator, and f is the phase space density of a particle specie. For the FRW model of cosmology (which describes an isotropic and homogeneous universe), this Liouville operator becomes

$$\hat{\mathbf{L}}[f(E, t)] = E \frac{\partial f}{\partial t} - \frac{\dot{R}}{R} |\vec{p}|^2 \frac{\partial f}{\partial E}. \quad (\text{C.2})$$

We define the number density of a particle a as $n_a \equiv N_a/V$, where N_a is the total number of a particles in a volume $V = V_0 R^3(t)/R^3(t_0)$. The comoving volume is denoted by V_0 . In the absence of any interaction, N_a is conserved and $\dot{N}_a = 0$, implying

$$\frac{d(n_a V)}{dt} = V(\dot{n}_a + 3H n_a) = 0, \quad (\text{C.3})$$

where we use $H \equiv \dot{R}/R$. In the radiation-dominated epoch, $H = 1.66 g_*^{1/2} T^2/M_{\text{Pl}}$, with $M_{\text{Pl}} = 1.2 \times 10^{19} \text{ GeV}$, and g_* as the effective number of degrees of freedom given by

$$g_* \equiv \sum_b g_b + \frac{7}{8} \sum_f g_f. \quad (\text{C.4})$$

In the early universe, when the temperature was high enough and all particles were in thermal equilibrium, we have $g_* = 28 + 7/8 \times 90 = 106.75$.

However, in the presence of interactions, the a particles will undergo collisions and their number density will change subsequently. For example, an interaction: $a + B \rightarrow C$ decreases the number by one unit, whereas the inverse of the interaction increases it by one unit. This change of particle number in a physical volume is governed by the collision operator.

The number density for a particle in thermal equilibrium at a temperature T is given by

$$n_a = g_a \int \frac{d^3 p}{(2\pi)^3} f(E, t) = g_a \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\exp[(E - \mu)/T] \pm 1}, \quad (\text{C.5})$$

where g_a is the number of internal degrees of freedom of a . Upon integrating eq. (C.2) and multiplying by $g_a/(2\pi)^3 E$ both sides, we get:

$$\frac{1}{E} \frac{g_a}{(2\pi)^3} \int d^3 p \hat{\mathbf{L}}[f(E, t)] = \frac{g_a}{(2\pi)^3} \int d^3 p \left[\frac{\partial f}{\partial t} - H \frac{|\vec{p}|^2}{E} \frac{\partial f}{\partial E} \right]. \quad (\text{C.6})$$

Using eqs. (C.1) and (C.6), we can thus write

$$\frac{g}{(2\pi)^3} \int \mathbf{C}[f] \frac{d^3 p}{E} = \frac{dn}{dt} + 3Hn. \quad (\text{C.7})$$

The collision operator for the process $a + B \rightarrow C$ is given by:

$$\begin{aligned} \frac{g_a}{(2\pi)^3} \int \mathbf{C}[f] \frac{d^3 p_a}{E_a} = & - \sum_{B,C} \int d\Pi_a d\Pi_B d\Pi_C (2\pi)^4 \delta^4(p_a + p_B - p_C) \\ & \times [f_a f_B (1 \pm f_C) |\mathcal{M}|_{a+B \rightarrow C}^2 - f_C (1 \pm f_a)(1 \pm f_B) |\mathcal{M}|_{C \rightarrow a+B}^2], \end{aligned} \quad (\text{C.8})$$

where the $+$ ($-$) applied to bosons (fermions), and f_x is the phase density of a particle specie x ; and

$$d\Pi_x \equiv g_x \frac{1}{(2\pi)^3} \frac{d^3 p}{2E}. \quad (\text{C.9})$$

Here, the δ -function ensures energy momentum conservation.

We further need to make the following assumptions to simplify eq. (C.8):

1. We use Maxwell Boltzmann statistics for all particle species, so that we have $f_x \propto e^{(E-\mu)/T}$.
2. In the absence of Bose condensation and Fermi degeneracy, we can also ignore the blocking and stimulated emission factors, so that $1 \pm f \simeq 1$ since $f_x \approx e^{-\langle E \rangle/T} \approx 0.05$, with $\langle E \rangle = 3T$ as the average energy.

The thermally averaged reaction density is defined as

$$\gamma(a + B \rightarrow C) \equiv \int d\Pi_a d\Pi_B d\Pi_C (2\pi)^4 \delta^4(p_a + p_B - p_C) \times e^{-E_a/T} e^{-E_B/T} |\mathcal{M}(a + B \rightarrow C)|^2. \quad (\text{C.10})$$

Therefore, under the aforementioned assumptions and using eq. (C.10) in eq. (C.8), from eq. (C.7), we get the generic form of BE for a particle a as

$$\dot{n}_a + 3Hn_a = \sum_{B,C} \left[\frac{n_C}{n_C^{\text{eq}}} \gamma(C \rightarrow a + B) - \frac{n_a}{n_a^{\text{eq}}} \cdot \frac{n_B}{n_B^{\text{eq}}} \gamma(a + B \rightarrow C) \right]. \quad (\text{C.11})$$

It is much more convenient to consider the evolution of the number of particles in a co-moving volume to scale out the effect of expansion of the universe. This is usually done by defining dependent variables $Y_a \equiv n_a/s$ and $x = m_a/T$ where s is the entropy density of the universe

$$s \equiv \left(\frac{2\pi^2}{45} \right) g_{*s} \frac{m_a^3}{x^3}, \quad (\text{C.12})$$

with the effective degrees of freedom associated with entropy, $g_{*,s}$, given by

$$g_{*,s}(T) \equiv \sum_b g_b \left(\frac{T_{b,\text{dec}}}{T} \right)^3 + \frac{7}{8} \sum_f g_f \left(\frac{T_{f,\text{dec}}}{T} \right)^3. \quad (\text{C.13})$$

Here, the T_{dec} denotes the decoupling temperature of the corresponding boson/fermion.

In terms of yield and entropy, we have

$$\frac{dY_a}{dx} = \frac{dt}{dx} \frac{d}{dt} \left(\frac{V n_a}{V s} \right) = \frac{dt}{dx} \frac{\dot{n}_a + 3Hn_a}{s}, \quad (\text{C.14})$$

where the total entropy Vs is conserved. In the radiation dominated epoch where $T \propto 1/\sqrt{t}$, one can write $dt/dx = 1/(Hx)$ and the BE of eq. (C.11) can be rewritten as

$$\frac{dY_a}{dx} = \frac{1}{sHx} \sum_{B,C} \left[\frac{Y_C}{Y_C^{\text{eq}}} \gamma(C \rightarrow a + B) - \frac{Y_a}{Y_a^{\text{eq}}} \cdot \frac{Y_B}{Y_B^{\text{eq}}} \gamma(a + B \rightarrow C) \right], \quad (\text{C.15})$$

where the equilibrium yield of a massive specie i is given by

$$Y_i^{\text{eq}} = \frac{45g_i}{4\pi^4 g^*} x^2 K_2(x). \quad (\text{C.16})$$

C.2 BE for freeze-out

The generic form above applies in a straightforward way to the case of decays. Here, we discuss the formalism $2 \leftrightarrow 2$ interactions such as $\chi\bar{\chi} \leftrightarrow X\bar{X}$ relevant for Section 4.3.

Using eqs. (C.7) and (C.8), we have

$$\begin{aligned} \frac{dn_\chi}{dt} + 3Hn_\chi = & - \int d\Pi_\chi d\Pi_{\bar{\chi}} d\Pi_X d\Pi_{\bar{X}} (2\pi)^4 |\mathcal{M}|^2 \\ & \times \delta^4(p_X + p_{\bar{X}} - p_\chi - p_{\bar{\chi}}) [f_\chi f_{\bar{\chi}} - f_X f_{\bar{X}}], \end{aligned} \quad (\text{C.17})$$

where we use $|\mathcal{M}|_{\chi\bar{\chi} \rightarrow X\bar{X}}^2 = |\mathcal{M}|_{X\bar{X} \rightarrow \chi\bar{\chi}}^2 = |\mathcal{M}|^2$ from CPT invariance.

We will assume that all the particles have thermal distributions with $\mu = 0$. Further, we assume that the particles $X, \bar{X}$ are in thermal equilibrium, so that we have

$$f_X = \exp[-E_X/T], \quad f_{\bar{X}} = \exp[-E_{\bar{X}}/T]. \quad (\text{C.18})$$

The δ -function enforces energy conservation: $E_\chi + E_{\bar{\chi}} = E_X + E_{\bar{X}}$, leading to the following relation

$$f_X f_{\bar{X}} = \exp[-(E_X + E_{\bar{X}})/T] = \exp[-(E_\chi + E_{\bar{\chi}})/T] = f_\chi^{\text{eq}} f_{\bar{\chi}}^{\text{eq}}, \quad (\text{C.19})$$

where $f_\chi^{\text{eq}} \equiv \exp[-E_\chi/T]$ and $f_{\bar{\chi}}^{\text{eq}} \equiv \exp[-E_{\bar{\chi}}/T]$. Under the assumption of no initial asymmetry so that $n_i = n_{\bar{i}}$ for all species, we can use the above relation to rewrite eq. (C.18) in terms of number densities

$$\frac{d_{ch}i}{dt} + 3Hn_\chi = -\langle\sigma v\rangle \left(n_\chi^2 - n_\chi^{\text{eq},2} \right),$$

where the thermally averaged annihilation cross-section is given by

$$\begin{aligned} \langle\sigma v\rangle = & (n_\chi^{\text{eq}})^{-2} \int d\Pi_\chi d\Pi_{\bar{\chi}} d\Pi_X d\Pi_{\bar{X}} (2\pi)^4 \\ & \times \delta^4(p_X + p_{\bar{X}} - p_\chi - p_{\bar{\chi}}) |\mathcal{M}|^2 \exp(-E_\chi/T) \exp(-E_{\bar{\chi}}/T). \end{aligned} \quad (\text{C.20})$$

Finally, we can rewrite the BE in terms of yield as

$$\frac{dY_\chi}{dx} = -\frac{s}{Hx} \langle\sigma v\rangle \left(Y_\chi^2 - Y_\chi^{\text{eq},2} \right), \quad (\text{C.21})$$

where $x = m_\chi/T$.

ABBREVIATIONS

- Λ CDM** Λ Cold Dark Matter (*p.* 12)
- ν MSM** Neutrino Minimal Standard Model (*p.* 87)
- $0\nu\beta\beta$** Neutrinoless Double Beta Decay ($0\nu\beta\beta$) (*p.* 25)
- 2HDM** Two Higgs Doublet Model (*p.* 46)
- ADM** Asymmetric Dark Matter (*p.* 81)
- ALPs** Axion-like Particles (*p.* 87)
- BAO** Baryon Acoustic Oscillations (*p.* 44)
- BAU** Baryon Asymmetry of the Universe (*p.* 41, 81)
- BBN** Big Bang Nucleosynthesis (*p.* 43)
- BEs** Boltzmann Equations (*p.* 52)
- BSM** Beyond the Standard Model (*p.* 1)
- CEvNS** Coherent Elastic Neutrino-Nucleus Scattering (*p.* 75)
- CHAMPs** Charged Massive Particles (*p.* 66)
- CKM** Cabibbo-Kobayashi-Mashikawa (*p.* 11)
- CMB** Cosmic Microwave Background (*p.* 44)
- DD** Direct Detection (*p.* 73)
- DM** Dark Matter (*p.* 13, 63)
- DW** Dodelson–Widrow (*p.* 86)
- EFT** Effective Field Theory (*p.* 16)
- EW** Electroweak (*p.* 5)
- EWPO** Electroweak Precision Observables (*p.* 12)
- EWPT** Electroweak Phase Transition (*p.* 46)
- FIMP** Feebly Interacting Massive Particle (*p.* 77)
- FOPT** First-Order Phase Transition (*p.* 46)
- GScM** Generalised Scotogenic Model (*p.* 38)
- GTR** General Theory of Relativity (*p.* 12)
- GUT** Grand Unified Theory (*p.* 15)
- HNLs** Heavy Neutral Leptons (*p.* 27)

- ID** Indirect Detection (*p.* 73)
IO Inverted Ordering (*p.* 24)
ISS Inverse Seesaw (*p.* 35)
- LFV** Lepton Flavour Violation (*p.* 11)
LHC Large Hadron Collider (*p.* 40)
LN Lepton Number (*p.* 12)
LNV Lepton Number Violation (*p.* 40)
LSS Linear Seesaw (*p.* 36)
- MACHOs** Massive Compact Halo Objects (*p.* 67)
MOND Modified Newtonian Dynamics (*p.* 63)
MSSM Minimal Supersymmetric Standard Model (*p.* 46)
- NO** Normal Ordering (*p.* 24)
NP New Physics (*p.* 16)
- PBH** Primordial Black Hole (*p.* 67)
PMNS Pontecorvo-Maki-Nagakawa-Sakata (*p.* 11)
- QCD** Quantum Chromodynamics (*p.* 4)
QED Quantum Electrodynamics (*p.* 3)
QFT Quantum Field Theory (*p.* 2)
- RHNs** Right-Handed Neutrinos (*p.* 30)
- ScM** Scotogenic Model (*p.* 37)
SIDM Self-Interacting Dark Matter (*p.* 68)
SM Standard Model (*p.* 1)
SSB Spontaneous Symmetry Breaking (*p.* 6)
SSI Type-I seesaw (*p.* 30)
SSII Type-II seesaw (*p.* 32)
SSIII Type-III seesaw (*p.* 33)
- VEV** Vacuum Expectation Value (*p.* 7)
- WDM** Warm Dark Matter (*p.* 68)
WIMPs Weakly Interacting Massive Particles (*p.* 71)

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