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New Physics Searches from Cosmic-Ray Air-Showers

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Declaran:

Que la presente memoria, New Physics Searches from Cosmic-Ray Air-Showers, ha sido realizada por Víctor Muñoz Albornoz bajo su dirección, y constituye su Tesis para optar al grado de Doctor en Física por la Universidad de València.

Y para que así conste, en cumplimiento de la legislación vigente, presenta en el Departamento de Física Teórica de la Universidad de València la referida Tesis Doctoral, y firman el presente certificado.

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Abstract

Departament de Física Teòrica

Doctorat en Física

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by Víctor Muñoz Albornoz

This thesis deals with the study of possible production and detection of new fundamental particles, created in atmospheric showers from the collision of cosmic-rays with the upper layers of our atmosphere.

We will distinguish two broad categories for the particles that could be searched in terrestrial experiments; long-lived particles and stable particles. In the first case, we consider particles with a large enough lifetime that could propagate through the atmosphere, and penetrate underground to decay inside the instrumented volume of large neutrino detectors. On the other hand, we consider stable particles that could scatter on the target material of underground experiments, leaving a detectable signal. The models considered in this thesis are motivated by both, theoretical puzzles, and experimental evidence of physics beyond the Standard Model of particle physics.

In the Chapter 1 of this thesis we provide a general overview of the Standard Model, as well as the motivation for the extensions that have been studied in this work. In Chapter 2, we provide an introduction to cosmic-rays and atmospheric showers. In Chapter 3 we discuss the atmospheric production and detection of new fundamental particles belonging to the models introduced on the first chapter, providing the main original results of this thesis.

A compendium in Spanish including the motivation, and results of this thesis is given in Chapter 4. Finally, the last piece of this work can be found in Chapter 5, which contains the original peer-reviewed publications that constitutes the main body of this dissertation.

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Although this thesis has been written only by me, it is the result of a collective effort and a long road that began several years ago.

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Chapter 1

The Standard Model of Particle Physics and Beyond

For over two thousand years, humans have been thinking about what matter is made of, giving birth to the idea that objects in our universe are composed of basic constituents or fundamental particles. Despite this long history, it was the development of quantum theory over the past century what really fueled the substantial progress in answering this question, by providing an accurate description of the elementary particles that compose matter, as well as their basic interactions. The current state of the classification of all elementary particles, and the description of three of the four fundamental interactions is contained in the Standard Model (SM) of particle physics [106, 187, 177, 82, 130, 129, 119, 140, 107, 94, 123]. In what follows we give a brief overview of its main ingredients and consequences.

1.1 Particles and Symmetries

All elementary particles described by the SM can be classified according to their spin, and there are two broad classes; particles with half-integer spin (fermions), and particles with integer spin (bosons). This spin classification of the elementary particles is connected with the fact that they fall into irreducible representations of the Poincaré group, which is defined by the set of transformations implementing the Minkowski space-time symmetry.

The SM contains twelve spin-1/2 fermions, plus their corresponding antiparticles. These fermions can be classified according to how they interact into six quarks including the up (u), down (d), charm (c), strange (s), top (t) and bottom quark (b), each coming in three colours, and six leptons including the electron (e), electron neutrino (ν_e), muon (μ), muon neutrino (ν_μ), tau (τ), and the tau neutrino (ν_τ).

On the other hand, the SM contains a total of thirteen bosons; twelve of them are spin-1 vector bosons, which includes eight gluons (G) acting as the force carriers of the strong interaction, three electroweak gauge bosons $W^\pm$ and Z , the photon γ as mediator of the electromagnetic interaction, and finally, a single spin-0 scalar H , the so-called Higgs boson, which was discovered at the LHC in 2012 [63, 3]. The Higgs boson plays a fundamental role in the generation of masses as we will describe in the next section.

The SM is a Quantum Field Theory. In this framework, each particle has an associated quantum field, $\phi(x)$, responsible for creating and annihilating particles at a given point in space-time x . The dynamics of the fields (or the interaction of the

associated particles) is contained in the Lagrangian density $\mathcal{L}[\phi(x), \partial_\mu \phi(x)]$, which is a functional of the field and its derivatives. The action, S , is the space-time integral of the Lagrangian density

$$S[\phi] = \int d^4x \mathcal{L}. \quad (1.1)$$

Just as in classical mechanics, requiring that the variation of the action vanishes, leads to the equations of motion of the system, while conserved quantities are related to the symmetries of the Lagrangian. Note that in this thesis we will work in natural units with $\hbar = c = 1$, so the Lagrangian density will have mass dimension four.

The possible interactions among particles in the SM are determined by the symmetries of the theory, in particular the transformations associated to the rephasing of the fields

$$\phi(x) \rightarrow e^{i\alpha} \phi(x), \quad (1.2)$$

are a very important class of symmetries for the SM. More generally, we can demand invariance under a group containing multiple generators

$$\phi \rightarrow e^{i\alpha_a T_a} \phi, \quad (1.3)$$

where T_a are the generators of the symmetry group, with $a = 1 \dots n$, where n is the number of generators of the group. The symmetry transformation in equation 1.2 can be either global, if α is constant, or local when $\alpha = \alpha(x)$. Local symmetries are also referred to as gauge symmetries. The relevant symmetry groups for the SM are the Unitary group $U(N)$ formed by $N \times N$ unitary matrices, and the Special Unitary group $SU(N)$ formed by $N \times N$ unitary matrices with unit determinant.

In particular, the SM gauge symmetry group is the direct product of three groups

$$G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y, \quad (1.4)$$

where $SU(3)_c$ describes the strong force, $SU(2)_L$ describes the weak force, and $U(1)_Y$ describes an electromagnetic-like force called hypercharge. The actual electromagnetic (EM) force corresponds to the subgroup of $SU(2)_L \times U(1)_Y$ that remains unbroken after electroweak symmetry breaking (EWSB), as we will see in the next subsection.

Local symmetry transformations for the fields in the SM Lagrangian imply the existence of massless gauge bosons, one for each symmetry generator. This is easily seen by requiring that the Lagrangian describing free fields, that includes space-time derivatives, is invariant under an infinitesimal gauge transformations of equation 1.3. A gauge field A_μ^a can be introduced transforming as

$$A_\mu \rightarrow G A_\mu G^{-1} + \frac{i}{g} (\partial_\mu G) G^{-1}, \quad (1.5)$$

where $A_\mu \equiv A_\mu^a T^a$ and $G \equiv e^{i\alpha^a T^a}$. The non-covariant transformation of the derivatives of the field can then be compensated by the transformation of the gauge field in such a way that the derivatives

$$D_\mu \phi \equiv \partial_\mu \phi + ig A_\mu \phi, \quad (1.6)$$

transform covariantly

$$D_\mu \phi \rightarrow e^{i\alpha_a T_a} D_\mu \phi. \quad (1.7)$$

It is then easy to build a Lagrangian that preserves the gauge symmetry, by substituting normal derivatives by covariant ones. This also fixes how gauge fields interact with matter fields. g is the coupling constant defining the strength of the interaction.

The SM contains all the Lagrangian terms that are not forbidden by the symmetry group of equation 1.4, and it is also restricted to contain only terms up to mass dimension four. The last property makes the SM a renormalizable theory [135, 144], ensuring that the predictions of the model at low energies are not affected by any possible new dynamics that might exist at very high energies, or equivalently that the physics at long distances does not depend on the dynamics that takes place at very short distances.

A summary of the charges of the SM fields under the symmetry group of 1.4 is provided in table 1.1. The index i runs from 1 to 3 and is used to denote the family index, while the a index is used to refer to the different gauge field components (2 components for the W and 8 for the gluon field). With this charge assignment we can simply use the prescription from equation 1.6 to write explicit expressions for each of the fields, for instance the covariant derivative for the q_L field is given by

$$D_\mu q_L = \partial_\mu q_L + i g_c \lambda^a G_\mu^a q_L + i g \tau^a W_\mu^a q_L + i \frac{g'}{6} B_\mu q_L, \quad (1.8)$$

where g_c , g and g' are the gauge coupling constant of $SU(3)_c$, $SU(2)_L$ and $U(1)_Y$ respectively. An analogous procedure applies to the rest of the fields of the SM, following the quantum number assignment given in table 1.1.

Finally, note that the subscripts L and R are used to denote the left and right chiral components of fermionic fields, respectively. The decomposition of fermion fields into these chiral components is made because the $SU(2)_L$ gauge group only acts non-trivially on left-handed fields, and hence the L subscript. The chiral decomposition treats these components in a different way, and the decomposition takes place through projector operators

$$\psi = \psi_L + \psi_R = (P_L + P_R)\psi \equiv \left(\frac{1 - \gamma_5}{2} + \frac{1 + \gamma_5}{2} \right) \psi, \quad (1.9)$$

where γ_5 is the fifth gamma matrix. As can be seen from table 1.1, the left-handed fields transform as doublets of weak isospin, $SU(2)_L$, while the right-handed fields are singlets. Note that since parity symmetry relates opposite helicity eigenstates, that are in one-to-one correspondence with chiral eigenstates in the massless limit, the SM is a chiral theory that breaks parity in a maximal way.

1.2 The Standard Model Lagrangian

Now that we have described the basic ingredients of the SM, we can proceed to write down the full Lagrangian that can be written as

$$\mathcal{L}_{\text{SM}} = \mathcal{L}_{\text{Kinetic}} + \mathcal{L}_{\text{Yukawa}} + \mathcal{L}_{\text{Higgs}}. \quad (1.10)$$

$\mathcal{L}_{\text{Kinetic}}$, which describes the free propagation of the fields and also includes the gauge interactions through the covariant derivatives, is given by

$$\begin{aligned} \mathcal{L}_{\text{Kinetic}} = & (D^\mu H)^\dagger (D_\mu H) + i \bar{\ell}_L \not{D} \ell_L + i \bar{e}_R \not{D} e_R + i \bar{q}_L \not{D} q_L + i \bar{u}_R \not{D} u_R + i \bar{d}_R \not{D} d_R, \\ & - \frac{1}{4} G^{\mu\nu a} G_{\mu\nu}^a - \frac{1}{4} W^{\mu\nu a} W_{\mu\nu}^a - \frac{1}{4} B^{\mu\nu} B_{\mu\nu}, \end{aligned} \quad (1.11)$$

Type	Notation	$SU(3)_C, SU(2)_L, U(1)_Y$
Quarks ($\times 3$ families)	$q_{iL} = (u_{iL} \ d_{iL})^T$	(3, 2, 1/6)
	u_R	(3, 1, 2/3)
	d_R	(3, 1, -1/3)
Leptons ($\times 3$ families)	$\ell_{iL} = (\nu_{iL} \ e_{iL})^T$	(1, 2, -1/2)
	e_R	(1, 1, -1)
Higgs	$H = (H_1, H_2)^T$	(1, 2, 1/2)
Gauge bosons	B_μ	(1, 1, 0)
	W_μ^a	(1, 3, 0)
	G_μ^a	(8, 1, 0)

TABLE 1.1: Fields of the SM with their gauge transformation properties. The bosons of the SM are a spin-0 field (the Higgs) and eleven spin-1 fields, while the fermions are all spin-1/2 fields represented by chiral spinors. The subscripts L and R refer to left- and right-handed fields. Bold symbols are used for non-abelian groups, and indicate the dimension of the representation under which each field transforms.

where $\mathcal{D} = \gamma^\mu D_\mu$, with D_μ the appropriate covariant derivative according to the gauge transformation properties summarized in table 1.1, and a is the gauge field component index that runs over the adjoint representation. The second line of equation 1.11 involves the kinetic term for the gauge bosons, which are given by the Lorentz-invariant combination of abelian and non-abelian strength tensors with

$$G_{\mu\nu}^a = \partial_\mu G_\nu^a - \partial_\nu G_\mu^a + g_c f^{abc} G_\mu^b G_\nu^c \quad (1.12)$$

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g t^{abc} W_\mu^b W_\nu^c \quad (1.13)$$

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \quad (1.14)$$

where f^{abc} and t^{abc} are the structure constants for $SU(3)_c$ and $SU(2)_L$ respectively. Finally, note that we have dropped the family index in the Lagrangian from equation 1.11, as adding the other generations only involves adding two more copies of the kinetic terms with proper names for the fields. Kinetic terms are family blind.

The second term, $\mathcal{L}_{\text{Yukawa}}$, contains Yukawa-like interactions involving fermions and the Higgs doublet:

$$\mathcal{L}_{\text{Yukawa}} = -Y_{ij}^e \bar{\ell}_{Li} H e_{Rj} - Y_{ij}^d \bar{q}_{Li} H d_{Rj} - Y_{ij}^u \bar{q}_{Li} \tilde{H} u_{Rj} + h.c., \quad (1.15)$$

where Y_{ij}^e , Y_{ij}^d and Y_{ij}^u are 3×3 complex Yukawa matrices in flavor space with explicit family indices, and $\tilde{H} \equiv i\sigma_2 H^*$, and σ_2 is the second Pauli matrix. These interactions are responsible for generating fermion masses in the Standard Model after the breaking of the electroweak symmetry takes place. The masses of the fermions will be obtained after a diagonalization procedure which will be explained below.

Finally, the Higgs Lagrangian contains the potential responsible for breaking spontaneously the electroweak symmetry down to electromagnetism. The Higgs potential is the most general renormalizable potential that can be written

$$\mathcal{L}_{\text{Higgs}} = -\mu^2 (H^\dagger H) - \lambda (H^\dagger H)^2, \quad (1.16)$$

with λ a positive constant of mass dimension $[\text{mass}]^0$ that ensures the potential is bounded from below. Regarding the sign of μ^2 , this will determine whether the

potential has a minimum at $H^\dagger H = 0$, or at a different position. The choice $\mu^2 < 0$ implies that the minimum is for field values which satisfy

$$H^\dagger H = -\frac{\mu^2}{2\lambda}. \quad (1.17)$$

Gauge transformations give us the freedom to choose a particular minimum such that the vacuum expectation value (VEV) of the Higgs field is

$$\langle H \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}, \quad (1.18)$$

where $v = \sqrt{-\mu^2/\lambda}$, has been measured to be equal to 246 GeV [163]. While the potential is invariant under the electroweak symmetry by construction, it is clear that the vacuum solution is not. This phenomenon is called spontaneous symmetry breaking. Since the Higgs field carries no color, only the electroweak symmetry is involved and a $U(1)$ subgroup remains unbroken, this is electromagnetism:

$$G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y \rightarrow SU(3)_c \times U(1)_{em}. \quad (1.19)$$

The Goldstone theorem [109] tell us that there will be one Goldstone boson for each generator of the symmetry that is broken (three in this case). We can fix the unitary gauge to get rid of these Goldstone bosons and expand the Higgs field around the true vacuum

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix}, \quad (1.20)$$

where h is the physical Higgs. The first term in equation 1.11 will yield to the term responsible for the mass of the gauge bosons

$$\mathcal{L}_{\text{Kinetic}} \supset g^2 \frac{v^2}{8} \left[(W_\mu^1)^2 + (W_\mu^2)^2 + \left(B_\mu (g'/g) - W_\mu^3 \right)^2 \right], \quad (1.21)$$

where the mass eigenstates are defined according to the combination of gauge fields by

$$\begin{aligned} W_\mu^\pm &= \frac{1}{\sqrt{2}} (W_\mu^1 \mp W_\mu^2), \\ Z_\mu &= \cos \theta_w W_\mu^3 - \sin \theta_w B_\mu, \\ A_\mu &= \sin \theta_w W_\mu^3 + \cos \theta_w B_\mu, \end{aligned} \quad (1.22)$$

where $\tan \theta_w \equiv g'/g$, with θ_w the Weinberg angle, and A_μ is the photon field. The masses of the gauge bosons are

$$m_W = g \frac{v}{2}, \quad m_Z = \frac{v}{2} \sqrt{g'^2 + g^2}, \quad m_A = 0. \quad (1.23)$$

This process of the generation of masses for the $W^\pm$, and Z weak gauge bosons through electroweak symmetry breaking is what is called the Brout–Englert–Higgs mechanism [82, 140, 119, 129, 130]. Note that as consequence of the spontaneous breaking of the electroweak symmetry, the electric charge Q is related to the hypercharge Y and the third component of the isospin T_3 by

$$Q = Y + T_3. \quad (1.24)$$

After EWSB the Yukawa interactions from equation 1.15 also yield mass terms for the fermions

$$\mathcal{L}_{\text{Yukawa}} = -\frac{v}{\sqrt{2}} \left(Y_{ij}^e \bar{e}_{Li} e_{Rj} + Y_{ij}^d \bar{d}_{Li} d_{Rj} - Y_{ij}^u \bar{u}_{Li} u_{Rj} \right) + h.c., \quad (1.25)$$

with mass matrices

$$M_{ij}^f = \frac{v}{\sqrt{2}} Y_{ij}^f. \quad (1.26)$$

There is no mass term for the neutrinos in the SM. These mass matrices are not diagonal.

The diagonalization can be achieved by making a unitary transformation for left- and right-handed fields independently

$$f_L^i = (V_L^f)_{ij} f_L^{j(m)}, \quad f_R^i = (V_R^f)_{ij} f_R^{j(m)}, \quad (1.27)$$

where the subscript m is used to refer to the mass eigenstates. The diagonal mass reads

$$M_{ij}^{f(m)} = (V_L^f)_{ik} M_{kl}^f (V_R^{f\dagger})_{lj}, \quad (1.28)$$

and this procedure of diagonalization will also yield to a mix of up and down components for charged-current weak interactions

$$\mathcal{L}_{CC} = \frac{g}{\sqrt{2}} \bar{u}_L^{(m)} \gamma_\mu (V_L^u V_L^{d\dagger}) d_L^{(m)} W^\mu. \quad (1.29)$$

The product $V_L^u V_L^{d\dagger}$ defines the CKM matrix, which can be parametrized in all generality by three angles and one CP violating phase [190, 142]. Since neutrinos are massless, it is possible to absorb the leptonic transformation matrix, however we will show in the next chapter that neutrino mixing analogous to that in the quark sector arises in extensions of the SM with massive neutrinos.

1.3 Beyond the Standard Model

There are several theoretical motivations and experimental observations that call for an extension of the SM. In this section we will go through some of the drawbacks of the SM that motivate the extensions that are considered in this thesis. From a theoretical point of view, we will comment on the hierarchy problem and the charge quantization problem, while on the observational side, we will discuss the nature of dark matter, as well as the origin of neutrino masses.

1.3.1 The Hierarchy Problem and Supersymmetry

Expanding the potential of the Higgs in the SM around the VEV, allows to identify the Higgs mass

$$m_h^0 = \lambda \frac{v^2}{2}. \quad (1.30)$$

However this is only the tree-level contribution to the mass. The physical mass of the Higgs boson is also affected by the quantum corrections from every particle that

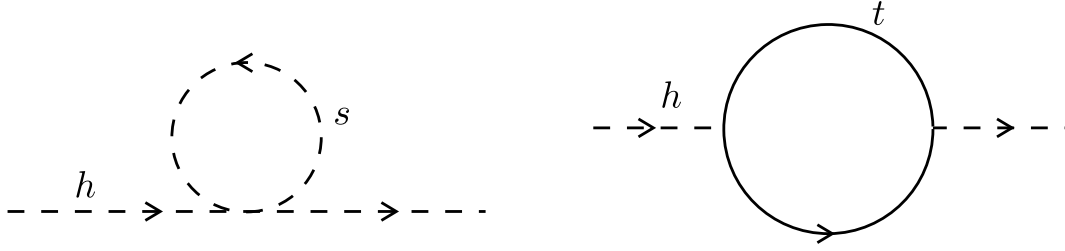


FIGURE 1.1: Loop diagrams obtained from the coupling of the Higgs with a scalar particle (s) in the left side, and a top quark in the right side.

couples directly or indirectly to the Higgs field. Consider the coupling of the Higgs with the top quark (right diagram from Figure 1.1). Applying a sharp momentum cut-off, Λ , as regulator of the loop integral, it is possible to show that this diagram gives a contribution to the Higgs mass proportional to

$$\delta m^2 \propto -y_t^2 \Lambda^2, \quad (1.31)$$

where y_t is the coupling of the Higgs with the top quark. Taking into account the value measured for the VEV, it is clear that this correction pushes the physical mass of the Higgs ($m_h^0 + \delta m^2$) to a value very close to the high-energy scale Λ , and in fact, since we expect that at least a new physics scale emerges when quantum gravity becomes relevant (around the Planck scale $\sim 10^{18}$ GeV), the measured value of the Higgs mass seems to imply a very fine tuning of the parameters. The problem posed by the gap between the electroweak scale ($\sim 10^2$ GeV) and the scale at which new physics emerges is known as the hierarchy problem.

Next, consider the case in which the Higgs boson couples to a complex scalar particle s by a $y_s |H|^2 |s|^2$ term. This kind of coupling leads to a loop diagram like the one in the left side of Figure 1.1. Once again using a sharp momentum cut-off as regulator, it can be shown that this loop diagram gives a mass contribution proportional to

$$\delta m^2 \propto +y_s \Lambda^2. \quad (1.32)$$

If each of the quarks and leptons of the SM is accompanied by two complex scalars fields with $y_s = y_f^2$, then the UV contributions of equations (1.31) and (1.32) can cancel against each other and there is no need of tuning the parameters related in order to explain the smallness of the Higgs mass compared to Λ . The idea of having an extra scalar (or fermionic) degree of freedom for each fermion (or boson) in the SM, is what lies behind supersymmetry [110, 189].

The space-time symmetry of the SM is described by the Poincaré group, which can be extended by adding generalized space-time transformations connecting fermions and bosons. Gauge invariant and supersymmetric field theories can be constructed respecting this extended space-time symmetry. The supersymmetric model that contains the SM and with the minimum set of new particles is called the Minimal Supersymmetric Standard Model (MSSM) [76]. Current research is conducted within extensions of the MSSM which can also address some of the others major problems

of the SM such as neutrino masses and dark matter. Nearly all of these extensions involve the existence of long-lived particles, meaning particles that could travel macroscopic distances before decaying. In particular, in this thesis we will study a supersymmetric extension involving long-lived, electrically neutral fermions called neutralinos.

We will study the production of these long-lived particles in atmospheric showers. In principle these particles could travel large distances, and decay inside the instrumented volume of large terrestrial neutrino detectors. It is expected that this kind of atmospheric searches can be able to probe particles with lifetimes that are of the order of a few kilometers in the laboratory frame. Such a scenario could not be probed in collider experiments, as particles with such a large lifetime would escape the detector, and in fact we will show that competitive constraints can be established for several models.

1.3.2 The Charge Quantization Problem and Millicharged Particles

As discussed previously, the electric charge of a given particle in the SM can be obtained from equation 1.24, and it depends on the third component of the isospin and the hypercharge. Notice that the latter is assigned to reproduce the experimental observation for the SM particle charges. We observe in nature that quarks are confined into hadrons which have charges that are multiples of the electron charge. Moreover, mesons are quark-antiquark pairs which have zero charge if made of two up, or two down pairs, or $\pm$ the electron charge if they mix up and down quarks. Similarly, baryons are made of three quarks, whether they are up or down, the charge is always a multiple of the electron charge. Therefore it appears that electric charge is quantized even though there is no principle dictating this rule. Natural explanations of the observed charge quantization principle are provided in extensions of the SM that involve the unification of forces [105], the existence of extra dimensions [141], or the existence of magnetic monopoles [77].

However, none of the above explanations has been demonstrated, and in fact, models with particles that have a tiny fraction of the charge of the electron,

$$Q = \epsilon e \ (\epsilon \lesssim 10^{-2}), \quad (1.33)$$

and which satisfy the symmetries of the SM have been built and studied in the literature for a long time, see for instance [171] and references therein. Particles with such a small unquantized electric charge are referred to as Millicharged Particles (MCPs). Direct searches for this kind of particle have been made by considering their effect in electromagnetic observables, as well as accelerator and astrophysical probes. Figure 1.2 shows a very early excluded region of parameter space considering different probes for MCPs in reference [79].

Searches for MCPs have been performed in many modern accelerator-based experiments [170, 64, 137, 13, 41], but also in the context of dark matter experiments [89]. In the case of colliders and beam-dump experiments, the masses of the MCPs that are usually probed lie within the MeV-GeV range, but the constraints for masses around (and above) 1 GeV are as weak as $Q \sim 10^{-2}e$ [20], leaving plenty of room for improvement in the future. On the other hand, the case of Millicharged Dark Matter usually focuses on masses below 1 MeV and it is highly constrained by several astrophysical observations including Big-Bang Nucleosynthesis, Stellar Evolution and the Cosmic Microwave Background [72].

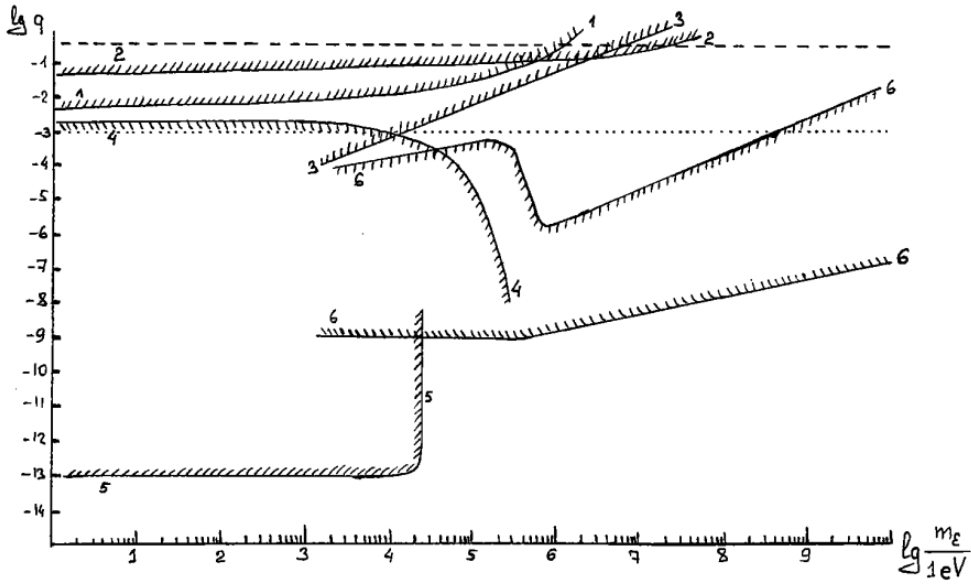


FIGURE 1.2: Region of parameter space in the plane charge-vs-mass of MCPs forbidden by consideration of their contribution to: 1 - anomalous -magnetic moment of the electron, 2 - anomalous magnetic moment of the muon, 3 - Lamb shift, 4 - heat transport in the core of a red giant, 5 - cooling time of a white dwarf, 6 - relic concentration. Figure from reference [79]

In this thesis we will study the possibility of searching for MCPs produced in the decay of mesons in cosmic-ray air-showers, and we will show that it is possible to establish stringent constraints using large underground neutrino detectors such as Super-Kamiokande or Juno [78].

1.3.3 The Nature of Dark Matter

There is a good amount of experimental evidence based on cosmological observations, which shows that 26% of the total energy budget of our universe corresponds to a form of non-baryonic matter which seems to interact only via the gravitational force and does not absorb nor emit light. This form of matter is the so-called Dark Matter (DM) [131, 48, 17, 16, 19]. The history of DM has dramatically evolved since the thirties, when it was found that the velocity dispersion of objects belonging to the distant Coma cluster were simply too large to be explained only by the distribution of visible matter in the cluster [193], up to nowadays where it is largely accepted that our model for cosmological evolution of the universe is well described by the Λ CDM model, which describes a universe filled with three major components, one of them being dark matter, playing a fundamental role in structure formation [93].

Observations of relic photons from the Cosmic Microwave Background (CMB) provide rich information about the matter distribution in the universe. Using information from the power spectrum of CMB anisotropies it is possible to fix the main parameters of the Λ CDM model. Of these parameters, the ratio $\Omega_i = \frac{\rho_i}{\rho_c}$ between the present value of energy density of species i , and the critical energy necessary to

make the Universe flat, has been measured accurately for the main components of the energy budget: [19]

$$\Omega_b h^2 = 0.02237 \pm 0.00015, \quad \Omega_c h^2 = 0.1200 \pm 0.0012, \quad \Omega_\Lambda h^2 = 0.3167 \pm 0.0082, \quad (1.34)$$

where $h = 0.678 \pm 0.009$ is the present expansion rate of the Universe in units of 100 km/Mpc/sec, and the subscripts b, c and Λ correspond to baryonic matter, cold (*i.e.* non-relativistic) dark matter, and dark energy, respectively. The later is by far the dominant component of our universe with the cosmological constant Λ being associated with the vacuum energy of space (or dark energy), which is responsible for the observed accelerating expansion of the universe [174, 164].

The mechanism by which DM populated the universe with the abundance that we observe today is still unknown, however there are two main popular mechanisms that we will briefly comment. The simplest mechanism is the so-called Thermal Freeze-Out [112], in which it is assumed that dark matter is composed of a new type of particle that at a very early epoch of the universe it was relativistic and in thermal equilibrium with the Standard Model bath. In this scenario, both DM and SM particles annihilate efficiently into each other $\text{DM} + \text{DM} \leftrightarrow \text{SM} + \text{SM}$. As the universe expands and cools down, there is a point in time where the interaction is no longer efficient, and the decoupling of species takes place when the interaction rate Γ drops below the Hubble expansion rate H . The abundance of DM freezes-out at this point. To obtain a quantitative estimate, a Boltzmann equation for the evolution of the number density must be solved. With a number of assumptions including the homogeneity and isotropy of our universe, it is possible to obtain a relatively simple expression for the evolution of the density of DM

$$\frac{dY}{dx} = -s \frac{x}{H} \langle \sigma v \rangle (Y^2 - Y_{eq}^2), \quad (1.35)$$

where Y is the comoving number density defined by $Y = \frac{n}{s}$, with n the DM number density and s the entropy density. The subscript $()_{eq}$ refers to the comoving number density at equilibrium. The variable $x \equiv m/T$ is the quotient between the DM mass and temperature, while $\langle \sigma v \rangle$ is the thermally-averaged DM pair-annihilation cross section times the relative velocity. The solid lines in figure 1.3 display the evolution of the comoving number density as a function of x . The drop takes place as the annihilation process is overwhelmed by the expansion of the universe, and the yield will decrease with the thermally averaged cross section.

On the other hand, there is a popular mechanism for attaining the DM relic density which does not involve an initial condition of equilibrium because the DM interacts with the SM particles too feebly to reach thermal equilibrium. This mechanism is called Freeze-in [121]. In this scenario, there is an initial negligible amount of DM, yet the feeble interaction would still yield a sizable production of DM through decays and annihilation of particles in the thermal bath. The DM progressively accumulates over time, until the temperature drops below its mass and the DM freezes-in with a yield that increases with the interaction strength of the particle with the bath. The dashed colored lines in figure 1.3 show the evolution of this process. Notice that in contrast with the freeze-out case, the yield increases with the coupling of DM with the particles in the thermal bath.

Many candidate particles have been proposed to explain DM, being the most popular ones those that are capable of solving other major problems in the SM, such as the

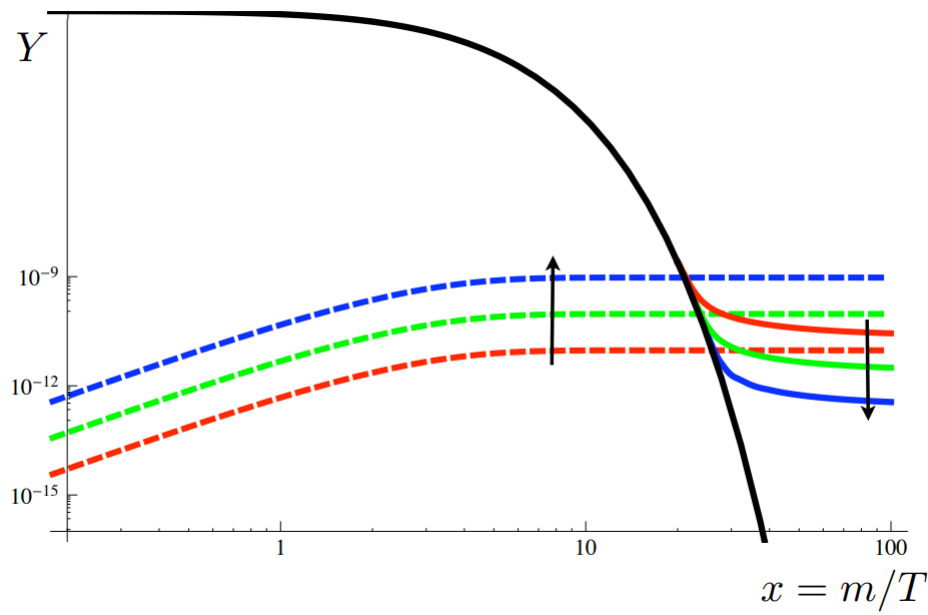


FIGURE 1.3: Evolution of the dark matter relic yields for conventional Freeze-out (solid lines) and Freeze-in (dashed lines), as a function of $x = m/T$. Figure taken from reference [121]. Curves corresponding to larger cross sections are indicated by the arrows.

hierarchy problem or the origin of neutrino masses among others. The possible DM mass range remains largely unconstrained. Some of the most popular lighter DM candidates includes light neutralinos, dark photons, axions and sterile neutrinos. Explaining the main features and implications of these candidates goes far beyond the scope of this thesis and we refer the reader to references [51, 147, 172, 90]. For higher masses at the GeV-TeV level, the most popular candidate are the so-called weakly interactive massive particles (WIMPs) that nicely fit the freeze-out scenario, and can provide a solution to the hierarchy problem at the same time. This is referred to as the “WIMP-miracle”, and good reviews on this subject can be found in references [176, 138, 179]. Finally for much larger DM masses, theorized black holes of primordial origin have been put forward as one of the most popular candidates [60, 116]. A Primordial Black Hole (PBH) is believed to be formed from the gravitational collapse of large density fluctuations in the early universe. Numerous cosmological and astrophysical constraints apply to this case, and a summary can be found in references [60, 58, 59]. It is generally believed that PBHs cannot constitute the whole DM.

At this point there is no doubt that DM exists and it is the seed for structure formation at large scales, however we still have very little knowledge about its nature and behavior. The main features of DM that we know are

- Dark matter is very long-lived, if not stable. We can observe the gravitational effects of DM today, and the observations of non-luminous matter in the universe implies that dark matter should have a lifetime at least as large as the current age of the universe (13.8 Gyr).
- Dark matter is non-baryonic. A Massive Compact Halo Object (MACHO) is

a kind of astrophysical body made of baryonic matter that emits very little or no radiation at all, and which travels through interstellar medium unbounded from any planetary system. Because of these features, it has been considered as a good baryonic candidate for DM, however searches for microlensing amplification of light have ruled out this possibility [162, 26, 183]. Regarding dark matter as primordial black holes (PBHs), this is a possibility that has not been ruled out completely, but since PBHs are formed in the radiation-dominated era, they are not subject to big-bang nucleosynthesis constraints, and are thus considered as non-baryonic [60].

- If dark matter does interact with the Standard Model, it must be through a feeble interaction. Strongly interacting dark matter is ruled out because it would yield to a heating rate for the earth far beyond what is measured [150]. On the other hand, DM with a standard electric charge would not be compatible with current astrophysical observations, yet a DM component with a small electric charge ϵq , ($\epsilon \ll 1$), is highly constrained but still possible [146, 10, 149]. Finally, because of their elusive nature, neutrinos interacting with DM is the less constrained scenario.
- Dark matter self-interaction is constrained. This is obtained from observations of merging clusters [68, 173] from which a constraint on the ratio of the DM-DM cross section and the dark matter mass can be obtained [125].

Dark matter searches are typically classified according to three broader categories; direct detection, indirect detection and collider searches. In direct detection experiments, the goal is to observe the signal from the recoil of a target nucleus, produced from the scattering of galactic DM with a target medium. These kind of experiments are usually located underground in order to suppress the cosmic-ray background. On the other hand, indirect detection experiments aim to observe SM particles produced from the decay or annihilation of dark matter in outer space, being the detection of radiation the most common type of signal. The most challenging aspect of this kind of search is related to the simulation of the background. Finally, in collider searches the most traditional strategy consists of searching for significant amounts of missing energy and momentum, as it is expected that once the proton-proton collision takes place, DM is produced but escapes the detector. This approach is usually model dependent, however a considerable effort is being devoted to these searches. In particular, light DM produced in extended dark sectors involving long-lived particles, is under active search with new proposed experiments such as The Search for Hidden Particles (SHiP) experiment [29, 22], and the MAAssive Timing Hodoscope for Ultra Stable neutral pArticles (MATHUSLA) [71, 67].

In this thesis we will consider scenarios with extended dark sectors, but focusing on searches in cosmic-ray air-showers.

1.3.4 The Origin of Neutrino Masses

It was not until the nineties when the observation of flavour change in solar and atmospheric neutrinos at Super-Kamiokande [96] and SNO [24], established that neutrinos are massive particles. As a result flavour eigenstates which take part in electroweak interactions are not identical to the mass eigenstates, the later being quantum superpositions of the former

$$|\nu_i\rangle = \sum_{\alpha} U_{i\alpha} |\nu_{\alpha}\rangle, \quad (1.36)$$

where $i = 1, 2, 3$ and α denotes the flavours e, μ, τ . The matrix U is the so-called Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix [169, 153], and encodes the unitary relation between the weak and massive eigenstates in a similar fashion as the CKM matrix. The PMNS matrix can be parametrized with three mixing angles and one CP violating phase δ ¹

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}c_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix}, \quad (1.37)$$

where $c_{ij} \equiv \cos \theta_{ij}$, and $s_{ij} \equiv \sin \theta_{ij}$, with $\theta \in [0, \pi/2]$. The oscillation of neutrinos was first introduced by Pontecorvo [167].

The probability that a neutrino of flavour α oscillates to a flavour β after traveling a distance (called baseline) L is

$$P(\nu_\alpha \rightarrow \nu_\beta) = \sum_{ij} U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^* \exp \left(-i \frac{\Delta m_{ji}^2 L}{2|\mathbf{p}|} \right), \quad (1.38)$$

where $\Delta m_{ji}^2 = m_i^2 - m_j^2$ is the neutrino mass squared difference, and $\mathbf{p}$ is the neutrino momentum. The formula clearly involves a superposition of waves with amplitude depending on the mixing through the PMNS matrix, and wavelengths that depend on the mass differences.

The study of neutrino oscillations provides the measurements of the neutrino parameters involved in equation 1.38, via experiments using neutrino sources of natural origin; like the sun or the atmosphere, and also produced on terrestrial experiments including beam-dump experiments and reactors. Of particular importance for this thesis is the study of atmospheric neutrinos which will be treated in the next chapter, however the reader can consult references [178, 114] for updated values on the parameters describing neutrino oscillations. It should be noted that neutrino oscillations do not provide information on the absolute value of neutrino masses, but on the mass squared differences.

Although the simplest extension of the SM that includes massive neutrinos is pretty straightforward to write down, by adding a right handed component ν_R for each of the neutrino species, the way to introduce a neutrino mass term is not unique. Unlike the case for charged fermions, it is possible to construct a gauge invariant neutrino mass term in two different forms depending on whether lepton number L is preserved or not. Notice that in the SM, both lepton number L and baryon number B are accidental symmetries of the SM.

If lepton number is preserved, a Dirac neutrino mass term can be built simply by adding three right-handed neutrinos ν_R which are SM gauge singlets. This allows to write a Yukawa coupling term

$$\mathcal{L} = -Y_{ij}^\nu \bar{\ell}_{Li} \tilde{H} \nu_{Rj} + h.c. \quad (1.39)$$

After EWSB, the neutrinos mass matrix obtained is proportional to the Higgs VEV, in complete analogy with the mass formula for the rest of the Standard Model's fermions in equation 1.26

$$m_D = \frac{v}{\sqrt{2}} Y_\nu. \quad (1.40)$$

¹This parametrization implicitly assumes that neutrinos are Dirac fermions, for Majorana fermions there are two additional phases.

On the other hand, if lepton number is violated, then it is possible to add a Majorana mass term

$$\mathcal{L}_s = -Y_{ij}^\nu \bar{\ell}_{Li} \tilde{H} \nu_{Rj} - \frac{1}{2} (m_M)_{ij} \bar{\nu}_{Ri}^c \nu_{Rj} + h.c., \quad (1.41)$$

where $\nu^c = C \bar{\nu}^T$, with C the charge conjugation matrix. The Lagrangian from equation 1.41 describes the so-called Seesaw type-I mechanism [156, 192, 104], which is the most popular alternative for generating neutrino masses.

After EWSB, the Lagrangian in equation 1.41 leads to the mass matrix

$$\mathcal{M} = \begin{pmatrix} 0 & m_D \\ m_D^T & m_M \end{pmatrix}. \quad (1.42)$$

This matrix can be easily diagonalised, and the resulting mass eigenstates are given by admixtures of the active and sterile fermions. Assuming that $m_M \gg m_D$, we find two approximate mass matrices given by

$$m_{light} \approx m_D (m_M)^{-1} m_D^T, \quad m_{heavy} \approx m_M. \quad (1.43)$$

In the equation above the subscripts *light* and *heavy* are not randomly assigned. The matrix m_M is expected to be associated with a very high energy scale, thus equation 1.43 implies a spectrum with heavy neutrino states N_j with mass $\sim m_M$, and three light mass eigenstates ν_i with masses that are suppressed by the inverse of m_M . We observe that m_{heavy} increases with m_M , while m_{light} decreases, and that is why the model is called seesaw.

We stress, however, that particular realisations of the seesaw mechanism described above can lead to phenomenologically meaningful scenarios with heavy neutrino masses at the GeV scale and below. An example of such realisations is given by the inverse seesaw model [158, 159]. Our approach in this thesis is phenomenological and model independent, so we do not make any specific assumption regarding the relationship between the neutrino mixing and the mass of the sterile neutrino. In future chapters we will discuss in detail about searches for relatively heavy sterile neutrino states, and we will report the results of atmospheric searches based on simple assumptions about their production and detection mechanisms.

Chapter 2

Cosmic Rays and Atmospheric Showers

Cosmic Rays (CRs), are highly energetic particles traveling through space-time at relativistic velocities. About 90% of them are protons, and the rest is mostly alpha particles ($\sim 9\%$), and other heavier nuclei including Carbon, Oxygen, among others. The energy spectrum covers a wide range of energy scales, from about 10^9 eV, up to 10^{20} eV and above. To measure such a broad energy range there are different detection techniques. For energies in the range $\sim 1 - 10^4$ GeV satellite experiments such as the Alpha Magnetic Spectrometer (AMS) [21], and the Payload for Antimatter Matter Exploration and Light-nuclei Astrophysic (PAMELA) experiment [18], have performed direct measurements of cosmic rays in outer space by using an array of components which provides information of the momentum and charge of the passing particles. On the other hand, very high energy CRs ($> 10^4$ GeV), are measured indirectly by terrestrial Air-Shower experiments such as the Pierre Auger Observatory [2], and KASCADE-Grande [30]. These terrestrial experiments can measure a fraction of the surviving particles produced as secondaries from the collision of the primary CRs with the atmosphere and use this information to infer and reconstruct the incoming cosmic ray properties including composition, energy and direction. Figure 2.1 shows the spectrum of cosmic rays as a function of the kinetic energy of the particles, as measured by several experiments including balloons, satellites, and ground based arrays. Notice that the vertical axis of the figure corresponds to the square of the energy of the particle, times the rate of particles arriving per unit of kinetic energy, area and solid angle. Moreover, it is worth stressing that for energies less than 10 GeV per nucleon, the measured CR flux can change during the solar cycle, an effect referred to as “solar modulation” [66], which will be discussed later.

The origin of cosmic rays is not yet known and it is an active field of research, however it is expected that most of them are originated within our Galaxy, as a result of the diffuse shock acceleration of particles in Supernova remnants [15]. Other sources of CRs include quasars, active galactic nuclei, and gamma-ray bursts. The spectrum in figure 2.1 cannot be explained by a single power law, indicating a multi-component origin and acceleration of CRs. The low energy part of the spectrum ($10-10^3$ GeV) is well modeled by a power law with $E^{-\gamma}$, with the spectral index, $\gamma \approx 2.7$ [102]. At higher energies around 3 PeV there is a steepening of the spectrum denominated “the knee”, which can be fitted by a different spectral index $\gamma \approx 3.1$. Finally, it is observed that for ultrahigh energy CRs at the EeV level there is a softening of the spectrum, the so-called “ankle”, with $\gamma \approx 2.6$. Notice that the spectrum of figure 2.1 indicates that the flux of cosmic rays is about 1 particle per $\text{m}^2\text{-s}$ below the knee, 1 particle per $\text{m}^2\text{-yr}$ around the knee, and only 1 particle per $\text{km}^2\text{-yr}$ around the ankle, so this is the reason why in order to detect high-energy CRs very large

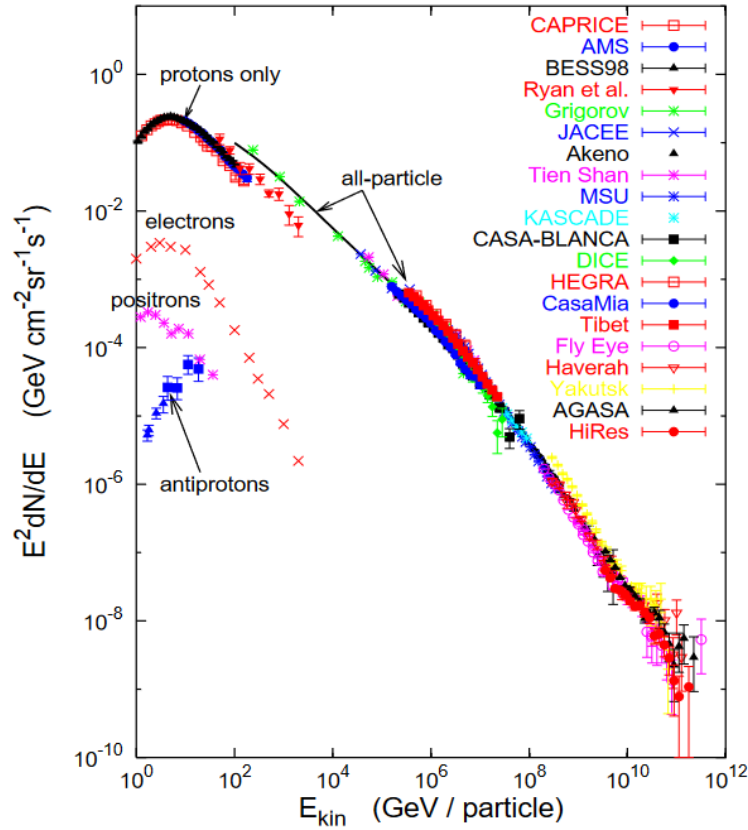


FIGURE 2.1: All particle cosmic ray energy spectrum as a function of the particle's kinetic energy, measured by different terrestrial and satellite experiments as indicated in the legend. Figure from reference [98].

terrestrial detector arrays are needed.

2.1 Air Showers

When cosmic ray particles reach the Earth, they collide with the atmosphere, primarily with nitrogen and oxygen atoms through a deep inelastic scattering process. This collision produces an extensive air shower or cascade of lighter particles including protons, neutrons, alpha particles, hadrons, photons (x-rays), and leptons [40, 117]. The height at which the collision of primary CRs with atoms of the atmosphere takes place varies with the atmosphere's density, but it is somewhere around 10-20 km above the surface of the Earth. The main components of any atmospheric shower are the electromagnetic and hadronic shower, which will be described below.

Hadronic showers are the core component of a cosmic ray-induced air shower, the hadrons are produced directly in the collision of the primary CRs with atmospheric nuclei through strong interactions. About 90% of the hadrons produced in these interactions are pions (with roughly the same proportion of $\pi^0, \pi^\pm$) [163], kaons and light baryons. As these secondary particles propagate through the atmosphere, they can decay and interact with other nuclei producing numerous new particles that drive the spread of the shower, which will take place until the energy per nucleon is

less than ~ 1 GeV (minimum energy for pion production). The decay of pions gives rise to two important components; muons and electromagnetic showers, the most important reactions are

$$\begin{aligned}\pi^+ &\rightarrow \mu^+ + \nu_\mu, \\ \pi^- &\rightarrow \mu^- + \bar{\nu}_\mu, \\ \pi^0 &\rightarrow 2\gamma.\end{aligned}\tag{2.1}$$

The right panel in figure 2.2 shows an illustration of an hadronic shower with pions and other hadrons being formed after the collision of a primary cosmic ray with an atmospheric nucleus, π^0 will give rise to the electromagnetic component, while $\pi^\pm$ will contribute to the neutrino and muon flux which can reach the surface of the Earth. Subsequent collisions of secondaries in the atmosphere will also contribute to the development of the shower.

On the other hand, the electromagnetic shower is fed by the hadronic shower through photoproduction. The photons interact with matter in the atmosphere generating electron-positron pairs, which at the same time emit more photons via bremsstrahlung radiation. The left panel in figure 2.2 illustrates these processes. Since most of the hadrons in the shower either quickly decay or re-interact, most of the energy of the primary CRs eventually finds its way into the electromagnetic component.

Many of the important features of the cascade can be explained by a simple toy model introduced by Heitler in the thirties [127], where the cascade is described by particles of the same type interacting in the medium with an interaction length λ . Such interactions will create two new particles, each with half of the initial energy E_0 from the primary particle. Once the shower reaches a certain depth $X = n\lambda$, there is a total of $N = 2^n$ particles, and this process continues until a critical energy per particle E_c is attained (below this the cross-section for the interaction is simply too small), and the number of particles ceases to grow. The maximum number of particles in the shower is

$$N_{max} = \frac{E_0}{E_c},\tag{2.2}$$

and the maximum depth of the shower is proportional to the logarithm of the ratio between the initial energy and the critical energy

$$X_{max} = \lambda \frac{\ln(E_0/E_c)}{\ln 2}.\tag{2.3}$$

The evolution of the number of particles is not described in this toy model and it depends on the strength of the energy loss. Although this model is extremely simple, it provides a good qualitative description of both electromagnetic and hadronic showers. In practice air showers are modeled either by Monte Carlo techniques [126] or by numerical solutions of the so-called Cascade Equations [154, 155, 97], with the later being the most popular alternative due to speed and efficiency in the resolution.

2.2 Cascade Equations

The development of the particle content in an air shower can be described by a set of coupled integro-differential equations referred to as cascade equations. The main

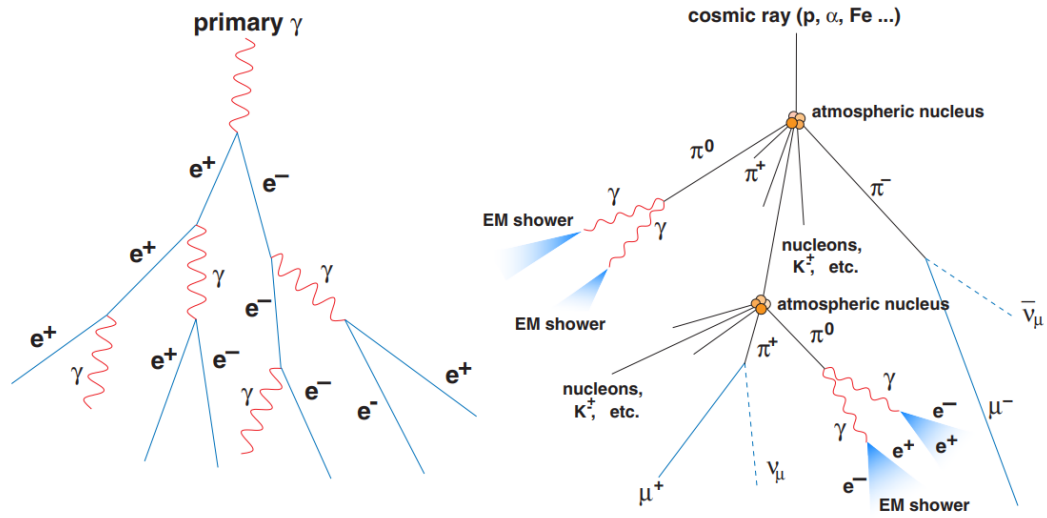


FIGURE 2.2: Representation of the main components of an air-shower. **Left:** Electromagnetic shower. **Right:** Hadronic shower. Figure from [186].

goal is to determine the spectral rate of particles of type h

$$\Phi_h = \frac{dN_h}{dEdtdAd\Omega}, \quad (2.4)$$

which is the number of particles per unit energy, time, unit area and solid angle, typically expressed in $[\text{GeV s cm}^2 \text{ sr}]^{-1}$.

The development of an air shower is usually described in terms of the slant depth X , which indicates the amount of atmosphere traversed by the particles in the shower. This quantity is calculated by integrating the density of air at distance ℓ from the ground

$$X(\ell) = \int_0^\ell dl \rho(l). \quad (2.5)$$

Consider the change in the number of particles of type h over a small piece of atmospheric volume. The number of particles can decrease when the particles either decay to another type of particle i , or when it annihilates by interacting with air nuclei. On the other hand, the number of particles of type h can also increase when a particle of type i enters the atmospheric volume and produces a particle of type h , via decay or scattering. The cascade equation takes into account all of these possible

interactions

$$\begin{aligned}
\frac{d\Phi_h}{dX}(E, X) = & -\frac{\Phi_h(E, X)}{\lambda_{int,h}(E)} \\
& -\frac{\Phi_h(E, X)}{\lambda_{dec,h}(E, X)} \\
& -\frac{\partial}{\partial E}(\mu(E)\Phi_h(E, X)) \\
& +\sum_i \int_E^\infty dE_i \frac{dN_{i \rightarrow h}^{int}}{dE}(E_i, E) \frac{\Phi_i(E_i, X)}{\lambda_{int,i}(E_i)} \\
& +\sum_i \int_E^\infty dE_i \frac{dN_{i \rightarrow h}^{dec}}{dE}(E_i, E) \frac{\Phi_i(E_i, X)}{\lambda_{dec,i}(E_i, X)}, \tag{2.6}
\end{aligned}$$

The first term in equation 2.6 involves the losses due to interactions. The interaction length λ_{int} measures the average distance traveled by the particle between two successive collisions

$$\lambda_{int,h}(E) = \frac{\langle m_{air} \rangle}{\sigma_{h-air}^{inel}(E)}, \tag{2.7}$$

with average air mass¹, $\langle m_{air} \rangle = 14.5 m_p$, and $m_p = 1.24 \times 10^{-24}$ g the proton mass. The interaction length is of the same order of magnitude for most of the hadrons produced in the cascade and it changes very slowly with the energy. Figure 2.3 shows the interaction length for pions, which practically overlaps with the interaction length of other mesons such as kaons.

On the other hand, the second sink term in equation 2.6 involves the decay length, which is a quantity that can change by several orders of magnitude with the energy as it is proportional to the distance traveled by the particle before decaying in the rest frame

$$\lambda_{dec,h}(E, X) = E \frac{c\tau_h}{m_h} \rho(X), \tag{2.8}$$

where τ_h and m_h are respectively the lifetime, and mass of the particle type h . The high variation in this quantity is due to the dilation of the lifetime of the relativistic particle. Figure 2.3 illustrates this feature for several mesons, notice that there is also a big change in the decay length from one particle species to another due to the value of the particle's lifetime itself.

Finally, the last sink term involves the energy loss due to bremsstrahlung and the ionization of charged particles. The Bethe-Bloch formula includes the stopping power $\mu(E)$ for the charged particle, with typical losses at the MeV/g/cm² level [88, 181]. However, since in this thesis we are usually concerned about flux calculations at higher energies (> 1.0 GeV), this loss term is sub-leading and is usually ignored [111].

On the other hand, the first source term in equation 2.6 involves the production of type h particles from the interaction of type i particles with air molecules. The

¹We consider an ordinary air molecule by means of the average atmospheric composition based on oxygen and nitrogen nuclei, with an average atomic number of $\langle A_{air} \rangle = 14.5$.

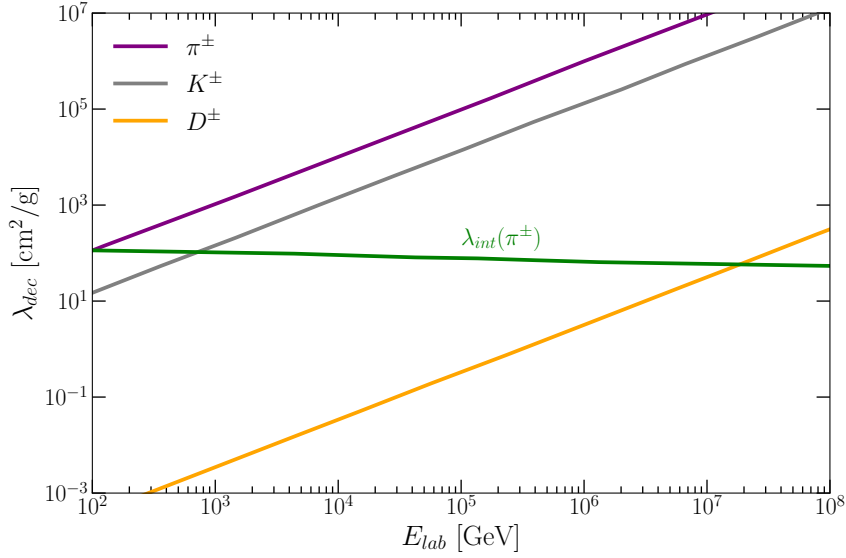


FIGURE 2.3: Decay length and Interaction length for several mesons produced in atmospheric showers as a function of the energy. Figure reproduced from reference [88].

distribution involves its differential cross-section normalized to the total inelastic cross-section of type i particles with air molecules, as given by

$$\frac{dN_{i \rightarrow h}^{int}}{dE}(E_i, E) = \frac{1}{\sigma_{inel, i-air}(E_i)} \frac{d\sigma_{i \rightarrow h}(E_i, E)}{dE}. \quad (2.9)$$

The second source term in equation 2.6 involves the fraction of particles that can decay to h , as well as the differential energy distribution of the particle produced in the decay. Taking into account that this can take place for different channels, this equation reads

$$\frac{dN_{i \rightarrow h}^{dec}}{dE}(E_i, E) = \sum_j \text{Br}_{j, i \rightarrow h} \frac{dN_{j, i \rightarrow h}}{dE}(E_i, E), \quad (2.10)$$

where the index j runs over every possible decay channel.

Integrating equation 2.6 with respect to the slant depth yields an estimate of the energy-dependent flux at the surface of the earth, which can be contrasted with experimental measurements. A qualitatively good solution for the cascade equations can be found in an analytical way by considering a power law primary spectra for the CRs, and a scale-invariant interaction cross sections [184, 148, 102, 136]. However, in this thesis we will focus on numerical solutions based on the Matrix Cascade Equation (MCEq) software [86, 87, 85], which provides state-of-the-art results for the flux spectrum of hadrons and leptons. MCEq uses an efficient numerical computing scheme by rewriting the cascade equations into matrix form, where the energy values are binned in discrete steps and each particle type can be grouped into a column

vector

$$\Phi = \begin{pmatrix} \Phi_{h_1} \\ \vdots \\ \Phi_{h_k} \\ \vdots \\ \Phi_{h_p} \end{pmatrix} \longrightarrow \Phi_{h_k} = \begin{pmatrix} \Phi_{h_k}(E_0) \\ \Phi_{h_k}(E_1) \\ \vdots \\ \Phi_{h_k}(E_N) \end{pmatrix} \quad (2.11)$$

where the indices in $h_1, h_2, \dots, h_p$ are used to identify different particle types, which includes more than 50 types of mesons, baryons and leptons. The energy integrals in equation 2.6 can be replaced by a restricted summation constrained to satisfy energy conservation, while the interaction and decay terms from equations 2.9 and 2.10 can be rewritten in terms of interaction and decay coefficients

$$c_{i(E_i) \rightarrow h(E_h)} = \Delta E_i \frac{dN_{i \rightarrow h}^{int}}{dE_h}(E_i, E_h), \quad (2.12)$$

$$d_{i(E_i) \rightarrow h(E_h)} = \Delta E_i \frac{dN_{i \rightarrow h}^{dec}}{dE_h}(E_i, E_h), \quad (2.13)$$

which will yield to matrices with a structure containing an upper-triangular form due to energy conservation

$$C_{i \rightarrow h} = \begin{pmatrix} c_{i(E_0) \rightarrow h(E_0)} & \cdots & c_{i(E_0) \rightarrow h(E_N)} \\ & & c_{i(E_1) \rightarrow h(E_N)} \\ & \ddots & \vdots \\ 0 & & c_{i(E_N) \rightarrow h(E_N)} \end{pmatrix}, \quad (2.14)$$

$$D_{i \rightarrow h} = \begin{pmatrix} d_{i(E_0) \rightarrow h(E_0)} & \cdots & d_{i(E_0) \rightarrow h(E_N)} \\ & & d_{i(E_1) \rightarrow h(E_N)} \\ & \ddots & \vdots \\ 0 & & d_{i(E_N) \rightarrow h(E_N)} \end{pmatrix}. \quad (2.15)$$

For each type of particle h produced in the atmospheric shower, we can incorporate these coefficient matrices and assemble a full interaction matrix $\mathbf{C}$, and a decay matrix $\mathbf{D}$.

Regarding the interaction lengths from equation 2.6, these can be arranged in a diagonal matrix Λ_{int} , sorted with the same order of the particle species appearing in the column vector flux of equation 2.11

$$\Lambda_{int} = \text{diag} \left(\frac{1}{\lambda_{int}^{h_1}}, \dots, \frac{1}{\lambda_{int}^{h_p}} \right), \quad (2.16)$$

where

$$\frac{1}{\lambda_{int}^{h_k}} \equiv \left(\frac{1}{\lambda_{int, E_0}^{h_k}}, \dots, \frac{1}{\lambda_{int, E_N}^{h_k}} \right) \quad (2.17)$$

for $k = 1, \dots, p$. Similarly, we can define a matrix for the decay lengths Λ_{dec} , however since they also depend on the slant depth X , we can factor out this dependence by making $\tilde{\lambda}_{dec} = \lambda_{dec}/\rho(X)$, such that the decay length matrix reads

$$\Lambda_{dec} = \text{diag} \left(\frac{1}{\tilde{\lambda}_{dec}^{h_1}}, \dots, \frac{1}{\tilde{\lambda}_{dec}^{h_p}} \right). \quad (2.18)$$

Using this notation and ignoring the energy loss term, the cascade equation can be expressed by

$$\frac{d}{dX} \Phi = \left[(C-1)\Lambda_{int} + \frac{1}{\rho(X)}(D-1)\Lambda_{dec} \right] \Phi. \quad (2.19)$$

To solve this equation numerically, several tables including hadron-air cross sections, and decays are needed as inputs, we will discuss these details later on.

2.3 Atmospheric Muons and Neutrinos

Atmospheric neutrinos and muons are copiously generated by two-body decays of charged pions and kaons, produced in the interaction of CRs with nucleons in Earth's atmosphere. Since the muon is long-lived enough and the neutrino is a stable particle, these are the main components from air-showers that arrive to the surface of the Earth. In this section we will discuss the production and detection of these leptons.

2.3.1 Atmospheric Muons

Muons are perhaps the most important link between cosmic rays and particle physics. They were discovered in the thirties by Anderson and Neddermeyer while using a cloud chamber to study cosmic radiation at sea level [160]. Multiple measurements shows that muons arriving to the surface of the Earth from vertical directions are given at a rate of one particle per cm^2 per minute [55, 73, 120]. Muons give the dominant signal at sea level and can even penetrate deeply underground, becoming the main source of background for many underground experiments.

The muon energy distribution results from the combination of production spectrum and energy losses. A general characterization of the spectrum can be established by a comparison of its total energy with respect to the critical energy to reach the Earth surface from a production point h_0 above: $\epsilon_\mu = \frac{m_\mu c^2}{c\tau_\mu} h_0 \sim 1.0 \text{ GeV}$ [102]. We can distinguish three different energy regimes for the muon spectrum:

1. For $E_\mu \leq \epsilon_\mu$, the energy spectrum is mostly flat, and steepens gradually around 10 GeV in the same way as the primary spectrum. In this energy regime the energy losses and the decay of the muon are relevant and cannot be neglected in the cascade equation [54, 133].
2. For $\epsilon_\mu \leq E_\mu \leq \epsilon_{\pi,K}$, where $\epsilon_\pi \sim 115 \text{ GeV}$, and $\epsilon_K \sim 850 \text{ GeV}$, most of the mesons decay so the spectrum of muons follows the same as of their parent mesons (which is the same as the primary CRs).
3. For $E_\mu \gg \epsilon_{\pi,K}$, there is a high boost factor. The decay probability for mesons is suppressed, and even though the meson spectrum has the same power law as the CRs, the rate of their decay steepen with one power of E_μ .

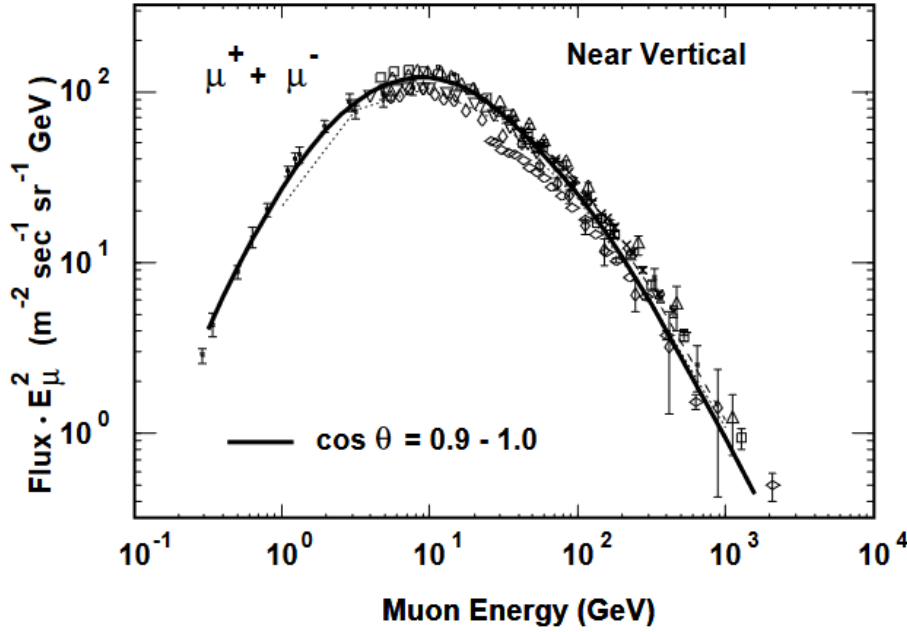


FIGURE 2.4: Low energy, near vertical muon flux spectrum measurements (dots), compared with detailed flux calculations made by Honda et al. From reference [133].

The low energy piece of the muon flux spectrum can be seen in figure 2.4, which shows an early, yet very detailed semi-analytical flux calculation made by M. Honda et al. [133], for muons arriving with directions given by the cosine of zenith angle θ , with $\cos \theta = 0.9 - 1.0$. This calculation includes relevant low-energy effects including solar modulation and the geomagnetic field, which will be discussed in the next section.

Although in these early studies the uncertainty in the flux was estimated to be at the level of 10%-20%, there has been progress in the description of the spectrum by the development of detailed Monte Carlo codes such as CORSIKA and FLUKA [92, 84], the later being a very popular tool for particle transport which is also used in collider studies. The uncertainties are dominated by the unknowns in the primary spectrum, and the treatment of hadronic interactions. A more detailed discussion regarding uncertainties will be made in the next chapter. We notice, however, that none of these tools were well suited for estimations at very high energies, and moreover the Monte Carlo methods involves a high computational cost. This motivates the use of numerical solvers which includes high energy simulations such as MCEq. The cascade equation for muons at high energies takes the simple form

$$\frac{d\Phi_\mu}{dX}(E, X) = \sum_M \int_E^\infty dE_M \frac{dN_{M \rightarrow \mu}^{dec}}{dE} \frac{\Phi_M(E_M, X)}{\lambda_{dec, M}(E_M, X)}, \quad (2.20)$$

where the sum in M runs over the mesons in the shower that decay to muons. MCEq allows the possibility to choose among different cosmic-ray models and hadronic interaction models, moreover it allows to estimate the partial contribution to the flux from each meson. For partial contributions to the muon flux we distinguish two components; the prompt component contribution originates from very short lived

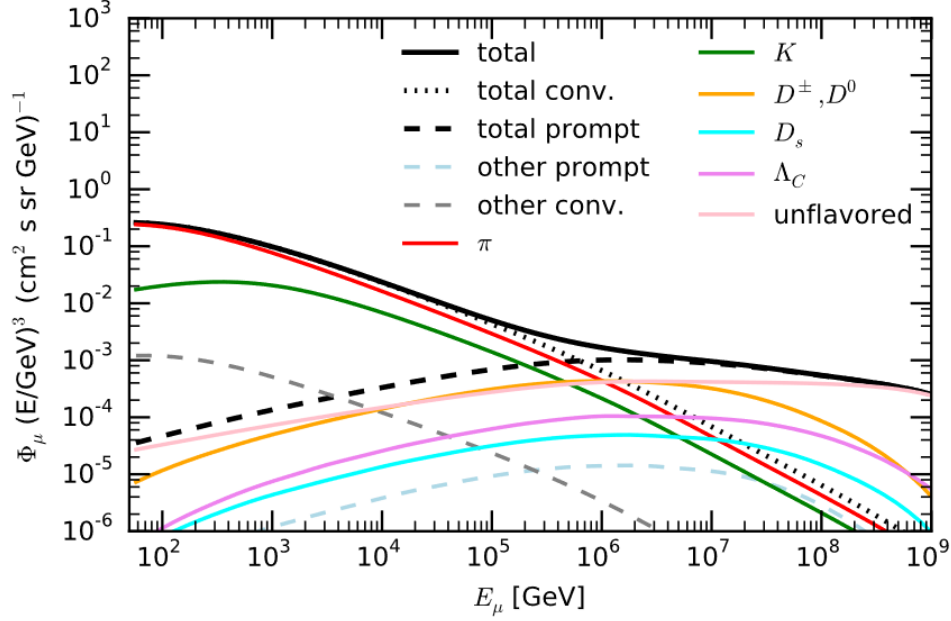


FIGURE 2.5: High energy, near vertical, flux spectrum of muons scaled by E^3 . The graphic includes partial hadron contribution as indicated in the legend. The solid black line denotes the total flux, while the black dotted lines indicate the total contribution from the so-called conventional component. The black dashed line is used to denote the total prompt component from charmed hadrons. Figure obtained using the MCEq software.

mesons and baryons containing charm quarks such as D-mesons and Λ_c^+ . While the conventional component is originated mainly from the decay of pions and kaons. Figure 2.5 illustrates the muon spectrum with conventional and prompt contribution for high energies up to 1 EeV. As can be seen from this figure, the prompt component becomes dominant at higher energies starting at 10^6 GeV. The muon spectrum in figure 2.5, is obtained using the Thunman-Ingelman-Gondolo (TIG) cosmic ray model [111], and the SIBYLL-2.3 hadronic interaction model [87]².

2.3.2 Atmospheric Neutrinos

Similar to muons, atmospheric neutrinos are mostly produced in the decays of pions and kaons. However, at lower energies there is also an important component from muon decays. The full decay chain from pions producing both electron and muon neutrinos is

$$\pi^\pm \rightarrow \mu^\pm + \nu_\mu(\bar{\nu}_\mu) \rightarrow e^\pm + \nu_e(\bar{\nu}_e) + \bar{\nu}_\mu(\nu_\mu) + \nu_\mu(\bar{\nu}_\mu), \quad (2.21)$$

which would imply that if most of the muons decay there should be a factor of 2 in the ratio of muon neutrinos to electron neutrinos. However, in 1994 the Kamiokande

²The impact of cosmic ray models and hadronic interactions in flux calculations will be discussed in the next chapter.

experiment found that this ratio is approximately equal to 1.0, giving rise to the so-called atmospheric neutrino anomaly [95, 132, 61, 47]. As we will see, the explanation for this observation relies on the oscillation of neutrinos. Starting from the general neutrino oscillation formula, it is possible to show that in the simplest case of two-family mixing the oscillation probability is [113, 139, 128]

$$P(\nu_\alpha \rightarrow \nu_\beta) = \sin^2 2\theta \sin^2(\phi_{ij}), \quad (2.22)$$

with $\alpha \neq \beta$, and ϕ_{ij} the neutrino oscillation phase

$$\phi_{ij} = 1.27 \frac{\Delta m^2 (\text{eV}^2) L (\text{km})}{E_\nu (\text{GeV})}. \quad (2.23)$$

This approximation provides a good description of the $\nu_\mu \rightarrow \nu_\tau$ oscillations relevant for atmospheric neutrinos. In equation 2.22 the amplitude of the oscillation depends on the mixing angle, and it is maximal for $\theta = \pi/4$. The frequency of the oscillation is controlled by Δm^2 and L/E_ν . The oscillation probability is appreciable if the kinematical phase is of order one or larger. For instance, an experiment studying neutrinos with energies around 1 GeV, which travel a distance comparable to earth's diameter $L \sim 10^4$ km, will be sensitive to neutrino mass squared splittings as small as 10^{-4} eV². Figure 2.6 displays the oscillation probability of muon neutrinos to tau neutrinos as a function of the energy for different baseline distance L , for $\sin^2 \theta_{23} = 0.44$, and $|\Delta m_{23}^2| = 2.4 \times 10^{-3}$ eV².

For atmospheric neutrinos the relevant difference of masses is of order $\Delta m^2 \sim 10^{-3}$ eV², and considering a typical neutrino energy of around 10 GeV, the oscillations are visible for neutrinos propagating over hundred of kilometers, going through the Earth from below the horizon. Super-Kamiokande established the existence of neutrino oscillations with almost maximal mixing of muon neutrinos and tau neutrinos, by finding that the number of upward going muon neutrinos, *i.e.* those generated on the other side of the Earth, are half of the number of downward going muon neutrinos, implying the disappearance of this type of neutrino in the former scenario. The atmospheric mass splitting value is presently measured to be $|\Delta m_{23}^2| = (2.40^{+0.11}_{-0.12}) \times 10^{-3}$ eV² [182], while the atmospheric mixing angle is $\sin^2 \theta_{23} = 0.44 \pm 0.008$, with both measurements reported for normal ordering, where ν_1 is the lightest state while ν_3 is the heaviest one.

At high enough energies, neutrino fluxes can be computed using a cascade equation of the type given in equation 2.20. Just as with muons, using MCEq it is possible to estimate the individual contribution from each progenitor to the neutrino flux. In the case of tau neutrinos, the dominant contribution is from tau leptons in the entire energy range. On the other hand, electron and muon neutrinos receive a dominant contribution from kaons in the $10^2 - 10^6$ GeV energy range, and at higher energies the dominant contribution comes from D-mesons, which dominates the prompt component of the flux.

The low energy part of the spectrum requires a more delicate treatment, since the atmospheric neutrino flux peaks around 100 MeV, where geomagnetic effects and solar activity are relevant. These effects will be discussed in the next section. Figure 2.7 shows the vertical flux spectrum of both, electron and muon neutrinos, for energies ranging from 10 MeV up to 10^8 GeV. Details of the partial contributions included can be found in the caption. This figure includes the geomagnetic and solar effects using local measurements reported by Super-Kamiokande.

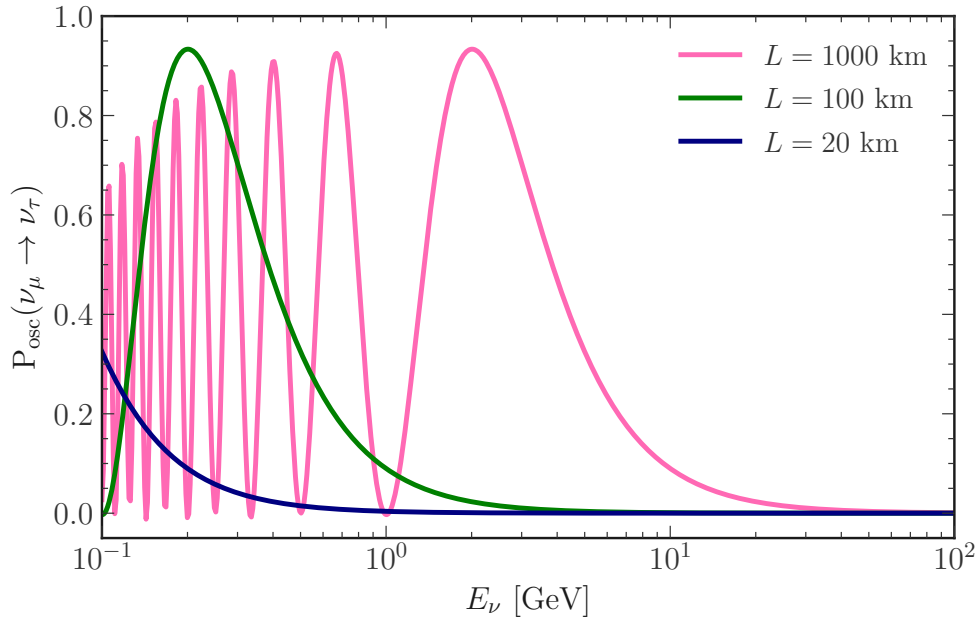


FIGURE 2.6: Probability of oscillation for a muon neutrino to become a tau neutrino as a function of the neutrino energy, for three different baseline distances as indicated in the legend. We have set $\sin^2 \theta_{23} = 0.44$, and $|\Delta m_{23}^2| = 2.4 \times 10^{-3} \text{ eV}^2$. Note that muon neutrinos produced in the atmosphere are unlikely to oscillate while reaching the surface, while the ones traversing the earth are very likely to oscillate.

2.3.3 Geomagnetic and Solar effects in the atmospheric flux

The sun emits a magnetized plasma where cosmic rays have to diffuse in order to reach the Earth's atmosphere, and as result during periods of high solar activity the low-energy portion of the CR spectrum is more suppressed. The total impact on the neutrino flux can be around 10% – 20%, and it depends on the location of the experiment. However, this is not really an uncertainty in the flux calculations, as detailed measurements have been performed to characterize this effect which is included in early estimates of neutrino fluxes [134].

On the other hand, the Earth's magnetic field has an impact on cosmic rays outside and inside the atmosphere. Inside the atmosphere it bends the charged particles produced in secondaries, while outside it acts as a filter by excluding those particles with low energies. The geomagnetic effect depends on the gyroradius of a particle, being the rigidity the relevant physical parameter that determines the impact of this effect. The geomagnetic rigidity cutoff is the minimum rigidity needed for a CR particle to be able to traverse the geomagnetic field up to a given location at earth's surface. The structure of the geomagnetic field causes variations in this cutoff, and the relevant energy range at which the effect produces an important impact depends

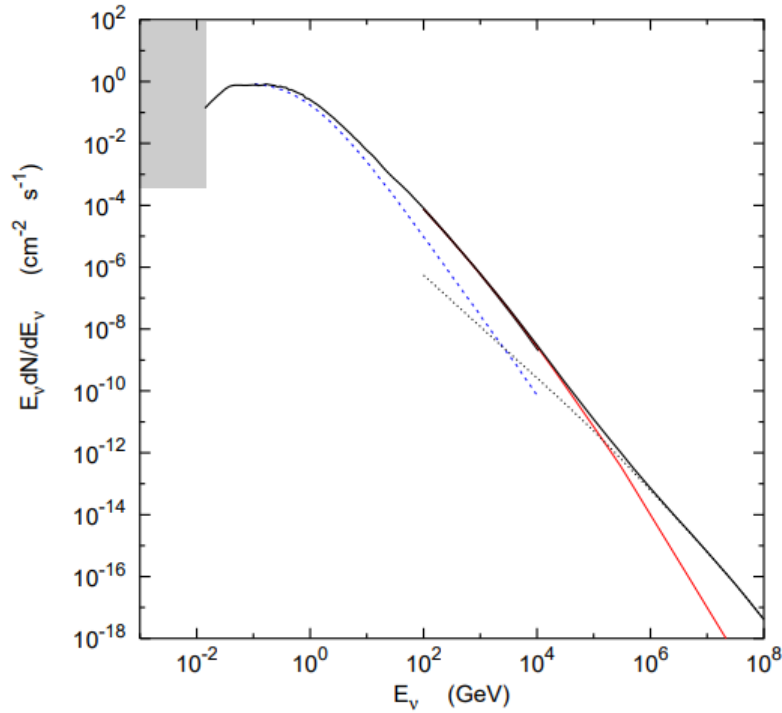


FIGURE 2.7: Neutrino spectrum: The vertical flux of $\nu_\mu + \bar{\nu}_\mu$ is depicted by solid lines, where the black curve represents the total flux, while the red curve represents only the non-prompt contribution, which is subdominant at high energies. The vertical flux of $\nu_e + \bar{\nu}_e$ is denoted by dotted lines, the blue curve is the non-prompt contribution, and the gray curve represents the prompt contribution. The gray area corresponds to the region where the solar flux dominates the spectrum. Figure taken from reference [99]

on the location and direction, in such a way that the flux must be calculated separately for each detector location. The impact of the geomagnetic effect on cosmic rays and neutrino fluxes has been measured for years, with Super-Kamiokande being one of the main collaborations that has performed multiple local measurements and simulations of this effect. A detailed atmospheric flux calculation including relevant low energy effects at Super-Kamiokande was performed by Gaisser and Honda in a seminal paper in 2002 [99].

Finally, we should note that the calculation of fluxes presents an additional challenge for determining the angular dependence at low energies, since one can no longer make the (often implicit) approximation that secondary particles go in the same direction as the primary cosmic ray. This would require a full three-dimensional calculation of the development of the shower including lateral displacement. However, in simulations involving MCEq, the lateral development of the shower is neglected as we are above the energy range in which this effect is relevant.

2.4 Detection of atmospheric muons and neutrinos

Muons have a relatively long lifetime and small interaction cross section with air nuclei, so it is not surprising that atmospheric muons are the most abundant charged

particles arriving at sea level. Moreover they are able to penetrate deeply underground, producing a relevant background for neutrino telescopes and other experiments searching for BSM particles.

A relevant quantity in the detection of atmospheric muons is their energy loss, which can be classified into continuous and discrete processes. For muons with energies below ~ 500 GeV the dominant energy loss process is due to ionization, which is almost constant in the relativistic regime. For higher energies, discrete processes dominate and include electromagnetic interaction with nuclei, bremsstrahlung and direct electron-positron pair production. All of these radiative processes are included in the total muon energy loss formula which is parameterized as [52]

$$\frac{dE_\mu}{dy} = -\alpha - \beta E_\mu, \quad (2.24)$$

where E_μ and y are the energy of the muon, and the thickness of crossed material in gr/cm^2 , respectively, while α accounts for the ionization loss. The β coefficient is the fractional energy loss by the three radiation processes mentioned above $\beta = \beta_{nuc} + \beta_{brem} + \beta_{pair}$. The thickness material y is usually given in units of meters of water equivalent (m.w.e.), with $1 \text{ m.w.e.} = 10^2 \text{ gr}/\text{cm}^2$. Moreover, since rock compositions are not the same for different underground experiments, a standard rock is defined as a common reference. This standard rock reference is characterized by an atomic mass $A = 22$, and a density of $\rho = 2.65 \text{ g}/\text{cm}^3$ [163]. Figure 2.8 displays the minimum average energy of muons at the surface E_{min} necessary to reach a certain depth $y = h$. For reference, reaching the surface of the Earth requires about 2 GeV (corresponding to 10 m.w.e), while a deep underground experiment such as Gran Sasso Laboratory requires about 2000 GeV (3800 m.w.e. or 1400 meters of rock overburden).

On the other hand, for atmospheric neutrinos the main experimental strategy of detection consists of observing high-energy interactions of these particles with a large target medium located underground, and then inferring the neutrino composition by the flavor of the outgoing lepton, which can be determined from the pattern of energy deposited within the detector. Atmospheric neutrino detectors are located underground to avoid the large atmospheric muon background, and require very large target volumes since their cross section is very small. The main detection technique for atmospheric neutrinos is by means of Cherenkov radiation [65], where the charged lepton produced in neutrino-nucleus interaction travels through the medium with a velocity that is greater than the phase velocity of the light in that substance, producing an emission of electromagnetic radiation that can be captured by photomultiplier tubes.

Neutrinos interact via the weak force, and this interaction is mediated by either charged gauge bosons $W^\pm$, or the neutral Z gauge boson. The former are denominated charged current (CC) interactions, while the latter are neutral current (NC) interactions. In CC interactions the neutrino is converted to the corresponding charged lepton as shown in the left panel of figure 2.9. The lepton carries an average of 50% to 80% of the neutrino energy [122], while the remainder is transferred to the nuclear target, producing a hadronic shower. On the other hand, in a NC interaction the neutrino is not converted into a lepton (see right panel of figure 2.9), but of course its energy/momentum changes since a fraction of its energy is transferred to the nuclear target, which is again released in the form of a hadronic shower. Since in this

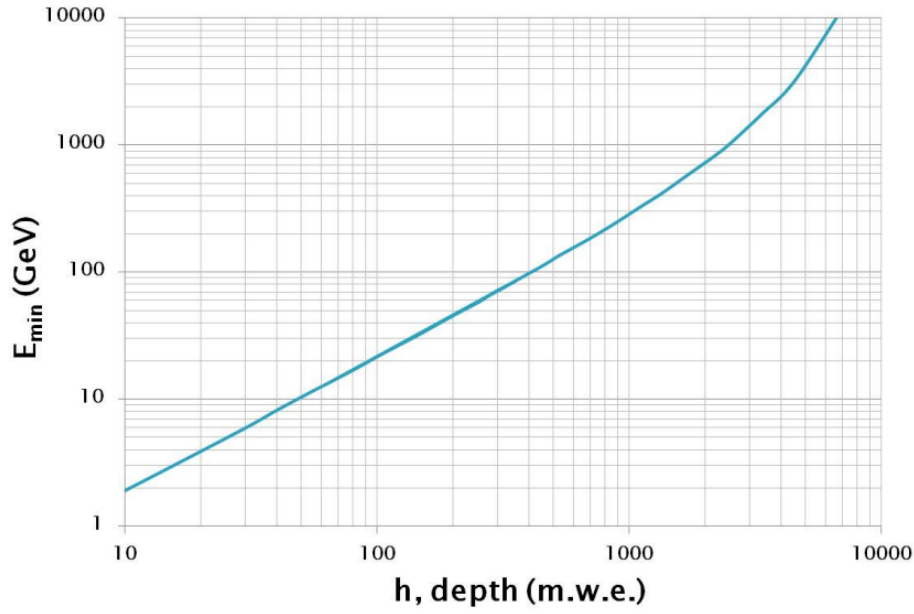


FIGURE 2.8: Minimum energy at surface E_{min} for muons to reach a given underground depth $y = h$. Figure adapted from [62].

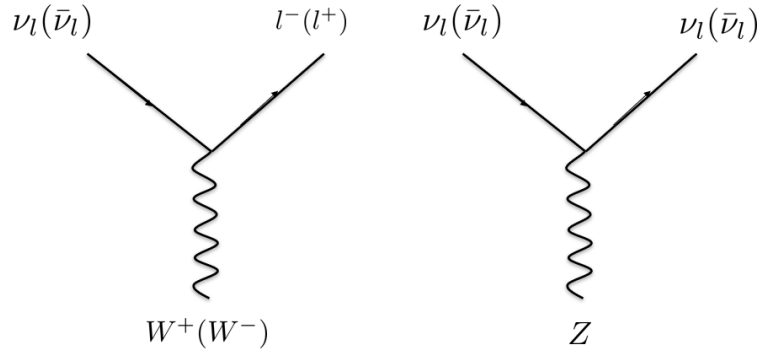


FIGURE 2.9: Left: Charged current neutrino interactions. Right: Neutral current neutrino interactions.

thesis we will also work with high energy data sets from Cherenkov detectors such as IceCube, it is worth pointing out that at high energies ($E_\nu \gtrsim 10$ GeV), neutrinos dominantly interact with hadrons through deep inelastic scattering (DIS), meaning that the neutrino scatters on the constituents of the hadrons within the atomic nucleus.

Recorded events in the detectors can be classified according to the type of measurement and the topology of the event. There are two types of measurements; contained events and partially contained events. The former are made of events completely registered where the neutrino interaction takes place in the fiducial volume and showers stop before exiting the detector. On the other hand, partially contained events are also originated in the fiducial volume, but one of the particles produced exit the detector and is detected on the veto counter. Almost all of this last kind of

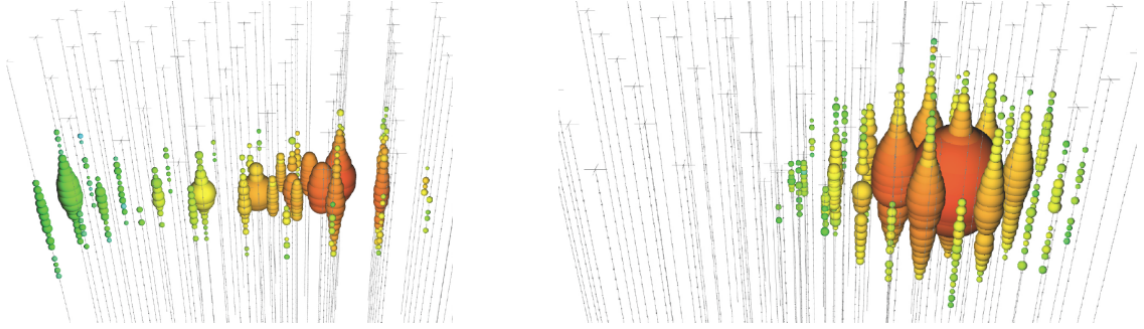


FIGURE 2.10: Left: Muon track crossing the detector. Each coloured dot represents a hit in a digital optical module (DOM). The size of the dot is proportional to the amount of light detected, and the colour represents the relative timing of light detection. Red dots represents earlier hits, while blue denotes later hits. Right: Cascade signal produced from CC interactions from either a ν_e or ν_τ , or from any flavour neutral-current interaction inside the detector. Figure adapted from [23].

events are CC interactions of ν_μ ($\bar{\nu}_\mu$) producing muons that exit the detector. Regarding the event topology, there are two main types of events, tracks and cascades. A track is produced from ν_μ charged-current interactions which generate muons. Muons have longer lifetime than the tau lepton, thus after production they propagate in a straight line and lose energy before they decay, leaving a track signature as shown in the left panel of figure 2.10. Notice that a down-going muon track which starts outside the detector is usually induced by cosmic rays and are thus a source of background. The reliable track signals are either the ones starting within the detector itself, or the ones from the up-going direction, since any up-going CR induced muon would be absorbed by the earth. On the other hand, a cascade signal is induced by either CC ν_e and ν_τ interactions, or by neutral-current interactions from neutrinos of all flavors. Any electromagnetic or hadronic shower produced in these interactions will yield this type of signal. In the case of CC ν_e interactions, the electron produced is scattered near the interaction vertex, creating an electromagnetic shower. In the case of a CC ν_τ interaction, the tau lepton produced decays almost instantaneously if it is not generated at ultra-high energies, creating a hadronic shower. Finally, when a NC interaction takes place, the products are an out-going neutrino of the same flavor and a hadronic shower. The right panel of figure 2.10 displays a typical cascade signal recorded in the IceCube experiment [7].

Now that we know how atmospheric neutrinos and muons are detected, we can move into the main chapter of the thesis in which we will discuss new physics searches of particles that could be produced in atmospheric showers. As the reader might imagine, the main background for such kind of searches comes precisely from the neutrino and muon components. In fact, it is not possible to distinguish an event from this source when looking for signals of decays of BSM particles, implying that the search strategy must be based on the observation of an excess of events on top of the expected atmospheric neutrino/muon signal. Notice that depending on the detector, it is possible to reconstruct the incoming direction of the particle, or the energy deposited in the detector, or both. Moreover, the events are typically classified according to energy range of the incident particle.

Chapter 3

Searching for new particles produced in the atmosphere

In the previous chapters, we have discussed the fundamental ingredients for the Standard Model and the description of the production of particles in cosmic-ray air showers. Now we are in a position to describe the main results of this thesis. In this chapter we will discuss the production and detection of certain types of BSM particles in atmospheric showers. We will classify these searches in two types: those involving particles that are long-lived, *i.e.* particles that once produced in the atmosphere can reach the surface, and decay within large underground detectors, and those BSM particles which are stable and can penetrate deep underground, producing a detectable signal by scattering on the target material.

3.1 Long-lived Particles

A long-lived particle (LLP) is a particle with a lifetime such that the distance it travels before decaying can be discerned by a detector. The existence of LLPs is predicted in many extensions of the Standard Model. Generically new particles that are relatively light and very weakly interacting can be long-lived. The search for LLPs in colliders and beam-dump experiments has become a field of intense activity in recent years, requiring dedicated search strategies. In this section we will focus on the atmospheric searches of long-lived heavy neutral leptons that can appear in theories that explain neutrino masses, as well as on light neutralinos that are predicted in some supersymmetric theories. The discussion will be based on the results found in the original publications [34, 69, 56].

3.1.1 Heavy Neutral Leptons

A Heavy Neutral Lepton (HNL) is a fermion with sterile neutrino quantum numbers. We will follow the notation from references [34, 69], and consider a minimal model realizing neutrino masses based on the Type-I seesaw, with at least two heavy Majorana singlets N_j , such that the Lagrangian contains the term

$$\mathcal{L}_N \supset \sum_j i\bar{N}_j \not{\partial} N_j - \left(Y_{\alpha j} \bar{\ell}_\alpha \tilde{H} N_j + \frac{m_{N_j}}{2} \bar{N}_j N_j^c \right), \quad (3.1)$$

where $\tilde{H} \equiv i\sigma_2 H^*$ is the complex conjugate of the Higgs field H , ℓ_α is the SM lepton doublet with flavor α , Y is a generic Yukawa matrix, and m_{N_j} is the Majorana mass of the singlet N_j . As we saw in section 1.3.4, after EWSB the mass eigenstates involves small mixture between active neutrinos and the HNL. As a consequence of the mixing, the HNL couples to the SM fields in the same way as active neutrinos but with a

suppressed coupling proportional to the small mixing angle¹ $U_{\alpha j} = m_{D\alpha} m_{Nj}^{-1}$, with m_D given in 1.40, and $\alpha = e, \mu$, or τ . Heavy neutral leptons can be produced in CC and NC interactions, and they can decay by the same kind of processes. When studying the production and detection of HNL, we will consider a simplified scenario where one mixing angle is assumed to be much larger than the other two, such that the model is described by only two parameters, the mass and one mixing.

The main production mechanism for HNL in atmospheric showers is similar to neutrinos, with the largest contributions to the flux coming from meson decays. Defining l as the distance traveled between the production point of the HNL to its point of entry in the detector, the HNL production profile (in a differential of distance dl), at high energies is given by a simple form of the cascade equation:

$$\frac{d\Pi_N}{dE d\Omega dl} = \sum_P \int_{E_P^{\min}}^{E_P^{\max}} dE_P \frac{1}{\gamma_P \beta_P c \tau_P} \frac{d\Phi_P}{dE_P d\Omega} \frac{dn}{dE} \quad (3.2)$$

where dn/dE stands for the differential distribution of the HNL energies in the decay, while Φ_P is the parent meson production rate. The parent mesons included in our simulations include $P = K^\pm, \pi^\pm, D^\pm, D_s^\pm$, and also the $\tau^\pm$ lepton. For each one of these possible channels one must take into account which of the couplings $|U_\alpha|$ can allow for the production. In Eq. 3.2, β_P and γ_P are the parent's boost factors, while τ_P is its lifetime in the rest frame. The differential distribution is given by

$$\frac{dn(E_P, E)}{dE} = \frac{\text{Br}(P \rightarrow N)}{\Gamma(P \rightarrow N)} \frac{d\Gamma(P \rightarrow N)}{dE}, \quad (3.3)$$

where $\text{Br}(P \rightarrow N)$ is the branching fraction of parent particles decaying to the HNL. The decay distribution term $\frac{1}{\Gamma(P \rightarrow N)} \frac{d\Gamma(P \rightarrow N)}{dE}$ must be computed independently for two-body and three-body decays. Expressions for the differential decay width of the two-body meson decays, and the three-body decay of the tau lepton can be found in references [180, 115], and were also checked in reference [34] with results presented in equations (2.5), and (4.3). Moreover, the expressions for the integral limits can be found in equations (2.7), and (2.13) of the same reference.

The different meson production rates, Φ_P were extracted from the MCEq software. This is not straightforward, but it can be done by switching off meson decay, and then correcting the rate at each point in space by including the probability of decay while propagating. More details can be found at reference [34]. The production rate of mesons is computed using a specific cosmic ray model, and a hadronic interaction model which are subject to uncertainties. We have used the SYBILL-2.3 hadronic interaction model [88], and the Hillas-Gaisser cosmic-ray model [100]. It is expected that the combined effect of the associated uncertainties is at the $\mathcal{O}(10 - 20\%)$ level, just as with atmospheric neutrinos. We will quantify this impact in the next subsection. Regarding the model of the atmosphere, we have used the NRLMSISE-00 atmospheric model [165]. Figure 3.1 depicts the production profile for HNL, over the branching fraction of production for different parent particles, including pions, kaons, D-mesons and the tau lepton. This simulation has been performed fixing $m_N = 0.1$ GeV, $h = 15.4$ km and $\cos \theta = 0.9$. The result shows that the low energy

¹We will adopt the simplified notation $|U_\alpha| \equiv U_{\alpha j}$, assuming that only one of these states is kinematically accessible.

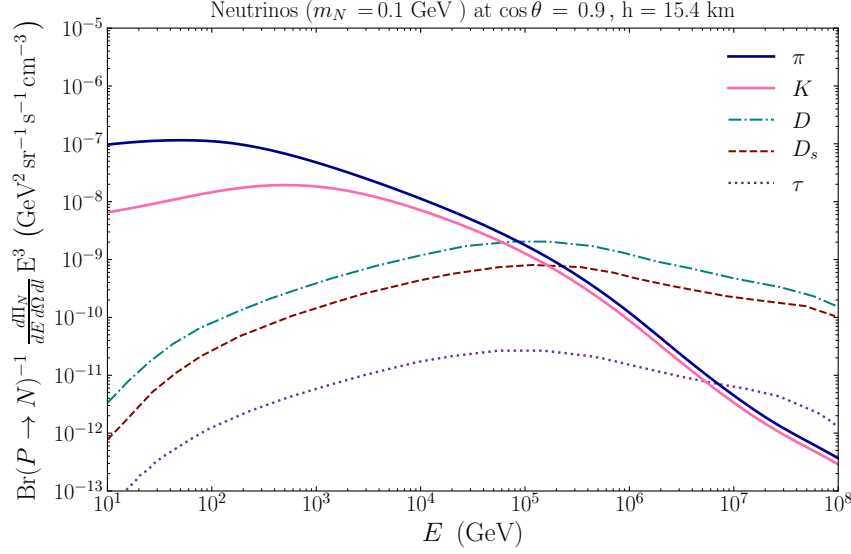


FIGURE 3.1: HNL production rate (for $m_N = 0.1$ GeV) per energy, solid angle and path length, at a height of 15.4 km and $\cos \theta = 0.9$, for different parent particles P . We consider production of π and K decays. For comparison we also show the corresponding result for two-body decays of heavier mesons (D, D_s) and leptonic (three-body) τ decays from Ref. [34]. For all lines shown here, effects due to the mass of the charged lepton produced together with the HNL have been neglected for simplicity.

part of the spectrum is dominated by the contribution from pions and kaons, just like is found for active neutrinos. At higher energies (above 10^5 GeV), the HNL production rate is dominated by D -mesons, with a contribution from tau leptons which is negligible.

From the production rate of HNL, we can estimate a signal at terrestrial experiments by considering the propagation of the HNL, and the decay probability within a detector. The first step involves computing the flux spectrum of HNL arriving to Earth, which can be accomplished by integrating the number of HNL produced at distances l , weighted by its survival probability

$$\frac{d\phi_N}{dE d\cos\theta} = \int_0^{l_{\max}} dl \frac{d\Pi_N}{dE d\cos\theta dl} e^{-\frac{l}{L_{\text{dec}}}}, \quad (3.4)$$

where $L_{\text{dec}} = \gamma_N \beta_N c \tau_N$ is the decay length of the HNL in the laboratory frame. Note that in equation 3.4 we focus on the flux per cosine of zenith angle rather than solid angle, since it is assumed that the flux is isotropic and the azimuthal component integrates trivially into a 2π factor.

The total number of HNL that can decay inside the detector, within a given time window ΔT , with energies in the interval $[E, E + dE]$ and incident direction defined by the cosine of the zenith angle between $[\cos\theta, \cos\theta + d\cos\theta]$, can be computed as

$$\frac{dN}{dE d\cos\theta} = \Delta T A_{\text{eff}}(E, \cos\theta) \frac{d\phi_N}{dE d\cos\theta}, \quad (3.5)$$

where $A_{\text{eff}}(E, \cos\theta)$ is the effective area for decay, which accounts for the probability that a decay takes place inside the detector, given a certain geometry. This effective

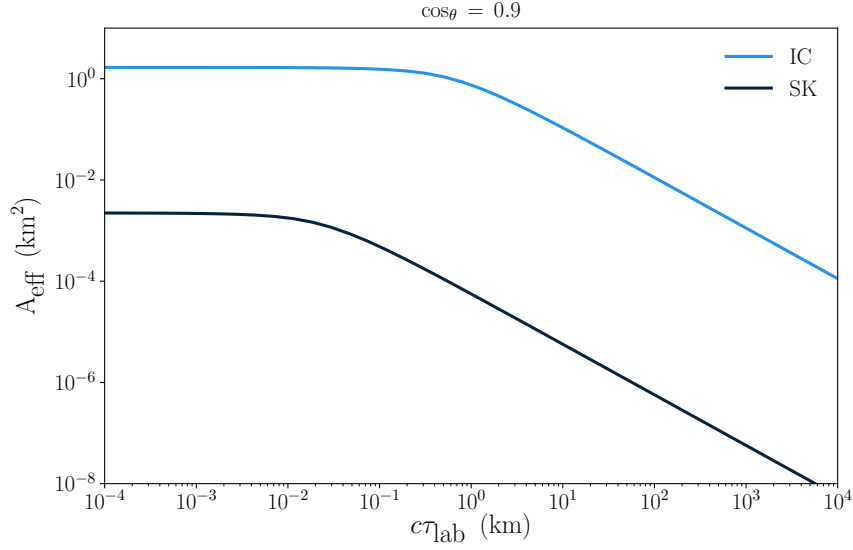


FIGURE 3.2: Effective decay area for the IceCube detector with radius ($R_{IC} = \frac{1}{\sqrt{\pi}}$ km and height $H_{IC} = 1$ km), and Super-Kamiokande ($H_{SK} = 0.04$ km and $R_{SK} = 0.02$ km), as a function of the decay length in the laboratory frame for down-going events with $\cos \theta = 0.9$.

area can be estimated integrating the surface of the detector normal to the flux direction, weighted by the decay probability of the HNL to take place inside the detector:

$$A_{\text{eff}}(E, \cos \theta) = \int dS_{\perp} \left[1 - e^{\frac{-\Delta \ell_{\text{det}}(\cos \theta)}{\ell_{\text{dec}}(E)}} \right]. \quad (3.6)$$

Here $\Delta \ell_{\text{det}}$ is the length of the segment of the HNL trajectory that cuts into the detector and $\ell_{\text{dec}}(E) \equiv c\tau_{\text{lab}}$ is the decay length in the laboratory frame. In Ref. [34] explicit expressions can be found assuming that the fiducial volume of the detector is a cylinder. In particular we consider two large underground neutrino detectors, IceCube (IC) and Super-Kamiokande (SK). For IceCube we consider a cylinder with 1.0 km height, and a radius of $\frac{1}{\sqrt{\pi}}$, which yields to a 1.0 km³ detector volume. For SK we take a height of 40 m and radius of 20 m. The effective area as a function of the decay length in the laboratory frame is shown in figure 3.2. We observe that the effective area for IC is about two orders of magnitude larger than for SK. The effective area is roughly constant up to $c\tau_{\text{lab}} \sim 1$ km in IC and about 10 m in SK and decreases with longer lifetimes, since the probability to decay within the detector falls as $\propto \ell_{\text{dec}}^{-1}$ in this regime.

We will now study under what conditions it is possible to observe an excess of HNL decay within these large volume experiments. The main idea behind is that if for a certain set of parameters involved in the production (and decay) of the HNL we predict an excess of events with enough significance that has not been observed, then these parameters can be excluded as it would imply that either IC or SK would have detected this excess on top of their expected background. Let us briefly describe the expected signal events for each of these neutrino detectors.

The IceCube Neutrino Observatory [7] is a 1.0 km³ neutrino detector at the Geographic South Pole, optimized for detecting neutrinos with energies above 100 GeV.

We consider the data sample corresponding to the analysis presented in reference [5], for which effective areas for interacting neutrino events are publicly available in the IceCube collaboration webpage [1]. The statistics corresponds to 641 days of contained atmospheric and astrophysical neutrino events. The reconstructed data samples are divided into tracks and cascade events, in the northern and southern hemispheres, and are binned in reconstructed energy. The experiment provides effective areas for each of these samples and for each particle, interaction type and true neutrino energy. From these effective areas we can infer the detector efficiency by dividing by the corresponding interaction cross sections for neutrinos on the detector. We therefore estimate the detector efficiencies as

$$\epsilon_v^{\alpha\beta}(E_{\text{rec}}^\nu, E_{\text{true}}^\nu, \theta) \simeq \frac{A_{\text{effIC}}^{\alpha\beta}(E_{\text{rec}}^\nu, E_{\text{true}}^\nu, \theta)}{\rho_{\text{ice}} V_{\text{IC}} N_A \sigma_\beta^\nu(E_{\text{true}}^\nu)} \quad (3.7)$$

where α selects the type reconstructed sample, for cascade/track, and north/south, while β selects the true interaction type. In the denominator we have the total cross section on the IceCube detector, which is estimated in terms of the nucleon-neutrino cross section, σ_β^ν , times the number of nucleons in the detector that can be estimated as the Avogadro number, N_A , times the ice's density, ρ_{ice} and the fiducial volume of the detector, V_{IC} . We assume that the HNL decay events can be measured with similar efficiencies as CC neutrino interactions, and also that the relationship between the energy of the decaying particle, and the deposited energy of all particles produced in the detector, is similar to the relationship between the true and reconstructed neutrino energies in CC neutrino interactions. This should be a good approximation for e-like events.

The number of signal events in the i -th energy bin at IC can then be obtained by

$$N_i^{\text{IC}} = \text{Br}(\text{HNL} \rightarrow e - \text{like}) \int d \cos \theta dE \epsilon_v^{\alpha\beta}(E_i, E, \theta) \mathcal{P}_{\text{pass}}(E, \theta) \frac{dN(E, \cos \theta)}{dE d \cos \theta}, \quad (3.8)$$

where we consider $0.2 \leq \cos \theta \leq 1$, corresponding to the southern sky, while $\alpha = \text{cascade}$ and $\beta = \text{CC } e\text{-like}$. We weight the signal by a branching fraction that includes the decay of the HNL to electrons, photons, τ -leptons or hadronic showers in the final state, which would be observed as e-like (or cascade-like) events. In equation 3.8, we included the passing fractions of the muon veto $\mathcal{P}_{\text{pass}}(E, \theta)$, which reduces events with a muon signal within a $3\mu\text{s}$ time window of the neutrino signal. This veto has been applied to the background and the data, so we need to include it in our calculations for consistency in the analysis. Since we focus on cascade-like events, the most appropriate passing fractions to use are those corresponding prompt ν_e , which are taken from Ref. [37]. The left panel of figure 3.3 shows the distribution of events expected in IceCube when including the signal from HNL produced from D_s mesons, that decay within the detector to cascade-like events. For this example we have fixed the mass and the lifetime of the HNL, as well as the total decay fraction of D_s mesons decaying to HNL, times the decay fraction of HNL to cascade-like events. The corresponding values are provided in the caption.

On the other hand, Super-Kamiokande is a smaller water Cherenkov detector, but it also has a lower energy threshold where the atmospheric flux is much higher. We will only consider cascade-like events, that is, events with an e-like topology as defined above. For this scenario we have used the fully-contained multi GeV e-like sample from reference [8], with reconstructed zenith angle $0 \leq \cos \theta \leq 1$. In this case, while we do have information regarding the zenith angle of the event sample, we do not have available information regarding the energy of the events. Thus,

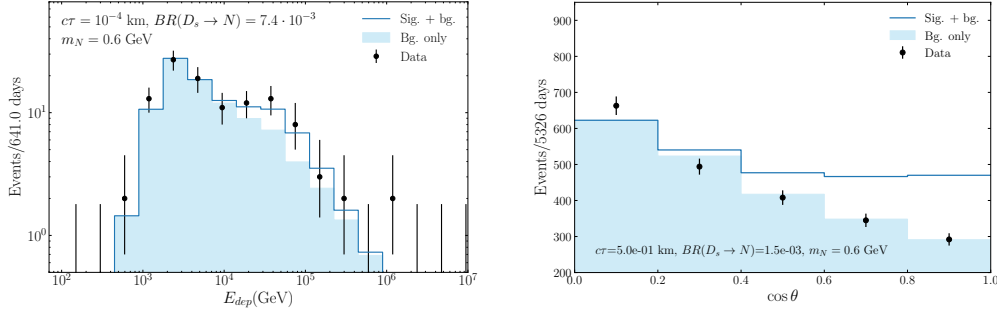


FIGURE 3.3: **Left:** Expected number of decays for the IC Medium Energy Starting Events (MESE) sample, as a function of the deposited energy in the detector and for the down-going sample only ($\cos \theta > 0$), for HNL with a mass $m_N = 0.6$ GeV, a $c\tau_N = 10^{-4}$ km and $Br(D_s \rightarrow N) Br(N \rightarrow CC \text{ e-like}) = 7.4 \cdot 10^{-3}$. The shaded histograms show the background prediction, including a fitted distribution to the astrophysical neutrino signal, the cosmic muon background, and atmospheric neutrino events. The black dots correspond to the observed data with error bars, as in Ref. [5]. **Right:** Expected number of decays in SK, as a function of the zenith angle. The shaded histograms show the atmospheric background prediction, while the solid histograms show the signal plus neutrino background. The signal in this case has been computed for a HNL with a mass $m_N = 0.6$ GeV, a $c\tau_N = 0.5$ km and $Br(D_s \rightarrow N) Br(N \rightarrow CC \text{ e-like}) = 1.5 \cdot 10^{-3}$. The data points, extracted from Ref. [8], are shown with statistical error bars for reference.

for this estimate we consider data binned only in $\cos \theta$, and integrate over all neutrino energies between the range 1.0 GeV and 90 GeV, which is the range of parent neutrino energies for multi-GeV, e-like contained events. We also assume that the detector reconstruction efficiencies for the decay of the HNL are similar to those of CC ν_e events in the multi-GeV range. Regarding detection efficiencies for Super-Kamiokande, we consider that they are flat in energy within the sample. The number of events in the i -th bin in $\cos \theta$ can be computed as:

$$N_i^{SK} = Br(\text{HNL} \rightarrow \text{e-like}) \int_{\cos \theta_i^{\min}}^{\cos \theta_i^{\max}} d \cos \theta \int_{1.0 \text{ GeV}}^{90 \text{ GeV}} dE \epsilon_{SK} \frac{dN(E, \cos \theta)}{dE d \cos \theta}, \quad (3.9)$$

where $\cos \theta_i^{\min}$ and $\cos \theta_i^{\max}$ are the lower and upper limits of the bin, while the efficiency is assumed to be $\epsilon_{SK} = 0.75$. The right panel of figure 3.3 shows the event distribution with respect to the cosine of the zenith angle that should be expected when considering a contribution from decays of HNL compatible with e-like events. For this figure we only consider the contribution of HNL produced from D_s mesons, and we have fixed the values of the mass, lifetime and total branching fraction of HNL decays, with values as indicated in the caption.

After the simulation of the expected number of events from HNL decays within the detector, we can proceed to search for the region of parameter space that can be constrained from the non observation of an excess of events over the expected background. The signal we are searching for is an excess of events, where the interaction vertex is contained in the fiducial volume of the detector. We will consider

an analysis for HNL with masses above, and below the kaon mass. The reason for adopting this strategy is that above the kaon mass the existing laboratory bounds are weaker. We consider a Poisson distribution for the events, which implies a chi-squared statistic of the form

$$\chi^2 = 2 \sum_{i,\alpha} \left(\mu_i^\alpha + B_i^\alpha - n_i^\alpha + n_i^\alpha \log \left(\frac{n_i^\alpha}{\mu_i^\alpha + B_i^\alpha} \right) \right), \quad (3.10)$$

where the sum runs over the energy (angular) bins for the IC (SK) detector analyses. We take $\alpha = \text{cascade(e-like)}$ for IC (SK). Here, n_i^α stands for the data observed in each bin, and N_i^α is the predicted number of signal events as explained above, see eqs. 3.9 and 3.8, while B_i^α is the background prediction that includes the expected number of atmospheric neutrino events². The parameters describing the lifetime, and branching fraction of the HNL are its mass m_N , and mixing U_α . In order to estimate both, the lifetime τ_N , and the total branching fraction $\text{Br}_{\text{tot}}(m_N, U_\alpha) = \text{Br}(P \rightarrow \text{HNL}) \times \text{Br}(\text{HNL} \rightarrow e - \text{like})$, it is necessary to compute the total decay width $\Gamma_N(m_N, U_\alpha)$. This will imply that the signal N_i^α is a function of these parameters, and there is a correlation between the production and decay through the mixing U_α .

For HNL with masses above the kaon mass, we find that since the atmospheric production of HNL from D-mesons and τ -lepton decays is very small, and the production and decay are very correlated, the limits that can be obtained in the $m_N - U_\alpha$ plane are very weak. On the other hand one could consider a more general scenario where both production and decay are independent, for example the HNL might be coupled to a dark sector that could increase their lifetime independently of U_α . We can then consider the limits as a function of the total branching fraction Br_{tot} , and lifetime τ_N as the free parameters that can be constrained. We consider that the data is chi-squared distributed with n degrees of freedom. Our exclusion limits indicate the region in parameter space where the χ^2 value exceeds the corresponding one at 90% CL. Our results are shown in Fig. 3.4 for the HNL, assuming $m_N = 1.0$ GeV. The thick lines correspond to the IC limits, while the thin lines indicates the SK result. For both IC and SK, the best limits are obtained when the decay length in the laboratory frame is $c\tau_{\text{lab}} \sim \mathcal{O}(10)$ km. Moreover, the two experiments observe events in different energy regimes: $\mathcal{O}(\text{TeV})$ for the case of IC, and $\mathcal{O}(\text{GeV})$ in the case of SK. Therefore the boost factor (which also depends on the mass of the HNL) will be very different. As a result, their sensitivities are complementary, and can probe a different range of $c\tau$. For IC, the best limit is reached for $c\tau \sim \mathcal{O}(10^{-3} - 10^{-2})$ km in the case of a HNL with mass $m \sim 0.5 - 1$ GeV, while for SK the best limit is obtained for $c\tau \sim \mathcal{O}(1 - 10)$ km.

On the other hand, we find that in the case of IC, the minimization of the χ^2 on these two-dimensional planes gives a minimum of the χ^2 that lies at a non-zero value of the HNL production rate. This is the result from a small excess in the data at around 30 TeV from an unknown source (see the left panel of figure 3.3). Figure 3.5 displays the 1σ regions for 2 degrees of freedom, where χ_{min}^2 is the global minimum, for a HNL produced from D, D_s and τ decays in the atmosphere.

Next we will discuss the results that can be obtained for HNL produced from pion and kaon decays in the atmosphere. The production rate of HNL from pions and kaons is orders of magnitude above the production from D-mesons. In this case we

²In the case of IC, we also consider the contribution from the fitted astrophysical neutrino flux.

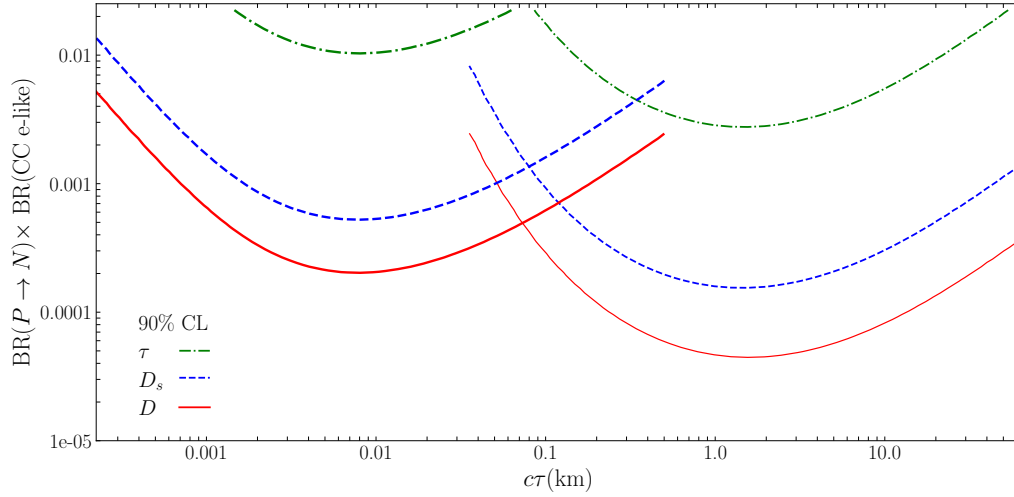


FIGURE 3.4: Limits on HNL with mass $m_N = 1.0$ GeV from IceCube (thick lines) and SuperKamiokande (thin lines) on the $BR(P \rightarrow N) \times BR(N \rightarrow CC\text{-}e \text{ like})$ vs $c\tau$ plane including production from the parent particles $P = D$ (solid), D_s (dashed) and τ (dash-dotted).

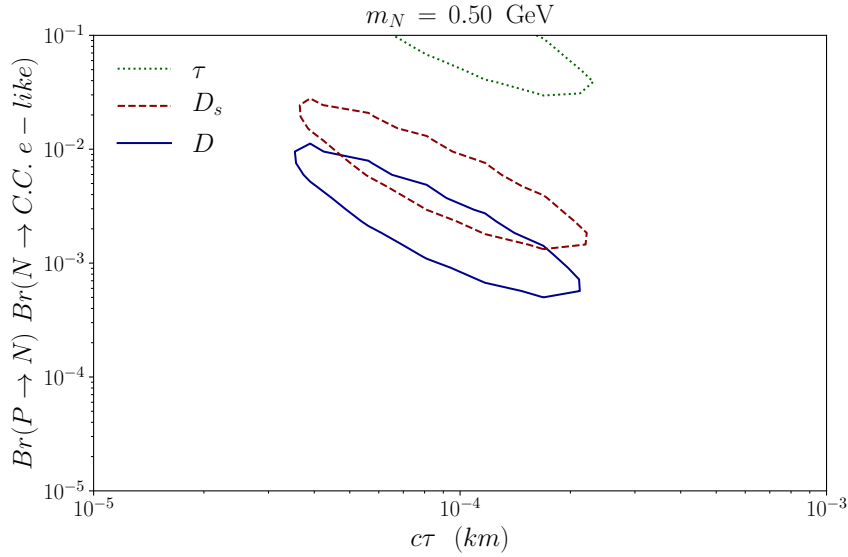


FIGURE 3.5: 1σ best-fit regions for HNL decays in IceCube assuming uncorrelated $Br(P \rightarrow N)$ vs $c\tau$ for $m_N = 0.5$ GeV, and considering separately the production from three parent particles $P = D, D_s, \tau$.

perform an analysis with the correlation, and explicit dependence of production and decay through the parameters m_N, U_α . Since we already know from the result found in figure 3.4, that the sensitivity for IceCube is worse than Super-Kamiokande, this analysis will focus in the data set of the later. We also notice that in this case we use both the μ - and e-like fully-contained multi-GeV events, adding the single- and multi-ring samples. We consider the same chi-squared test from equation 3.10, and take only one non-vanishing mixing U_α at a time. Our results are shown in Fig. 3.6 for $|U_e|^2$, and for $|U_\mu|^2$ in Fig. 3.7, at 90% confidence level (C.L.). In the case of $|U_e|^2$ the contribution from $\pi^\pm$ decays is clearly dominant for $m_N < 140$ MeV, as can be

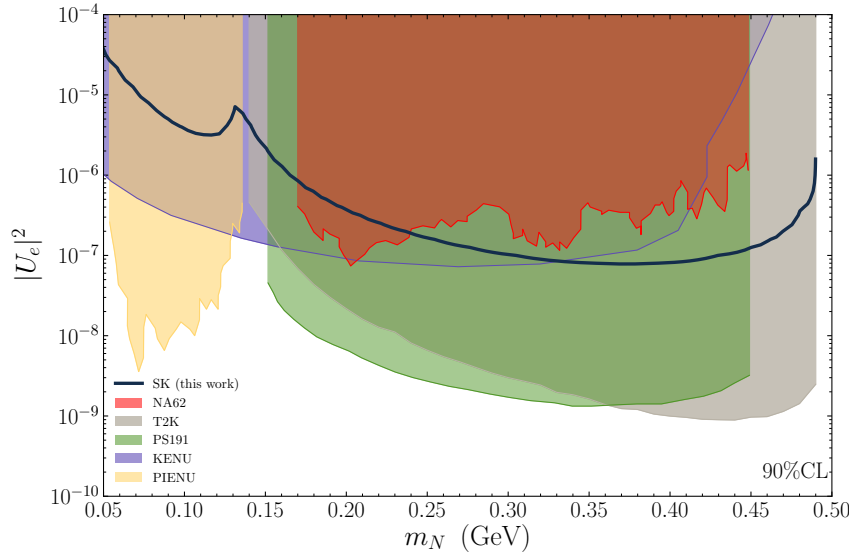


FIGURE 3.6: SK constraints on the minimal HNL scenario at 90% C.L., projected onto the plane $|U_e|^2$ vs m_N (for $U_{\mu,\tau} = 0$). Our results (solid black lines) are compared to corresponding limits obtained for the T2K near detector [9], NA62 [70], PS191 [50, 49], and PIENU [157]. The line labeled as KENU was derived in Ref. [53] using precision measurements of leptonic decay channels of the kaon.

seen from the peak in sensitivity at around 0.1 GeV. We find that the limits derived from our simplified analysis are already able to set competitive constraints on the mixing of HNL with the electron, and muon neutrino sectors. Our limits are also compared with those previously obtained in accelerator-based experiments: from displaced decay searches at PS191 [50, 49], and at the T2K near detector [9], as well as from peak searches in E949 [39], PIENU [157], and NA62 [70]. We also show the bound derived in reference [53] from measurements of kaon decays into electrons or muons [145]. In the case of $|U_e|^2$, the limits obtained from SK are comparable, or even better than analogous limits from peak searches, while they are not competitive with those from displaced decay searches. In the case of $|U_\mu|^2$, peak searches in E949 also yield better constraints than our limits from SK data.

3.1.2 Neutralinos

Neutralinos are electrically neutral fermions which arise in the framework of supersymmetric models. The MSSM is the simplest supersymmetric extension of the SM that is phenomenologically viable, preventing proton decay through a discrete Z_2 symmetry called R-parity. In this model, there are four neutralinos, with the lightest one being the lightest supersymmetric particle (LSP) of the model, which provides a natural candidate for fermionic dark matter because of its stability and lack of electromagnetic interactions. However, it is possible to extend this model, including R-parity violating (RPV) interactions, respecting the bounds on the proton lifetime, as long as baryon number or lepton number are preserved. In such a case, the neutralino LSP $\tilde{\chi}_1^0$ is no longer stable and can decay into SM particles. The R-parity violating couplings are small and can make the decay length of the neutralino macroscopic.

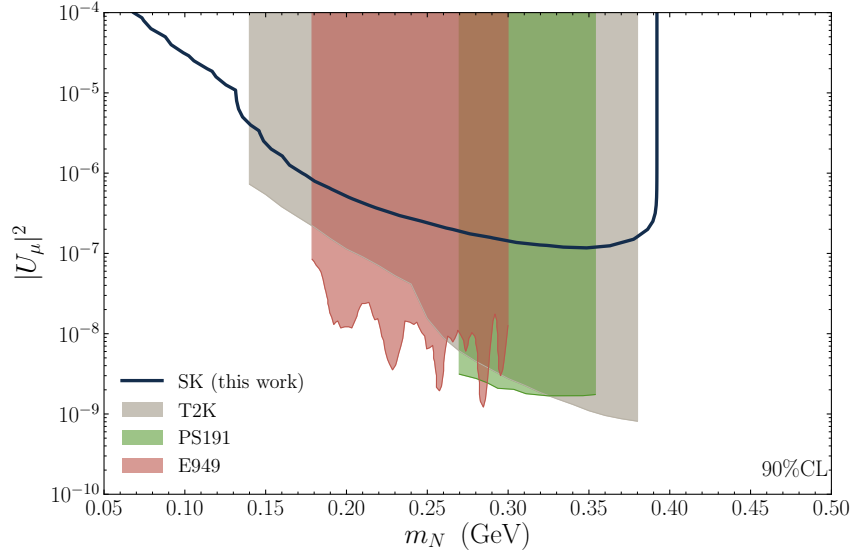


FIGURE 3.7: SK constraints on the minimal HNL scenario at 90% C.L., projected onto the plane $|U_\mu|^2$ vs m_N (for $U_{e\tau} = 0$). Our result (solid black line) are compared to corresponding limits obtained for the T2K near detector [9], E949 [39], and PS191 [50, 49].

The search for these long-lived neutralinos has been considered before in collider experiments, as well as future detectors in the vicinity of the LHC such as SHiP. However, we have studied for the first time the possibility of detecting terrestrial signals at underground neutrino detectors, which could be generated from the decay of neutralinos produced in our atmosphere.

Without imposing the R-parity symmetry, the most general Lagrangian that respects gauge and space-time symmetries in the MSSM contains the term [188]

$$\mathcal{L} \supset \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k^c, \quad (3.11)$$

where the hatted symbols denote gauge multiplets of superfields which includes leptons ℓ_i , and sleptons $\tilde{\ell}_i$ in $\hat{L}_i$, left-handed quarks q_{jL} , and squarks $\tilde{q}_{jL}$ in $\hat{Q}_j$, and right-handed antiquarks $\bar{d}_{kR}$, and antisquarks $\tilde{d}_{kR}^*$ in $\hat{D}_k^c$. The family indices i, j and k run from 1 to 3, leading to 27 independent trilinear couplings parameters λ'_{ijk} . The semileptonic operators implies unstable neutralinos that can be produced in atmospheric showers from the decay of mesons. The specific neutralino couplings included in equation can be found in equations (21)-(23) of reference [185]. The relevant Feynman diagrams involved in the production of neutralinos from charged meson decays that are relevant for our study can be seen in figure 3.8. We focus on the trilinear couplings λ'_{121} , and λ'_{112} , which allows to produce neutralinos from the decays of D -mesons and kaons. These decays are accompanied always with an electron, as detailed in the two benchmark scenarios depicted in table 3.1. The inclusion of RPV parameters also allows neutralino two-body decays into a meson and a lepton. In our first benchmark scenario, we study neutralinos with masses in the range from the mass of the kaon to the mass of the D -meson, while in the second benchmark case we restrict the neutralino mass to the interval between the mass of the pion, and the mass of the kaon. Therefore, for the second benchmark scenario,

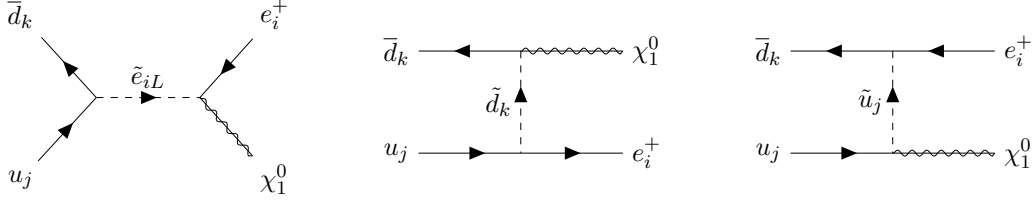


FIGURE 3.8: Neutralino production diagrams. We display the specific case where production occurs via the decay of a charged meson M_{kj}^+ composed of quarks d_k^C and u_j . For the cases of our interest (see table 3.1), we set the flavour indices of the initial states in the diagram such that M_{kj}^+ corresponds to D^+ ($k = 1, j = 2$) for benchmark B1 and K^+ ($k = 2, j = 1$) for benchmark B2.

we include the trilinear coupling λ'_{111} , which is necessary to make the neutralino decay when its mass is lower than the kaon mass.

Similarly as with the case of neutrinos and HNLs, we simulate the production of light neutralinos in the shower by solving the cascade equation involving only source terms as in equation 3.2, considering D-mesons and kaons as parent particles. To estimate the decay rate of these mesons to neutralinos, we follow reference [185]. Moreover, we consider different hadronic interaction models for the production of mesons that are updated with LHC data. In particular, we focus on the SYBILL-2.3 [88], QGSJET-II-04 [161], EPOS-LHC [166], and DPMJET-III [175] models. These models might yield non-negligible differences in the production rate of mesons, and can be a relevant source of uncertainty in our simulations³. As an illustration of our numerical simulations, we show in figure 3.9 the production rate for a 0.31 GeV neutralino, at a height h of 15.51 km, and with a vertical direction of 25.84° degrees from the zenith, produced in the decay of both D -mesons and kaons. The spread of the orange band in the plot covers the most pessimistic and optimistic cases given the uncertainties associated with the choice of a given hadronic interaction model. We were not able to assess the uncertainty for the simulation of neutralinos produced from D -mesons, since currently the only updated hadronic interaction model for charmed hadrons is SYBILL-2.3 [88, 6, 25, 11]. We will illustrate the impact of different hadronic interaction models in the final results. We will provide the excluded region of parameter space that can be probed using the same SK atmospheric neutrino data from reference [8].

Once produced, the neutralinos that reach the Super-Kamiokande detector can decay to SM particles inside its instrumental volume. Among the possible decay products of the neutralinos there are charged leptons and mesons, as well as neutral pseudoscalar and vector mesons. The different decay products can correspond to different kinds of signals in the detector; showering (or e -like), and non-showering (or μ -like) events. We focus on the showering (e -like) signals originated from the decay of neutralinos within two benchmark scenarios summarized in table 3.1. Although for the production of neutralinos from D -mesons and kaons we only need

³Another relevant input in this calculation is the cosmic-ray model chosen for the primary spectra. We checked that the uncertainties from this choice are sub-leading in the energy regime considered in this work. We use the Hillas-Gaisser cosmic-ray model H3a [101] in our calculations.

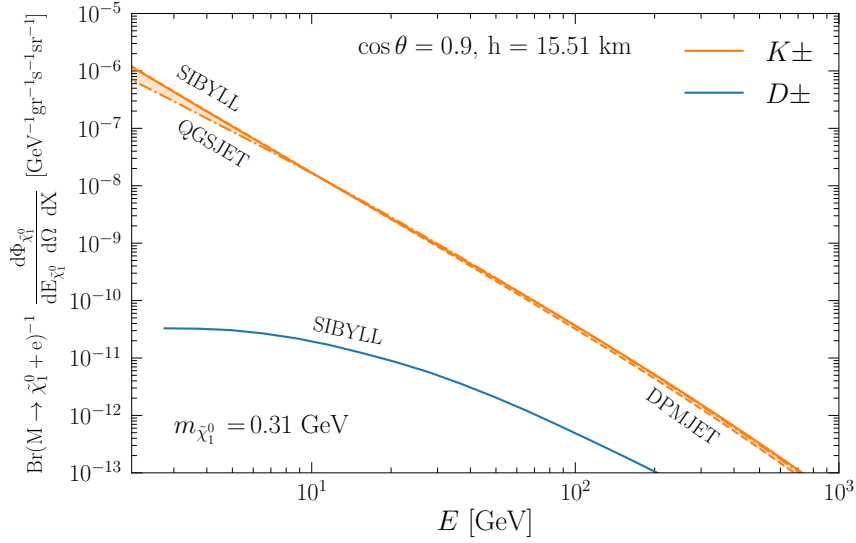


FIGURE 3.9: *Neutralino Production Rate.* Energy spectrum of the atmospheric production rate of light neutralinos from charged D -mesons (blue curve), and kaons (orange band). The spread of the orange band reflects the difference in the production originated from the election of different hadronic event generators; SIBYLL (solid line), QGSJET (dot-dashed line), and DPMJET (dashed line). We choose a representative mass for the neutralino with a value of 0.31 GeV. The angular direction is fixed at $\cos \theta = 0.9$, while the height was chosen around 15 km, where a maximum production rate is expected.

	RPV coupling	Production	Decay mode
B1	$\lambda'_{121}, \lambda'_{112}$	$D^\pm \xrightarrow{\lambda'_{121}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K_S^0 + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K^{*0} + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K^{(*)+} + e^-$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K_S^0 + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K^{*0} + \nu_e$
B2	$\lambda'_{112}, \lambda'_{111}$	$K^\pm \xrightarrow{\lambda'_{112}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^+ + e^-$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^0 + \nu_e$

TABLE 3.1: Parameter choices that define our benchmark scenarios B1 and B2. The CP conjugate processes of all the decays shown are also allowed and therefore contribute to the neutralino decay width. The decay modes displayed in this table are compatible with a shower (or e -like) signal in Super-Kamiokande. See text for details.

the trilinear couplings λ'_{121} , and λ'_{112} , the decay of neutralinos produced from kaons into visible showers in SK proceeds via the coupling λ'_{111} . For benchmark B1 we consider neutralino masses in the range $m_{K^+} + m_e \leq m_{\tilde{\chi}_1^0} \leq m_{D^+} - m_e$, while for scenario B2 the neutralino mass lies in the range $m_{\pi^0} + m_e \leq m_{\tilde{\chi}_1^0} \leq m_{K^+} - m_e$. The choice of B1 allows for a direct comparison with references [185, 75], where the sensitivity reach for future experiments aiming to explore the lifetime frontier was estimated only in the context of collider searches. Moreover, the B2 scenario allows us to study the production channel from kaon decays that has not been considered before, and that can give rise to an abundant production of neutralinos in atmospheric showers.

For the case of neutral mesons in the final states, we assume that they decay to a showering signal inside the volume of the detector. The parameter choice of the benchmark scenario B1 also allows K_L^0 as a final state for neutralino decays, but we do not consider this as part of our signal since this particle will typically decay outside the detector due to its large decay length.

The event distribution can be computed using the formula from equation 3.5, and with this it is possible to estimate the number of events using equation 3.9. The signal computed depends on the neutralino mass and the RPV couplings through its lifetime, branching fraction of production from mesons, and the fraction of neutralinos that decay to an e -like signal in the detector. Once again, we adopt a Poisson distribution for the signal events, and perform a chi-squared test using equation 3.10. We consider as background the corresponding e -like signal from atmospheric neutrinos provided by reference [8]. The results for the two benchmark scenarios are displayed in figures 3.10, and 3.11. Notice that it is possible to combine the dependence of the decay widths on the RPV parameters, and the sfermion mass in the ratio $\lambda'_{ijk}/m_{\tilde{f}}^2$, so we adopt this combination as free parameter when presenting the results. We use the decay width of reference [185] for both, meson and neutralino decay (see equation (38)-(41) of this reference).

For the benchmark scenario B1, the strongest bound found in the literature comes from meson oscillation observables [57]. The tree-level contribution to kaon oscillations induced by the choice of parameters in B1 implies the sneutrino mass-dependent bound [80]

$$|\lambda'_{112}\lambda'_{121}| \leq 2.2 \times 10^{-8} \left(\frac{m_{\tilde{\nu}_e}}{1 \text{ TeV}} \right)^2, \quad (3.12)$$

where $m_{\tilde{\nu}_e}$ is the sneutrino mass. In our results for B1, we evaluated the limits for a sneutrino and squark mass of 1 TeV. As it can be seen on the left panel, the limits imposed by kaon oscillations for degenerated sneutrino and squark masses exclude all the parameter space within the projected sensitivity reach at SHiP [185], our resulting constraints from SK, and the limit from the Drell-Yan (DY) monolepton process $pp \rightarrow \ell \bar{\nu}$ at the LHC (labeled as ‘Colliders’)[43].

On the other hand, since in Super-Kamiokande all the neutralino decay products listed in table 3.1 are visible as e -like events, we could still observe a signal when the parameter λ'_{112} is set to zero, and we are left with neutral kaons in the final states. In such a case, the limits from kaon oscillations do not apply, as can be seen in figure 3.10 right panel. For this case we are left only with the collider constraint from the DY dilepton process $pp \rightarrow \ell^+ \ell^-$ for the parameter λ'_{121} . In this last scenario we find that the limit obtained from Super-Kamiokande data is stronger than the current constraint in the corresponding mass range.

On the other hand, in the case of the scenario B2, while there is a stringent limit from

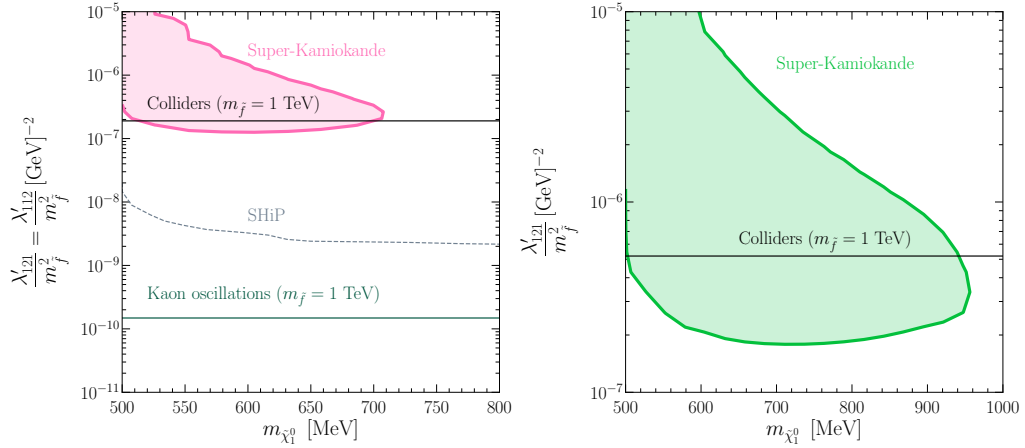


FIGURE 3.10: **Limits for benchmark B1.** Limits derived at 90% confidence level in the parameter space defined by the neutralino mass $m_{\tilde{\chi}_1^0}$ and the trilinear RPV parameters $\lambda'_{121}, \lambda'_{112}$. On the left panel, the pink contour indicates the parameter space excluded by SK when the two RPV couplings are set to the same value, while the dashed line shows the projected SHiP sensitivity reach [185]. On the right panel, the green contour indicates the excluded parameter space when $\lambda'_{121} \neq 0$ and $\lambda'_{112} = 0$. Solid lines indicate constraints from neutral kaon oscillations [80] in green and DY processes at the LHC (labeled as ‘Colliders’) [43] in black, both evaluated at a sfermion mass of 1 TeV (for details, see text).

kaon oscillations on the product of the parameters [80], the most stringent constraint comes from $0\nu\beta\beta$. The current limit for the half life of this process is determined by the KamLAND-Zen experiment [103], which imposes the upper limit [74]

$$\lambda'_{111} \leq 2.2 \times 10^{-3} \left(\frac{m_{\tilde{q}}}{1 \text{ TeV}} \right)^2 \left(\frac{m_{\tilde{g}}}{1 \text{ TeV}} \right)^{1/2}, \quad (3.13)$$

where $m_{\tilde{q}}$ and $m_{\tilde{g}}$ are the masses of the squarks and gluinos, respectively. Due to the higher kaon flux with respect to D -mesons, the limits from Super-Kamiokande on B2 are more stringent than the bounds from both, kaon oscillations and $0\nu\beta\beta$, as shown in Figure 3.11. On the left panel we present the bound on for $\lambda'_{112} = \lambda'_{111}$. The excluded parameter space is roughly independent of the choice of hadronic interaction model. Furthermore, our limit does not depend on the value of the sfermion masses. However, the bounds in equations (3.12–3.13) become more strict as sfermion masses increase, and become less strict when the mass is lower. On the right panel, we present the limit leaving both couplings free.

The excluded combinations of values for the neutralino masses and RPV couplings allow us to determine the region of the parameter space that can be probed in the plane defined by the total branching fraction, which is the production branching ratio of the neutralino, times the branching ratio of its decay to a visible signal, versus the lifetime of the neutralino. The results can be seen on the right panel of figure 3.12. Given that for benchmark B1 we are considering the same production and decay channels as in references [185] and [75], we can compare the current region that can be excluded by our Super-Kamiokande search, with the sensitivity reach

expected for FASER [91], and other possible future experiments including CODEX-b [108], MATHUSLA [67] and SHiP [27] (left panel of figure 3.12). This is achieved by translating the bounds on the fundamental parameters from these previous studies, into the plane of the total branching fraction and the lifetime of the neutralino. In the case where the neutralinos are produced in kaon decays, there are no studies on the expected sensitivity reach at the next generation of detectors. The fact that the limit for neutralinos produced from kaon decays is much better than for neutralinos produced from D-meson decays, is mostly due to the difference in the atmospheric flux of the parent meson.

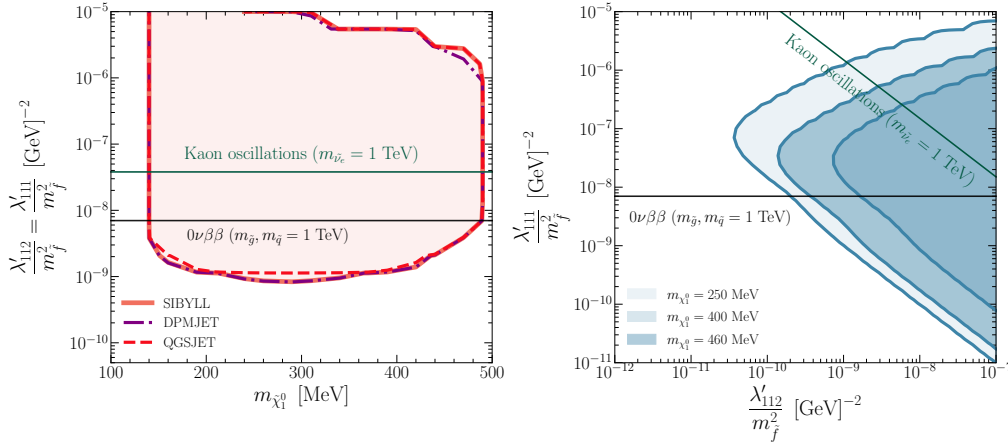


FIGURE 3.11: **Limits for benchmark B2.** Limits derived at 90% confidence level in the parameter space defined by the neutralino mass $m_{\tilde{\chi}_1^0}$ and the trilinear RPV parameters λ'_{112} , λ'_{111} . On the left panel, the red contours indicate the parameter space excluded by SK when the two RPV parameters included in B2 have the same value. The different types of lines are used to distinguish the contours obtained with three different event generators, SIBYLL-2.3, QGSJET-II-04, and DPMJET-III. On the right panel, the SIBYLL-2.3 hadronic interaction model was used to generate three blue contours that indicate the parameter space excluded by SK when both RPV parameters vary freely. The different shades of blue represent different values of the neutralino mass. Solid lines indicate constraints from kaon oscillations [80] in green and neutrinoless double beta decay ($0\nu\beta\beta$) [74] in black. The bounds are evaluated for degenerated sfermion masses set to a value of 1 TeV.

3.2 Stable Particles

In this section we will discuss searches for BSM particles that can be produced in atmospheric showers, and which are stable and do not decay after being produced in the sky. Two different scenarios will be considered; one model in which particles with a tiny electric charge can be produced directly from the decay of mesons, and another model with hadrophilic dark matter that could be produced from the decay of a scalar coupled to the Standard Model sector via up-type quarks. In both scenarios we will find that these particles can leave a signal at underground experiments, which can yield to competitive constraints in their respective parameter space. The discussion will be based on the results found in the original publications [38, 36].

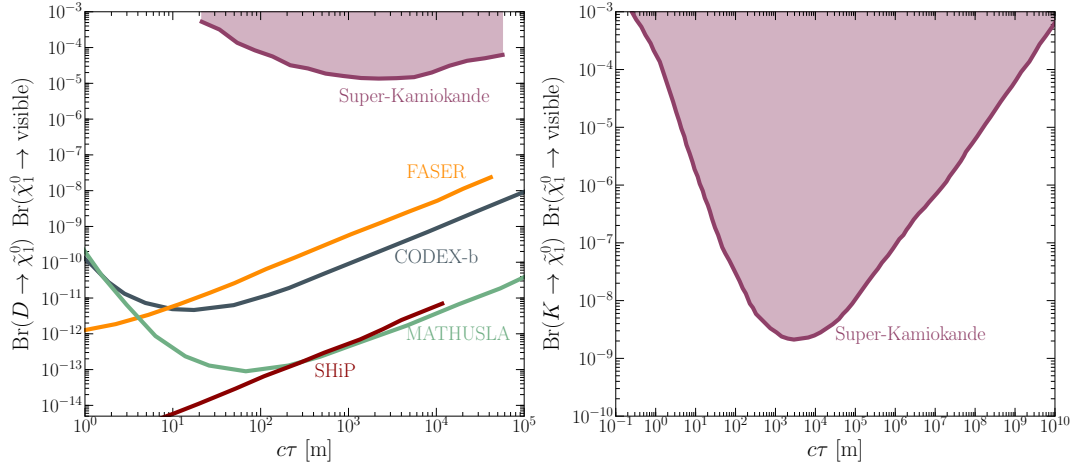


FIGURE 3.12: *Comparison with lifetime frontier experiments.* Limits at 90% confidence level in the plane defined by the total branching ratio and the proper decay length of the long-lived neutralino. On the left panel, the purple contour shows the limits for SK derived in this thesis for benchmark B1, and the coloured lines displays the projected sensitivity reach at future experiments for the same benchmark scenario derived in references [185, 75]. On the right panel, the limit obtained using Super-Kamiokande data for the benchmark scenario B2 is shown. There are no other studies for this case in the literature.

We note that while for both of these models there are previous investigations related to atmospheric searches, we consider state-of-the-art simulations and broaden the scope of the searches by including multiple detectors.

3.2.1 Millicharged Particles

As mentioned in the introduction, a millicharged particle or MCP is a theorized particle with a small electric charge $Q = \epsilon e$, with $\epsilon \ll 1.0$ a small number. We will explore the atmospheric constraints that can be placed on the charge fraction. We will make the standard phenomenological assumption regarding the nature of the millicharged particle by considering that it has some mass, and small coupling to the Standard Model photon via a small electric charge. MCPs can be produced from the decay of mesons generated in atmospheric showers. We include all the relevant channels: pseudoscalar (π^0, η) and vector mesons (ρ, ω, ϕ and J/Ψ)⁴.

In principle, heavier mesons such as Y and direct, Drell-Yan production of MCP could also allow for searches for MCP particles above $m_{J/\Psi}/2$. However, it is expected that the production rate of these mesons is limited, so we do not expect a good sensitivity from these production mechanisms.

The production of MCPs through meson involves two-body and three-body decays, with explicit formulas for the decay rates and branching fractions given in references [34, 38]. We simulate the production of MCPs by solving the cascade equation containing only source terms from the decay of mesons like in equation 3.2. For our benchmark results we have used the SYBILL-2.3 hadronic interaction model and the Hillas-Gaisser cosmic-ray model *H3a*. Moreover, since millicharged particles are stable, a direct integration of the production rate in the column depth X can be performed to find the differential energy spectra for a fixed zenith angle θ . In Fig. 3.13,

⁴The η' contribution to the MCP flux can be shown to be suppressed and has been neglected.

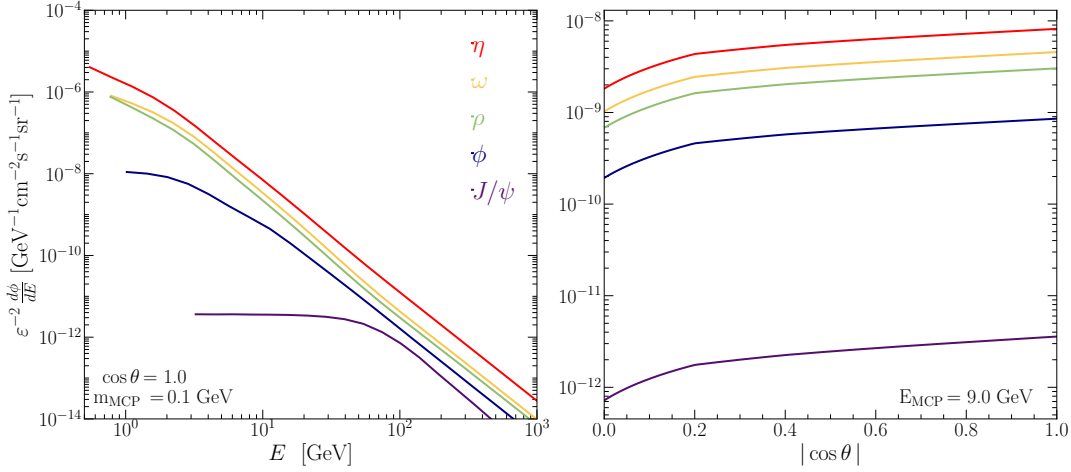


FIGURE 3.13: *Surface flux of millicharged particles from different parent mesons.* Left: Spectra, which scales as ε^2 , of millicharged particles arriving at the surface of the Earth with zenith angle 0° , for each one of the different parent mesons. Right: Angular dependence of the flux at 9 GeV. In both panels the MCP mass is set to 100 MeV.

we show the expected energy and angular dependence of the MCP flux at the surface of the Earth, for each one of the parent mesons considered in this work. We consider a 100 MeV MCP to illustrate our results, and we do not take into account MCPs with energies higher than 1 TeV, as the flux drops significantly. Uncertainties in the flux were evaluated by considering the impact of selecting different cosmic-ray models, and also different hadronic interaction models. We found that the largest source of uncertainty is in the model dependence of the hadronic interaction. The impact tends to be higher at higher energies, and depending on the energy range, and parent meson considered, it could reach over $\sim 60\%$.

We will show that underground neutrino detectors such as JUNO, and Super-Kamiokande, provide a good opportunity to search for MCPs produced in atmospheric showers. These experiments use the Earth's crust to shield from atmospheric muons, which usually constitutes the main source of background. Thus, the detection probability of MCPs depends on the flux arriving at the detector, which requires a simulation of the energy loss as the MCP propagates through the earth. It can lose energy by ionizing the medium or by scattering on nuclei. The average energy loss along the MCP trajectory can be parametrized by the same formula given for muons in equation 2.24, with the caveat that the electromagnetic interactions is suppressed by the charge fraction. We assume that the energy loss formula scales with ε^2 , and neglect any mass-dependence of the loss terms.

The probability that a MCP with an energy E_i at the surface, arrives at a detector located at depth d , with an energy E_f , depends on the coupling ε and the incident direction $\cos \theta$ as

$$\mathcal{P} = \exp \left(- \frac{D(\cos \theta)}{R(E_i, \varepsilon, E_f)} \right), \quad (3.14)$$

where the overburden length at depth d is given by

$$D = \sqrt{(R_\oplus - d)^2 \cos^2 \theta + d(2R_\oplus - d)} - (R_\oplus - d) \cos \theta, \quad (3.15)$$

where $R_\oplus$ is the radius of the earth. The average distance range $R(E_i, \epsilon, E_f)$ is given by

$$R = \frac{1}{\epsilon^2 b} \ln \left(\frac{1 + \frac{a}{b} E_i}{1 + \frac{a}{b} E_f} \right). \quad (3.16)$$

with $a = 0.22 \text{ GeV/mwe}$, and $b = 4.64 \times 10^{-4} \text{ mwe}$, being the energy loss parameters, which we take from standard muon measurements [143, 118].

The flux at the detector can be obtained by a convolution of the flux at the surface with the survival probability. Our simulation shows that for a relatively large charge fraction $\epsilon^2 \sim 10^{-2}$, the upward-going flux, $-1 \leq \cos \theta \leq 0$ is highly attenuated. On the other hand, for smaller charge fraction $\epsilon^2 \sim 10^{-5}$, the deflection and absorption of the MCPs is negligible.

Once the MCPs arrives at the detector they can interact with the medium by scattering off electrons. The differential cross section between MCPs and electrons is given by [152]

$$\frac{d\sigma}{dE_r} = \epsilon^2 \alpha_{EM}^2 \pi \frac{m_e(E_r^2 + 2E_\chi^2) - E_r(m_e(2E_\chi + m_e) + m_\chi^2)}{E_r^2 m_e^2 (E_\chi^2 - m_\chi^2)}, \quad (3.17)$$

where E_χ is the energy of the incoming MCP, E_r is the recoil energy, while α_{EM} is the electromagnetic coupling constant, and m_χ , and m_e are the masses of the MCP and the electron, respectively. It is possible to verify that in the ultrarelativistic limit $E_\chi \gg E_r, m_\chi, m_e$, the differential cross section scales like E_r^{-2} . A lower energy threshold therefore implies a larger signal.

The event rate from MCPs can be obtained by a convolution of the millicharged particle flux with the differential cross section, multiplied by the number of targets, and the detection efficiency of the experiment

$$\frac{dN}{dE_r d\cos\theta} = n_e \epsilon(E_r) \int dE_\chi \frac{d\phi}{dE_\chi d\cos\theta} \frac{d\sigma}{dE_r}, \quad (3.18)$$

where n_e is the total number of electrons in the detector, and $\epsilon(E_r)$ is the reconstruction efficiency of electron recoil events.

We will search for recoil electron signals in three different underground experiments: Super-Kamiokande, JUNO [78] and XENON1T [33]. The first two are underground neutrino detectors, while the later is a dark matter detector. It is interesting to compare the result from both kind of experiments, as neutrino detectors have larger target volume than XENON1T, yet the latter has a much lower energy threshold where the signal peaks. Moreover, we will divide the search strategy into two different types of analyses, those relying on single-hit signals, and those that rely on multiple energy depositions from a single MCP particle traversing the detector. We will show that the former is advantageous for large-volume, high-energy-threshold experiments (*e.g.*, water Cherenkov detectors like Super-Kamiokande), whereas the latter is strongest in scintillator experiments with low-energy thresholds like JUNO.

Next, we will discuss the different sources of background events that contribute to MCP searches in water Cherenkov, and liquid scintillator detectors.

Penetrating muons: Our atmosphere is flooded with muons generated from the decays of charged mesons that are dominated by pions and kaons at low energies. The most efficient way to suppress this, and other possible cosmogenic backgrounds is to have a sufficient amount of overburden over the location of the detector. This background may be rejected by measuring the opening angle of the Cherenkov cone of the candidate electron in the event [46].

Neutral-current neutrino events: Atmospheric neutrinos induce a source of background due to neutral-current interactions with nuclei in the medium. In the case of liquid scintillator detectors, the background is induced by interactions with ^{12}C where their de-excitation produces radiation. In Cherenkov detectors, elastic, neutral-current scattering of neutrinos off the target nuclei can lead to small energy depositions with a similar spectrum to the millicharged particle signature.

Charged-current neutrino events: Specifically in Cherenkov detectors, both $\nu_e/\bar{\nu}_e$, and $\nu_\mu/\bar{\nu}_\mu$ charged-current scattering events from the atmospheric neutrino flux may mimic the signal of a low-energy electron. The former directly creates low-energy electrons indistinguishable from the signal as long as the hadronic particles produced in the interaction are below the Cherenkov threshold. Despite the fact that they cannot be separated on an event-by-event basis, the recoil energy distribution of the electrons has a spectral shape different to that of the MCP signal, which allows to distinguish them statistically. Similarly, ν_μ events can mimic the background if the outgoing $\mu^\pm$ is below the Cherenkov threshold, but the Michel $e^\pm$ coming from the $\mu^\pm$ decay are visible. As in the previous case, the Michel electron energy distribution is different from a MCP statistically.

Radioactive and anthropogenic backgrounds: Radioactive backgrounds are produced by the materials in and around the detector. These can produce electrons with energies similar to those struck by MCPs, *e.g.*, in radioactive beta decays of nuclei. The dominant radioactive backgrounds in JUNO in the energy range relevant for this analysis, are ^8B , ^{10}C , ^{11}C , and ^{11}Be , which we reproduce from Ref. [28]. We note that this rate is significantly larger than currently allowed MCP signatures for $E_e \lesssim 12\text{ MeV}$, and will be a limiting factor for single-scattering searches. However, as we will discuss, multiple-hit signatures can be used to reduce these radioactive backgrounds.

Finally, we will briefly describe the data from the experiments considered in this work, and the constraints that can be established from a lack of a significant MCP signal. We derive our constraints by performing a binned-likelihood analysis, where the data and signal are organized in equally sized bins of electron recoil energy. To construct our test statistic, we compute for each one of the experiments considered, the number of signal events expected in a given bin of electron recoil energy. The number of signal events in reconstructed electron recoil energy is given by

$$\mu_i^s = T \int_{E_i-E_b/2}^{E_i+E_b/2} dE_r \int_{-1}^{+1} d\cos\theta \frac{dN}{dE_r d\cos\theta}, \quad (3.19)$$

where the symbols inside the inner most integral are defined in equation (3.18), T is the time exposure, E_b is the bin size, E_i is the bin center, and $\epsilon(E_i)$ is the average detection efficiency for a given bin.

To compute the test statistic comparing the observed data of recoil-electron-energies

with those expected from an MCP with mass m_χ , and mixing ε . We do not envision any additional triggers/cuts beyond those of the respective experiments searching for events with recoil electrons in the final state.

We adopt this strategy for Super-Kamiokande, and the XENON1T experimental results for reasons which will be detailed below.

The test statistic is calculated using the bin-by-bin likelihood function [35]

$$\mathcal{L}_i = \begin{cases} P(d_i|\mu_i^s) & d_i < \mu_i^s \\ 1 & d_i \geq \mu_i^s \end{cases}, \quad (3.20)$$

where d_i and μ_i^s are the data and expected signal (given m_χ and ε) in bin i , respectively, and $P(d_i|\mu_i^s)$ is the Poissonian likelihood of observing d_i given expectation μ_i^s . Our test statistic then is given by

$$\mathcal{TS} = -2 \log \left(\frac{\prod_i \mathcal{L}_i(m_\chi, \varepsilon)}{\prod_i \mathcal{L}_i(\varepsilon = 0)} \right), \quad (3.21)$$

where the denominator is the signal-free likelihood function and will return 1 when using the likelihood of equation (3.20).

Super-K: Our analysis uses the event selection in Ref. [46], which was designed to search for diffuse supernova background neutrinos. We use the data from SK-I, SK-II, and SK-III, and perform a joint likelihood analysis combining the three phases. For our analysis, we include the signal efficiencies provided in Ref. [46]. The background expectations estimated by the collaboration was extracted from the same reference.

XENON1T: Dark matter direct-detection experiments can measure electron recoil energies down to a few keV. We use data from the recent results of the XENON1T experiment [32] to search for signatures of MCPs produced in the earth's atmosphere. We note that the background models presented in reference [32] are insufficient to explain the observed rate of electron recoil events.

JUNO: In this case we perform a forecast where we calculate our test statistic assuming that the data collected will be consistent with the background expectation from reference [28], and perform a comparison with the expected signal-plus-background distributions which would arise from the signal produced from MCP scattering in the detector. The bin-by-bin likelihood for this case is

$$\mathcal{L}_i = P(\mu_i^b | \mu_i^b + \mu_i^s), \quad (3.22)$$

where μ_i^b (μ_i^s) is the expected background (signal) in bin i . We use the same expression for the test statistic $\mathcal{TS}$ in Eq. (3.21) above, and we assume two degrees of freedom when projecting JUNO sensitivity at 90% CL. We assume an exposure of 170 kty, which corresponds to approximately ten years of data collection.

Fig. 3.14 shows the constraints at 90% CL, compared with existing upper limits from the SLAC-mQ experiment [171], ArgoNeuT [14], the milliQan Demonstrator [42], and colliders [72]. As shown in this figure, the strongest current limit for 100 MeV $\lesssim m_\chi \lesssim 500$ MeV comes from this Super-Kamiokande analysis. However, a similar

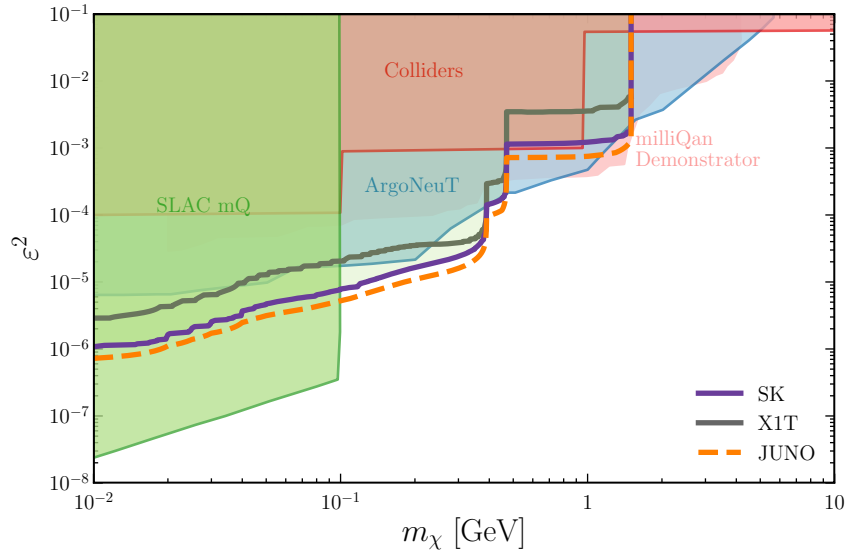


FIGURE 3.14: Constraints on millicharged particle parameter space from SLAC mQ [171], ArgonneNeuT [14], the milliQan Demonstrator [42], and colliders [72], compared with our 90% CL single-hit constraints from Super-Kamiokande (purple), JUNO (projection, orange), and XENON1T (grey).

duration of a JUNO single-hit analysis (dashed orange line) will be able to improve on this, due to JUNO's ability to reach very low electron recoil energy. On the other hand, although XENON1T (grey) can reach keV-scale electron recoils in this analysis, its small volume-exposure limits its capabilities relative to SK or JUNO.

We have found that the rate of signal events grows with smaller electron recoil energy E_r . However, in the case of both Super-Kamiokande and JUNO, the background do as well. Especially for JUNO, the background rate (dominated by radioactive emissions) became prohibitive for $E_r \lesssim 12$ MeV, preventing a single-scattering analysis from reaching lower recoil energies. One way of reducing or even eliminating these low-recoil-energy backgrounds, is to require multiple scatterings within one small time window, since the radioactive backgrounds are each associated with some relatively large half-life, such that the possibility of two emissions in a $\sim 10^{-6}$ s time window is very unlikely. Consequently, this type of analysis requires that a given MCP deposits energy at least twice as it traverses the detector. We will adopt this strategy inspired from Ref. [124], and we will focus on the JUNO experiment.

Assuming that a given MCP travels through a region with electron density n_e and length $L_{\text{det.}}$, we can estimate the probability that the MCP interact two or more times along its propagation by [38]

$$P_{n \geq 2} = 1 - e^{-\frac{L_{\text{det.}}}{\lambda_S}} \left(1 + \frac{L_{\text{det.}}}{\lambda_S} \right), \quad (3.23)$$

where $\lambda_S = 1/(n_e \sigma_S)$ is the soft-scattering⁵ mean free path, with the cross section estimated as [152]

$$\sigma_S \approx \frac{2\pi\alpha_{\text{EM}}^2 \epsilon^2}{m_e T_S} = 2.6 \times 10^{-25} \text{ cm}^2 \frac{\epsilon^2}{(T_S/1 \text{ MeV})}. \quad (3.24)$$

Here T_S is the minimum recoil energy capable of detecting events with multiple soft scatters in a small time window, α_{EM} is the electromagnetic coupling constant, and m_e is the electron's mass. Notice that T_S depends on the detector configuration. For JUNO, we estimate that $T_S \approx 0.01 - 1 \text{ MeV}$ due to the 10^4 photons/MeV and $\mathcal{O}(10^3)$ photo-electrons/MeV of the detector's liquid scintillator [83, 12].

Regarding the large background rates present at low recoil energy for JUNO from ^8B , ^{10}C , ^{11}C , and ^{11}Be radioactive decays, we assume that these backgrounds will be suppressed by requiring multiple scatterings in a sufficiently small time window. For JUNO, the emission timescale of the liquid scintillator is at most $\sim 200 \text{ ns}$ [28], so if we can restrict the time window to be $\mathcal{O}(10^{-6} \text{ s})$, then these backgrounds are reduced to being $\mathcal{O}(1)$ for the ten years of data collection we take for JUNO.

Including the probability of multiple soft-scattering events in our signal event rates, we project the sensitivity considering different recoil thresholds. The results are shown in Fig. 3.15. Note that in order to determine the detector length of JUNO that we discussed above, we simply average the path length of incoming MCP over the 35-m diameter spherical vessel. We assume that the backgrounds in JUNO are small enough such that 10 signal events are statistically significant in the 170 kt-yr exposure. We do not use any spectral information about the energy of the electron events. We take three different assumptions about the minimum threshold energy of electrons that JUNO could detect; 1 MeV (dot-dashed), 100 keV (dashed), and 10 keV (dotted). We observe that if the 10-keV threshold is attainable, we could expect that using the multiple-hit strategy the results will exceed those of single-scattering searches for atmospheric MCPs.

3.2.2 Hadrophilic Dark Matter

In this last scenario, we consider the production of atmospheric dark matter from the decay of a scalar mediator which can be produced from the decay of mesons in atmospheric showers. We consider a minimal realization within the framework of flavor-specific, hadrophilic, scalar mediator [44, 45]. The effective Lagrangian includes

$$\begin{aligned} \mathcal{L} \supset & i\bar{\chi}(\not{D} - m_\chi)\chi + \frac{1}{2}\partial_\mu S\partial^\mu S - \frac{1}{2}m_S^2 S^2 \\ & - \left(g_\chi S\bar{\chi}_L\chi_R + g_u S\bar{u}_L u_R + h.c.\right), \end{aligned} \quad (3.25)$$

where S is a singlet scalar mediator with mass m_S , while χ is the dark matter represented by a singlet Dirac fermion with mass m_χ , that is stabilized by a discrete Z_2 symmetry. On the other hand, g_u and g_χ are the couplings for the scalar with the SM, and the DM, respectively. The scalar mediates the interactions between the dark sector and the visible sector which we assume to be the up quark.

The scalar mediator S can be abundantly produced in meson decays in cosmic-ray air showers. We consider the production of the scalar from two-body decay of η , η' , and kaons. Explicit formulas for the branching fractions and decay widths can

⁵soft meaning a scattering with recoil electrons of too low an energy to be useful in single-hit searches, but energetic enough to be detected and used in a multiple-scattering search.

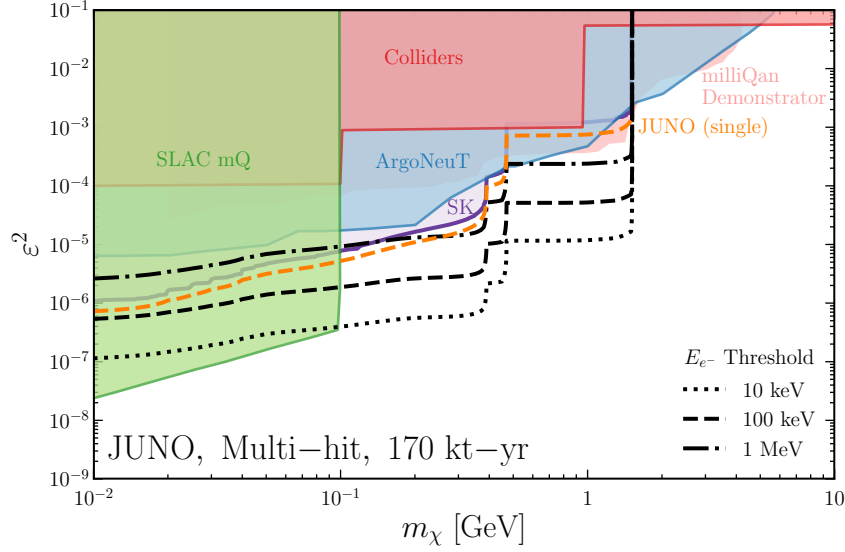


FIGURE 3.15: Black lines: Expected 90% CL sensitivity to millicharged parameter space using the multiple hit strategy for JUNO with a 170 kt-yr exposure. Three different assumptions about the minimum observable electron recoil energy are assumed: 10 keV (dotted), 100 keV (dashed), and 1 MeV (dot-dashed). Existing constraints (including our Super-Kamiokande analysis in purple) are shown as filled in regions, and the single-hit analysis of JUNO from Fig. 3.14 is shown as a dashed orange line.

be found in [36]. The production of the scalar is simulated by solving the cascade equation with a source term originated from the two-body meson decays. We consider a strong coupling between the dark matter and the mediator ($g_\chi = 1.0$), such that the production rate of dark matter can be obtained by integrating the energy distribution in the instant decay of the scalar, convoluted with its production rate. Thus we assume that the DM is produced at the same point (height) as the scalar mediator. If the coupling between the scalar and the dark matter is not strong enough, the mediator could be made a long-lived particle and the production of DM would be suppressed. Since the dark matter is stable, the expected flux at the surface of the earth can be obtained by integrating the production rate in the column depth X . The resulting DM flux at the surface of the Earth is shown in Fig. 3.16. We show each of the progenitor mesons η , η' and K^+ , considering $g_u = 10^{-3}$, $g_\chi = 1$, $m_S = 0.2$ GeV, and $m_\chi = m_S/3$.

The flux of DM with sizable interaction strength is attenuated as it traverses the medium from the top of the atmosphere towards the experiment. At large cross-sections the medium becomes “optically thick,” causing direct DM detection experiments to lose their sensitivity in this part of the DM parameter space (e.g. [151, 81]). The DM energy loss per unit depth z can be described by

$$\frac{dT_\chi}{dz} = - \sum_N n_N \int_0^{T_r^{\max}} T_r \frac{d\sigma_{\chi N}}{dT_r} dT_r, \quad (3.26)$$

where T_r is the energy lost by a DM particle in a collision with the nucleus N , while n_N is the number density, and $\frac{d\sigma_{\chi N}}{dT_r}$ is the differential cross-section between the dark

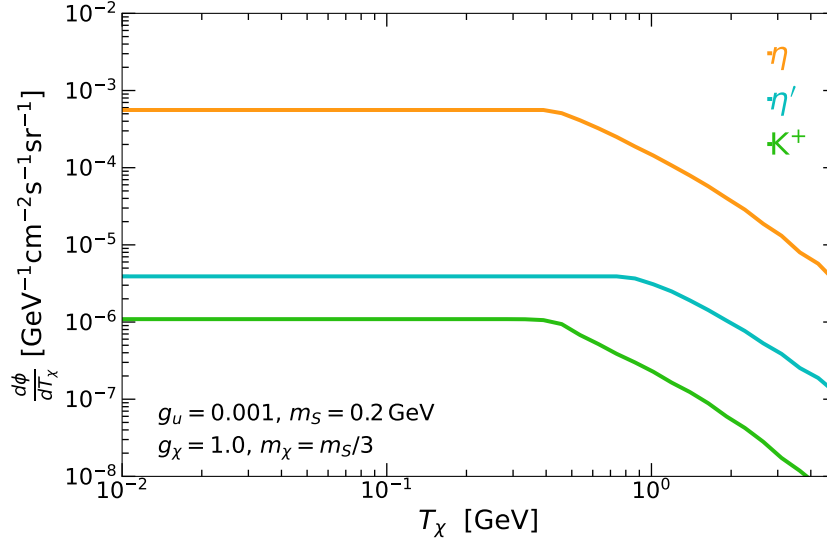


FIGURE 3.16: Dark matter flux spectrum at surface of the Earth as a function of the dark matter energy T_χ , produced from atmospheric decays of mesons as shown in the legend.

matter and the target nucleus. The expression for this cross-section with all the details involved, can be found in equations (2)-(4) of reference [36].

To account for DM flux attenuation, we numerically solve Eq. (3.26), which takes into account the relevant effects of the nuclear form factor. We note that the effects of the nuclear form factor in attenuation can significantly impact the DM parameter space to which direct DM detection experiments are sensitive to, as has been recently highlighted in Ref. [191]. The standard assumption for the energy loss of DM uses the approximation for the differential cross-section $\frac{d\sigma_{\chi N}}{dT_r} \sim \frac{\sigma_{\chi N}}{T_r^{max}}$, with T_r^{max} the maximum recoil energy, resulting in an overestimation of the DM absorption.

Once we have computed the energy loss using equation 3.26, we can simulate the expected flux at the detector for a given set of parameters m_χ, g_u . The prediction for the flux at detector can be used to estimate the differential event rate for a detector with N_t nucleons as

$$\frac{dN}{dT_r} = N_t \int_{E_\chi^{min}} dE_\chi \epsilon(T_r) \frac{d\sigma_{\chi N}}{dT_r} \frac{d\Phi_\chi}{dT_\chi}, \quad (3.27)$$

where $\epsilon(T_r)$ is the detection efficiency taken from reference [31]. In Fig. 3.17 we show the event spectra as a function of the nuclear recoil energy T_r for different target materials, including xenon (Xe), oxygen (O), hydrogen (H) and argon (Ar). The bumps in the spectrum follow the shape of the nuclear Helm factor in the interaction cross-section. We consider the model parameter values of $g_u = 10^{-3}$ and $m_\chi = 0.067$ GeV. This result highlights the importance of low $\mathcal{O}(1)$ keV experimental thresholds available for large DM experiments. We note that hadrophilic DM could also induce interactions in large neutrino experiments. However, such interactions require significantly higher thresholds and the resulting event rates are expected to be suppressed even taking into account the larger available fiducial mass.

Finally, we set limits on the model up-quark coupling g_u from XENON1T, by integrating Eq. (3.27) in the recoil energy window with a threshold of 5 keV, and a

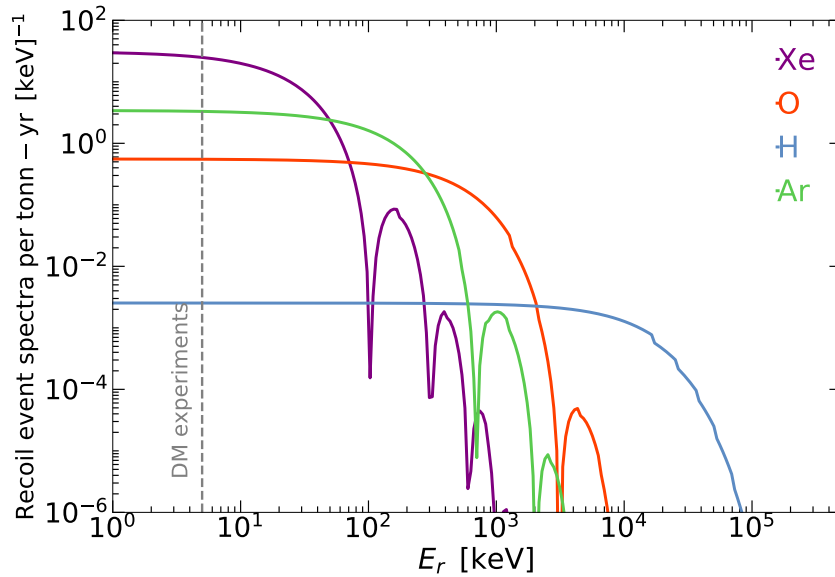


FIGURE 3.17: Event spectra for different target materials as a function of nuclear recoil energy E_r , including xenon (Xe, purple), oxygen (O, orange) and hydrogen (H, blue). The total exposure is one ton-year. The considered model values are $g_u = 10^{-3}$ and $m_\chi = 0.067$ GeV.

maximum recoil value of 40.9 keV. Following Ref. [31], we set a limit by requiring to observe more than 3.56 events for an exposure of 278.8 live-days. The limit we obtain on the plane $g_u - m_S$ is shown in Fig. 3.18. For comparison, we also display constraints on this model from beam dump experiments analyzed in reference [45] which includes the E787/E949, and MiniBooNE experiments. We note that the bounds from E787/E949 provide the strongest limits. This limit comes from measurements from the kaon decay $K \rightarrow \pi \nu \bar{\nu}$, which can be used to constraint $K \rightarrow \pi S$ through the same coupling g_u . Moreover, we also display the sensitivity reach projections for the future upgrade of XENON1T, the xenon-based DARWIN experiment [4], which is expected to have a size that is about 50 times larger with a fiducial mass of ~ 50 tons. The effects of implementing an exact numerical treatment of attenuation that also includes the nuclear form factor are shown in Fig. 3.18 with the line “absorption,” above which DM flux would be attenuated if the approximate treatment has been used instead. This further highlights the importance of attenuation for identifying and exploring the optimal DM parameter space with direct DM experiments. We note that our analysis results are independent of cosmology and the mechanism of DM production, which are not detailed in this study.

Having analyzed multiple new physical scenarios in this thesis in which we simulate the production of BSM particles in atmospheric showers for both stable and long-lived particles, it is clear that competitive constraints can be derived in both cases. Even in some (narrow) windows of the parameter space, these boundaries may cover new regions not yet tested. However, since atmospheric particle fluxes for BSM particles are generally smaller than those obtained from colliders, search optimizations with custom dedicated strategies should be considered for future studies, in order to improve the impact of these searches.

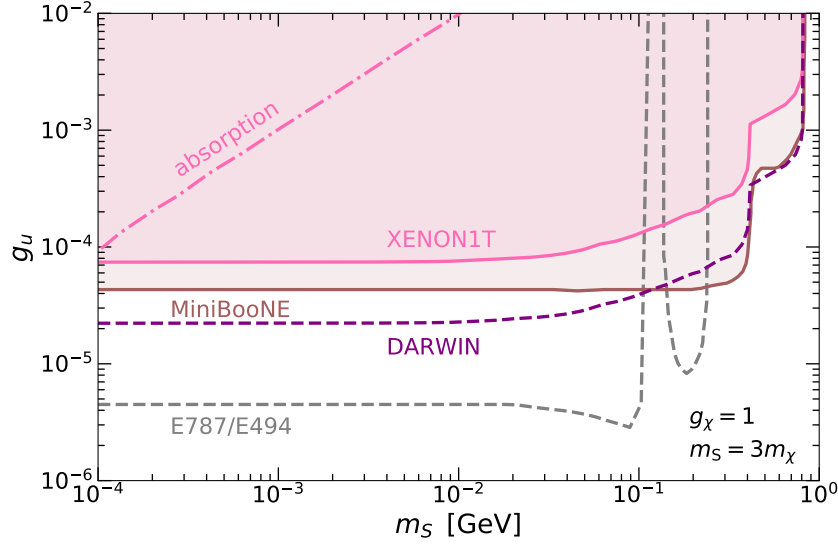


FIGURE 3.18: Limits from atmospheric DM flux from meson decays (red), for a hadrophilic scalar mediator of mass m_S on up-quark coupling g_u , considering $g_\chi = 1$ and $m_S = 3m_\chi$. We include the limits of E787/E949 and MiniBooNE which were derived in reference [45]. The “absorption” curve shows the limit obtained from overestimating the absorption of DM using a standard approximation for the differential cross section.

For long-lived particles, we found that it is possible to probe particles with lifetimes on the order of 1 to 10 km in large neutrino experiments such as Super-Kamiokande and IceCube. This is a notable feature and encourages these experiments to seriously consider dedicated atmospheric searches, mainly because our analysis is based on the statistical significance of the number of events and does not include more detail on the kinematics of the decays within the detector, which could allow for a better discrimination of the signal against the background, thus obtaining better restrictions.

Furthermore, the latter could be especially relevant for stable particles capable of making multiple impacts as they pass through the detector, as we show for the case of millicharged particles. We also note that for stable particles it is quite relevant to make an accurate estimate of the energy losses during propagation through the Earth, as this can have a large impact on the region of parameter space that can be constrained.

Finally, the work reported in the following articles provides a major effort to describe and implement a systematic search methodology for atmospheric BSM particles that includes state-of-the-art flux simulations as well as detailed elements for signal detection and analysis. The same methodologies described in this thesis could be applied to multiple new physics scenarios and experimental setups.

Chapter 4

Resumen de la Tesis

4.1 Introducción: Física más allá del Modelo Estándar y lluvias atmosféricas

Durante más de dos mil años, los seres humanos hemos estado pensando de qué está hecha la materia. Esta pregunta ha dado origen a la idea de que los objetos en nuestro universo están compuestos por partículas fundamentales. A pesar de que esto ha sido planteado desde hace mucho tiempo, ha sido el desarrollo de la teoría cuántica durante el siglo pasado lo que realmente impulsó el progreso sustancial hacia una respuesta a esta pregunta, al proporcionar una descripción precisa de las partículas elementales que componen la materia, así como la mayoría de sus interacciones básicas. El estado actual de la clasificación de todas las partículas elementales, y la descripción de tres de las cuatro interacciones fundamentales, se describen mediante el Modelo Estándar (SM) de la física de partículas. Sin embargo, existen varias motivaciones teóricas y observacionales que demandan una extensión del Modelo Estándar.

En esta tesis estudiaremos extensiones del Modelo Estándar que contienen nuevas partículas fundamentales, las cuales pueden ser producidas en la atmósfera producto de la colisión de los rayos cósmicos sobre ésta. Motivados desde una perspectiva teórica, estudiaremos un modelo supersimétrico, y además consideraremos un modelo con partículas milicargadas, el cual está motivado por el problema de la cuantización de la carga eléctrica. Por otra parte, estudiaremos modelos motivados por observaciones experimentales que incluyen el estudio de la materia oscura, así como también un modelo con leptones neutros pesados que podría resolver el problema sobre el origen de la masa de los neutrinos.

A continuación se dará una breve introducción a la motivación física detrás de los escenarios considerados en esta disertación.

4.1.1 Física más allá del Modelo Estándar

4.1.2 El problema de la jerarquía y la supersimetría

El bosón de Higgs es una de las partículas claves en el Modelo Estándar, puesto que su potencial es responsable de generar el mecanismo que otorga masa a las partículas fundamentales. La masa del bosón de Higgs ha sido medida por un valor en torno a los 125 GeV, sin embargo, teóricamente es de esperarse que la masa del Higgs reciba grandes correcciones cuánticas. Una forma de entender el por qué las correcciones cuánticas hacen tender a la masa del Higgs hacia un valor mucho más alto que el observado experimentalmente, es considerar el acoplo del Higgs a fermiones. En

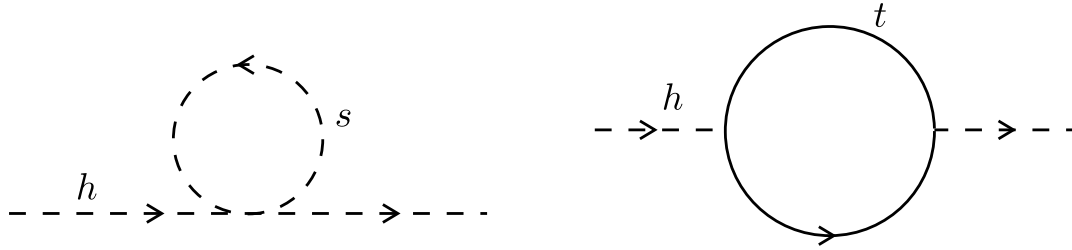


FIGURE 4.1: Diagramas a 1-loop obtenidos a partir del acoplamiento del bosón de Higgs con una partícula escalar s (lado izquierdo), y un quark Top t (lado derecho).

particular el acoplo con el quark Top como el que se muestra en lado derecho de la figura 4.1, implica una corrección a la masa del Higgs δm dado por

$$\delta m^2 \propto -y_t^2 \Lambda^2, \quad (4.1)$$

donde y_t corresponde al acoplamiento del Higgs con el quark Top, mientras que Λ corresponde a una escala de altas energías en donde se esperan nuevos efectos físicos, tales como efectos cuánticos de la gravedad ($\sim 10^{14}$ GeV) o la unificación de las fuerzas de la naturaleza ($\sim 10^{16}$ GeV). En virtud de esto, resulta enigmático entender el por qué la masa del Higgs y la escala electrodébil tienen un valor tan pequeño respecto a lo que se esperaría considerando las correcciones cuánticas. Esta diferencia se conoce como el problema de la jerarquía, y en el Modelo Estándar requiere de un ajuste muy fino de los parámetros para poder explicar este fenómeno.

Una posible solución a este problema es propuesta por la supersimetría, en donde se postula una simetría que relaciona las propiedades entre bosones y fermiones. En una teoría supersimétrica cada bosón tiene un supercompañero fermiónico y viceversa. De este modo, por cada fermion que se acople al Higgs se tendrá un escalar con una interacción como la del lado izquierdo de la figura 4.1, que dará origen a una corrección dada por

$$\delta m^2 \propto y_s \Lambda^2. \quad (4.2)$$

Si cada uno de los quarks y leptones del Modelo Estándar está acompañado por dos campos escalares complejos con $y_s = y_f^2$, entonces las contribuciones de las ecuaciones (4.1) y (4.2) pueden cancelarse entre sí y no es necesario ajustar los parámetros relacionados con la masa de Higgs para explicar su valor.

Supersimetría es construida mediante una extensión del algebra de Poincaré, añadiendo transformaciones espacio-temporales generalizadas que conectan fermiones y bosones. Esto se conoce con el algebra de Super-Poincaré. Teorías de campo supersimétricas invariantes de gauge pueden ser construidas respetando esta álgebra, y el modelo con el conjunto mínimo de partículas que se puede construir se llama Modelo Estándar Supersimétrico Mínimo (MSSM) [76]. Aunque este modelo es compatible con medidas electrodébiles, sigue siendo un modelo defectuoso. Por ello la investigación de frontera en éste tópico se lleva a cabo considerando extensiones del MSSM que también pueden resolver algunos de los otros problemas principales del Modelo Estándar, como las masas de neutrinos y la materia oscura. En particular, en esta tesis se estudiará un modelo supersimétrico que incluye neutralinos de larga vida, los cuales pueden ser producidos en nuestra atmósfera y detectados en experimentos subterráneos tales como Super-Kamiokande.

4.1.3 El problema de cuantificación de carga y partículas milicargadas

En el Modelo Estándar de la física de partículas la carga eléctrica se obtiene a partir de la tercera componente del isospin y la hipercarga. Esta última se asigna para reproducir la observación experimental de las cargas de las partículas del SM. Observamos en la naturaleza que los quarks están confinados en hadrones que tienen cargas que son múltiplos de la carga del electrón. Además, los mesones son pares quark-antiquark que tienen carga cero si están formados por dos pares up o dos down, o $\pm$ la carga del electrón si mezclan quarks up y down. De manera similar, los bariones están formados por tres quarks, ya sea top o up, la carga siempre es múltiplo de la carga del electrón. Por lo tanto, parece que la carga eléctrica está cuantificada, aunque no existe ningún principio que dicte dicha regla. Las explicaciones más populares para un principio de cuantización de carga se basan en extensiones del Modelo Estándar, incluida la unificación de las fuerzas de la naturaleza [105], la existencia de dimensiones extra [141], y la existencia de monopolos magnéticos [77]. Ninguna de estas posibles explicaciones para la cuantización de la carga eléctrica ha sido demostrada y, de hecho, los modelos con partículas que tienen una pequeña fracción de la carga del electrón,

$$Q = \epsilon e \ (\epsilon \lesssim 10^{-2}), \quad (4.3)$$

han sido contruidos y estudiados en la literatura desde hace mucho tiempo, véase, por ejemplo la referencia [171] y las referencias citadas ahí. Las partículas con una carga eléctrica tan diminuta se denominan partículas milicargadas (MCP). Las búsquedas de partículas milicargadas se ha extendido principalmente hacia experimentos modernos de colisión de partículas [170, 64, 137, 13, 41], y también dentro del contexto de la materia oscura [89, 163]. En el caso de estudio en aceleradores (es decir, colisionadores y experimentos de beam bump), las masas de los MCP que se prueban se encuentran dentro del rango MeV-GeV. Sin embargo, las restricciones para masas superiores 1.0 GeV son bastante débiles, en torno a $Q \sim 10^{-2}e$ [20].

Por otro lado, para el caso de MCPs que son a la vez materia oscura, las búsquedas se enfocan en estados más ligeros por debajo de 1.0 MeV. Este escenario está muy limitado por varias observaciones astrofísicas, incluyendo la nucleosíntesis del Big-Bang, la evolución estelar y el fondo cósmico de microondas [72].

En esta tesis estudiaremos la producción y detección de MCPs que pueden ser generarse en lluvias atmosféricas. Estas partículas pueden propagarse hasta ser detectadas en experimentos subterráneos tanto de neutrinos como de materia oscura. Dichas MCPs no tienen necesariamente que ser parte de la materia oscura.

4.1.4 La naturaleza de la materia oscura

Existe una gran cantidad de evidencia experimental basada en observaciones cosmológicas, que muestran que del total de la energía de nuestro universo alrededor de un 26% corresponde a una forma de materia no bariónica que parece interactuar solo a través de la fuerza gravitatoria y no es capaz de absorber o emitir luz. Esta forma de materia es la llamada Materia Oscura (DM) [131, 48, 17, 16, 19]. La historia de la materia oscura ha evolucionado dramáticamente desde los años treinta, cuando se descubrió que la velocidad de dispersión de los objetos pertenecientes al distante cúmulo de Coma era demasiado grande para ser explicada únicamente por la distribución de la materia visible en el cúmulo [193], hasta hoy en día donde es ampliamente aceptado que nuestro modelo para la evolución cosmológica del universo está bien descrito por el modelo Λ CDM. Este modelo describe un universo

lleno de tres componentes principales, uno de ellos es la materia oscura, la cual juega un papel fundamental en la formación de estructuras a grandes escalas en nuestro universo [93].

Observaciones de fotones del Fondo Cósmico de Microondas (CMB) proporcionan información valiosa sobre la distribución de la materia del universo. A partir de ciertas mediciones en la anisotropía y ajustes del espectro del CMB, es posible obtener los principales parámetros del modelo Λ CDM. De estos parámetros, la relación $\Omega_i = \frac{\rho_i}{\rho_c}$ entre el valor presente de la densidad de energía de la especie i , y la energía crítica necesaria para observar un Universo plano, ha sido medido para las principales componentes del universo [19]

$$\Omega_b h^2 = 0.02237 \pm 0.00015, \quad \Omega_c h^2 = 0.1200 \pm 0.0012, \quad \Omega_\Lambda h^2 = 0.3167 \pm 0.0082, \quad (4.4)$$

donde $h = 0.678 \pm 0.009$ es la tasa de expansión actual del Universo en unidades de 100 km/Mpc/seg, y los subíndices b , c y Λ corresponden a materia bariónica, materia oscura fría, y energía oscura, respectivamente. Este último es el componente dominante de nuestro universo.

En este momento de la historia no hay duda de que la materia oscura existe, y es la semilla para la formación de estructuras a gran escala. Sin embargo, todavía tenemos muy poco conocimiento sobre su naturaleza y comportamiento. Las principales características de DM que conocemos son:

- La materia oscura es muy longeva, si no estable. Podemos observar los efectos gravitatorios de DM hoy, y las observaciones de materia no luminosa en el universo implican que la materia oscura debería tener un tiempo de vida al menos como la edad actual del universo (13,8 Gyr).
- La materia oscura no es bariónica. Un Objeto de Halo Compacto Masivo (MACHO) es una especie de cuerpo astrofísico hecho de materia bariónica que emite muy poca o ninguna radiación, y que viaja a través del medio interestelar sin límites de ningún sistema planetario. Debido a estas características, se ha considerado como un buen candidato bariónico para DM. Sin embargo, las búsquedas de amplificación de la luz mediante microlentes han descartado esta posibilidad [162, 26, 183].
- Si la materia oscura interactúa con el Modelo Estándar, debe ser a través de una interacción débil. Materia oscura interactuando fuertemente con el SM ha sido descartada puesto que daría lugar a una tasa de calentamiento de la tierra mucho mayor que la que ha sido medida [150]. Por otro lado, DM con una carga eléctrica estándar no sería compatible con las observaciones astrofísicas actuales. Sin embargo, una componente de DM con una pequeña carga eléctrica ϵq , ($\epsilon \ll 1$), es aún un escenario que ha sido descartado por completo [146, 10, 149]. Finalmente, debido a su naturaleza elusiva, la interacción de neutrinos con DM es el escenario menos restringido.

Las búsquedas de materia oscura generalmente se clasifican de acuerdo con tres categorías más amplias; detección directa, detección indirecta y búsquedas en colisionadores. En los experimentos de detección directa, el objetivo es observar la señal del retroceso de un núcleo objetivo, producido por la dispersión de DM galáctico con un medio en el material. Este tipo de experimentos generalmente se ubican bajo tierra

para suprimir el fondo de rayos cósmicos. Por otro lado, los experimentos de detección indirecta tienen como objetivo observar partículas SM producidas a partir de la desintegración o aniquilación de materia oscura en el espacio exterior, siendo la detección de radiación el tipo de búsqueda más común. Finalmente, en las búsquedas de colisionadores, la estrategia tradicional consiste en buscar cantidades significativas de energía y momento faltantes, ya que se espera que una vez que se produzca la colisión protón-protón, se produzca DM que escape del detector. Este enfoque suele depender mucho más del modelo escogido, sin embargo, se está realizando un esfuerzo considerable en esta dirección. En particular, materia oscura liviana producida en sectores oscuros extendidos, los cuales involucran partículas de larga vida media, se encuentra en búsqueda activa, con nuevos experimentos propuestos tales como SHiP [29, 22], y MATHUSLA [71, 67].

En esta tesis discutiremos en detalle sobre este último tipo de búsqueda que involucra sectores oscuros extendidos, sin embargo, el escenario será diferente ya que involucra la producción de nuevas partículas de larga vida media, y sectores oscuros producidos en lluvias atmosféricas.

4.1.5 El origen de la masa de los neutrinos

Durante el desarrollo del Modelo Estándar, no era un hecho establecido que los neutrinos tuvieran masa, de hecho, esto no se estableció hasta los años noventa cuando se detectó el cambio de sabor en los neutrinos atmosféricos en Super-Kamiokande [96], implicando que los neutrinos son efectivamente partículas masivas. Los autoestados de sabor que participan en interacciones electrodébiles no son idénticos a los autoestados de masa, sino que estos últimos son superposiciones cuánticas de los autoestados electrodébiles.

$$|\nu_i\rangle = \sum_{\alpha} U_{i\alpha} |\nu_{\alpha}\rangle, \quad (4.5)$$

donde $i = 1, 2, 3$ y α denota los sabores e, μ, τ . La matriz U es la denominada matriz de mezcla Pontecorvo-Maki-Nakagawa-Sakata (PMNS) [169, 153], y codifica las combinaciones unitarias entre estados propios débiles y masivos de manera similar a la matriz CKM en el sector de los quarks. La matriz PMNS se puede parametrizar con tres ángulos de mezcla, y una fase δ que viola la simetría CP ¹

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix}, \quad (4.6)$$

donde $c_{ij} \equiv \cos \theta_{ij}$, y $s_{ij} \equiv \sin \theta_{ij}$, con $\theta \in [0, \pi/2]$. La oscilación de neutrinos fue introducida por primera vez por Pontecorvo [168]. Existen dos requisitos para que se produzca este fenómeno de oscilación; los neutrinos deben presentar coherencia cuántica durante su propagación sobre distancias macroscópicas, y la incertidumbre relacionada con su momentum en la producción, y el punto de detección, debe ser lo suficientemente grande como para permitir la producción de un estado de sabor coherente. La probabilidad de que un neutrino de sabor α oscile a un sabor β después de viajar una distancia L (denominada baseline), viene dada por

$$P(\nu_{\alpha} \rightarrow \nu_{\beta}) = \sum_{ij} U_{\alpha i}^* U_{\beta i} U_{\alpha j} U_{\beta j}^* \exp \left(-i \frac{\Delta m_{ji}^2 L}{2|\mathbf{p}|} \right), \quad (4.7)$$

¹Esta parametrización asume implícitamente que los neutrinos son fermiones de Dirac, para los fermiones de Majorana existen dos fases adicionales.

donde $\Delta m_{ji}^2 = m_i^2 - m_j^2$ es la diferencia del cuadrado de la masa del neutrino, y p es el momentum del neutrino. Esta fórmula refleja claramente una superposición de ondas con amplitud que depende de la mezcla a través de la matriz PMNS. El estudio de las oscilaciones de neutrinos se basa en las mediciones de los parámetros relevantes involucrados en la ecuación 4.7, a través de experimentos que involucran fuentes de neutrinos, ya sean de origen natural, tales como el sol o la atmósfera, o también producidos en experimentos terrestres como colisionadores y reactores. De especial importancia para esta tesis es el estudio de los neutrinos atmosféricos que será tratado en el siguiente capítulo, sin embargo el lector puede consultar las referencias [178, 114], para verificar los valores actualizados de los parámetros que describen las oscilaciones de neutrinos. Notamos que las oscilaciones de neutrinos no pueden proporcionar información sobre el valor absoluto de las masas de los neutrinos, sino solo sobre las diferencias del cuadrado de masa.

Aunque la extensión más simple del SM que incluye neutrinos masivos es bastante sencilla puesto que se consigue agregando un singlete fermionico de quiralidad derecha ν_R , la forma de introducir un término de masa de neutrino no es única. A diferencia del caso de los fermiones cargados, es posible construir un término de masa de neutrino invariante de gauge en dos formas diferentes dependiendo de si el número de leptones L se conserva o no. Si se conserva el número de leptones, se puede construir un término de masa de neutrino de Dirac simplemente agregando tres neutrinos de quiralidad derecha ν_R , singletes de gauge bajo el grupo del SM. Esto permite escribir un término de acoplamiento de Yukawa

$$\mathcal{L} = -Y_{ij}^v \bar{\ell}_{Li} \tilde{H} \nu_{Rj} + h.c. \quad (4.8)$$

Luego del rompimiento espontáneo de la simetría electrodébil (EWSB), la matriz de masas de neutrinos obtenida es proporcional al valor de expectación del vacío (VEV) del campo de Higgs

$$m_D = \frac{v}{\sqrt{2}} Y_v. \quad (4.9)$$

Por otro lado, si se viola el número de leptones, entonces es posible agregar un término de masa de Majorana

$$\mathcal{L}_s = -Y_{ij}^v \bar{\ell}_{Li} \tilde{H} \nu_{Rj} - \frac{1}{2} (m_M)_{ij} \bar{\nu}_{Ri}^c \nu_{Rj} + h.c., \quad (4.10)$$

donde $\nu^c = C \bar{\nu}^T$, siendo C la matriz de conjugación de carga. El lagrangiano de la ecuación 4.10 describe el llamado mecanismo de balancín [156, 192, 104], que es la alternativa más popular para generar masas de neutrinos. Después de EWSB, el Lagrangiano en la ecuación 4.10 conduce a la matriz de masa

$$\mathcal{M} = \begin{pmatrix} 0 & m_D \\ m_D^T & m_M \end{pmatrix}. \quad (4.11)$$

Esta matriz de masas se puede diagonalizar, y los estados propios de masas de neutrinos son mezclas de los activos con los singletes añadidos. Para $m_M \gg m_D$, los valores propios de masa son

$$m_{light} = m_D (m_M)^{-1} m_D^T, \quad m_{heavy} = m_M. \quad (4.12)$$

En la ecuación anterior, los subíndices $()_{light}$ y $()_{heavy}$ no han sido asignados al azar. Las masas de neutrinos producen modificaciones considerables en las perturbaciones del CMB, y en el espectro de potencia de las estructuras a gran escala. Las

mediciones actualizadas apuntan a que la suma de las masas de neutrinos activos es muy pequeña $\sum_i m_{\nu_i} < 0.12$ eV [19], y dado que se espera que m_M esté asociado con una escala de alta energía, la ecuación 4.12 implica un espectro con tres estados de neutrinos más ligeros proporcionales a $(m_M)^{-1}$. Observamos que m_{heavy} aumenta proporcionalmente con m_M , mientras que m_{light} disminuye, y por eso el modelo se llama balancín (seesaw). En capítulos futuros discutiremos en detalle las búsquedas de estados pesados de neutrinos estériles e informaremos sobre los resultados de las búsquedas atmosféricas basadas en las suposiciones más simples para los mecanismos de producción y detección.

4.1.6 Rayos cósmicos y lluvias atmosféricas

Puesto que estudiaremos la producción de nuevas partículas en la atmósfera, daremos una breve revisión a los rayos cósmicos y las lluvias atmosféricas, prestando especial atención a las principales componentes que representan el fondo para búsquedas de nueva física en experimentos subterráneos donde se espera que estas partículas atmosféricas dejen una señal detectable.

Los rayos cósmicos (CR) son partículas altamente energéticas que viajan a través del espacio-tiempo a velocidades relativistas. Cerca de 90% de ellos son protones, y el resto son en su mayoría partículas alfa ($\sim 9\%$), y otros núcleos más pesados como el Carbono y el Oxígeno. El espectro de energía cubre una amplia gama de escalas de energía, desde alrededor de 10^9 eV, hasta 10^{20} eV, e incluso mayores. Para poder medir un rango de energía tan amplio, existen diferentes técnicas de detección. Para energías alrededor de $\sim 1 - 10^4$ GeV, experimentos satelitales como el Espectrómetro Magnético Alpha (AMS) [21], han realizado mediciones directas de los rayos cósmicos en el espacio exterior mediante el uso de una matriz de componentes que proporciona información sobre el momento y la carga de los núcleos. Por otro lado, los CR de muy alta energía ($> 10^4$ GeV) se miden indirectamente mediante experimentos terrestres de lluvias atmosféricas como el Observatorio Pierre Auger [2] y KASCADE-Grande [30]. Estos experimentos terrestres pueden medir una fracción de las partículas supervivientes producidas como secundarias por la colisión de los CR primarios con la atmósfera terrestre, y utilizar esta información para inferir y reconstruir las propiedades de los rayos cósmicos entrantes. Esta información incluye la composición, la energía y la dirección. La figura 4.2 muestra el espectro de los rayos cósmicos en función de la energía cinética de las partículas, medida por varios experimentos que incluyen globos satélites, arreglos de detectores terrestres y satélites. Notar que el eje vertical de la figura corresponde al cuadrado de la energía de la partícula, multiplicado por la tasa de partículas que llega por unidad de energía cinética, área y ángulo sólido. Además, vale la pena resaltar que para energías menores a 10 GeV por nucleón, el flujo de CR medido puede variar durante el ciclo solar, este es el efecto denominado “modulación solar” [66], que se discutirá más adelante.

El espectro de la figura 4.2 no puede explicarse mediante una sola ley de potencia, lo que indica un origen de múltiples componentes para los rayos cósmicos. La parte de baja energía del espectro (10–1000 GeV) está bien modelada por una ley de potencia con $E^{-\gamma}$, siendo $\gamma \approx 2.7$ el índice espectral [102]. A energías más altas alrededor de 3 PeV, se puede observar un empinamiento del espectro denominado “la rodilla”, el cual puede ser ajustado por un índice espectral diferente $\gamma \approx 3.1$. Finalmente, se observa que para CR de ultra alta energía a nivel de los EeV, existe

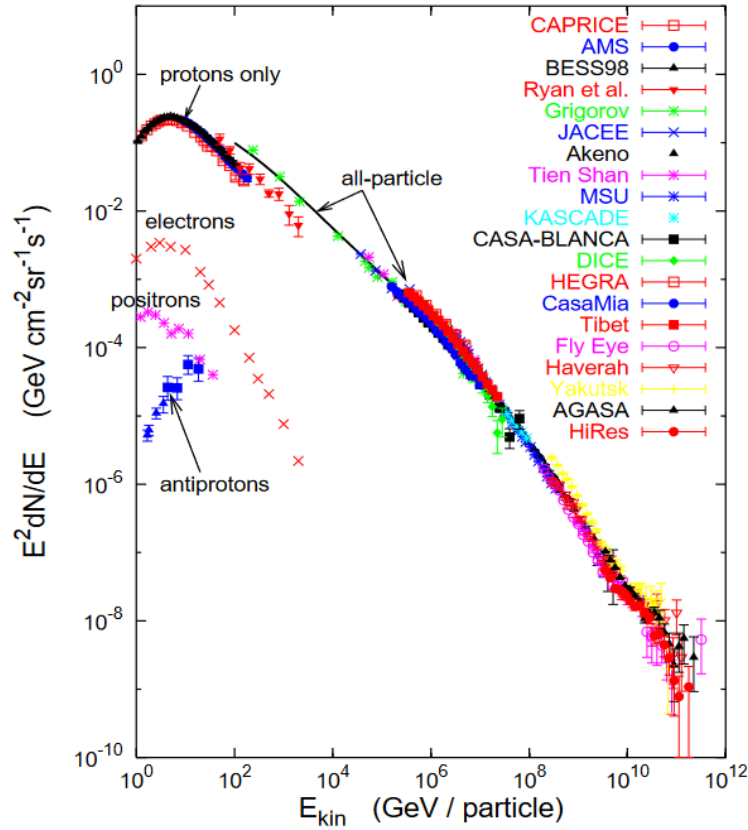


FIGURE 4.2: Espectro de energía de rayos cósmicos de todas las partículas en función de la energía cinética de la partícula, medida por diferentes experimentos terrestres y satelitales como se indica en la leyenda. Figura de la referencia [98].

una suavidad del espectro, el denominado “tobillo”, con $\gamma \approx 2.6$. Observe que el espectro de la figura 4.2 indica que el flujo de rayos cósmicos es de aproximadamente 1 partícula por $\text{m}^2\text{-s}$ por debajo de la rodilla, 1 partícula por $\text{m}^2\text{-año}$ alrededor de la rodilla, y solo 1 partícula por $\text{km}^2\text{-año}$ alrededor del tobillo, por lo que para detectar rayos cósmicos de alta energía se necesitan conjuntos de detectores terrestres muy grandes.

Cuando las partículas de rayos cósmicos llegan a la tierra, chocan con nuestra atmósfera, principalmente con átomos de nitrógeno y oxígeno a través de una interacción de dispersión inelástica. Esta colisión produce una gran cascada de partículas más ligeras, incluidos protones, neutrones, partículas alfa, hadrones, fotones y leptones [40, 117]. La altura a la que tiene lugar la colisión de los CR primarios con los átomos de la atmósfera varía con la densidad de la atmósfera, pero se encuentra entre 10 y 20 km sobre la superficie de la Tierra. Los componentes principales de cualquier lluvia atmosférica corresponden a las denominadas lluvia electromagnética y hadrónica. Las lluvias de hadrones son el componente central de las lluvias atmosféricas, estos se producen directamente en la colisión de los CR primarios con los núcleos atmosféricos que participan en interacciones fuertes. Aproximadamente 90% de los hadrones producidos en estas interacciones son piones [163], y los demás componentes contienen principalmente kaones y bariones ligeros. Mientras estas partículas secundarias se propagan a través de la atmósfera, pueden decaer e interactuar

con otros núcleos produciendo numerosas partículas nuevas que impulsan la propagación de la lluvia. Este proceso tendrá lugar hasta que la energía por nucleón sea inferior a ~ 1 GeV (energía mínima para la producción de piones). La descomposición de los piones da lugar a dos componentes importantes incluyendo muones y lluvias electromagnéticas.

Por otro lado, la lluvia electromagnética es alimentada por la lluvia hadrónica a través de la fotoproducción. Los fotones interactúan con la materia en la atmósfera generando pares de electrones y positrones, que al mismo tiempo emiten más fotones a través de la radiación bremsstrahlung. Dado que la mayoría de los hadrones en la lluvia se descomponen rápidamente o vuelven a interactuar, la mayor parte de la energía de los CR primarios finalmente encuentra su camino hacia la componente electromagnética.

La evolución de las partículas en una lluvia atmosférica puede describirse mediante un conjunto de ecuaciones integro-diferenciales acopladas denominadas ecuaciones en cascada. El objetivo principal es encontrar una solución para la tasa de espectro de partículas de un cierto tipo h

$$\Phi_h = \frac{dN_h}{dEdtdAd\Omega}, \quad (4.13)$$

que corresponde al número de partículas por unidad de energía, tiempo, unidad de área y ángulo sólido, típicamente expresado en unidades de $[\text{GeV s cm}^2 \text{ sr}]^{-1}$. El desarrollo de una lluvia atmosférica generalmente se describe en términos de la profundidad de inclinación X , que indica la cantidad de atmósfera atravesada por las partículas en la lluvia en unidades de $[\text{gr/cm}^2]$. Esta cantidad se calcula integrando la densidad del aire ρ a una distancia ℓ del suelo

$$X(\ell) = \int_0^\ell dl \rho(l). \quad (4.14)$$

Considere el cambio en el número de partículas de tipo h en una pequeña parte del volumen atmosférico. El número de partículas puede disminuir cuando las partículas se descomponen en otro tipo de partícula i o cuando se aniquila al interactuar con núcleos de aire. Por otro lado, el número de partículas de tipo h también puede aumentar cuando una partícula de tipo i ingresa al volumen atmosférico y se desintegra en una partícula de tipo h , y también cuando interactúa con el medio. La ecuación de cascada tiene en cuenta todas estas posibles interacciones.

$$\begin{aligned} \frac{d\Phi_h}{dX}(E, X) = & -\frac{\Phi_h(E, X)}{\lambda_{int,h}(E)} \\ & -\frac{\Phi_h(E, X)}{\lambda_{dec,h}(E, X)} \\ & -\frac{\partial}{\partial E}(\mu(E)\Phi_h(E, X)) \\ & + \sum_i \int_E^\infty dE_i \frac{dN_{i \rightarrow h}^{int}}{dE}(E_i, E) \frac{\Phi_i(E_i, X)}{\lambda_{int,i}(E_i)} \\ & + \sum_i \int_E^\infty dE_i \frac{dN_{i \rightarrow h}^{dec}}{dE}(E_i, E) \frac{\Phi_i(E_i, X)}{\lambda_{dec,i}(E_i, X)}, \end{aligned} \quad (4.15)$$

donde λ_{int} es el largo de la interacción, λ_{dec} es el largo del decaimiento, $\mu(E)$ es la potencia de frenado, mientras que las funciones $dN_{i \rightarrow h}/dE$ son distribuciones diferenciales de energía originadas a partir de procesos de desintegración y dispersión.

La suma cubre todos los tipos de partículas i que pueden producir el tipo de partículas h .

La ecuación de cascada 4.15 puede ser integrada con respecto a la profundidad de inclinación X , dando una estimación del flujo dependiente de la energía en la superficie de la tierra, que se puede comparar con mediciones experimentales. Se puede encontrar una solución cualitativamente buena para las ecuaciones de cascada de forma analítica, considerando un espectro primario modelado por una ley de potencia para los CR, y una interacción de escala invariante para las secciones transversales [184, 148, 102, 136]. Sin embargo, en esta tesis nos centraremos en soluciones numéricas basadas en el software Matrix Cascade Equation (MCEq) [86, 87, 85], que proporciona resultados de última generación para el espectro de flujo de hadrones y leptones.

Finalmente, comentaremos sobre las principales componentes de las lluvias atmosféricas que llegan a la superficie de la tierra; neutrinos y muones atmosféricos.

Los muones tienen una vida media relativamente larga ($\sim 10^{-6}$ s), y una pequeña sección transversal de interacción con los núcleos de aire, por lo que no sorprende que los muones atmosféricos sean las partículas cargadas más abundantes que llegan al nivel del mar. Además, son capaces de penetrar profundamente bajo tierra, produciendo un fondo relevante para los telescopios de neutrinos y otros experimentos subterráneos que buscan partículas BSM. Los muones son producidos mediante el decaimiento de hadrones, con contribuciones parciales que pueden ser clasificadas en dos componentes; una contribución originada por desintegraciones ‘prompt’, producidas a partir de mesones y bariones de vida media muy corta que contienen quarks charm, como por ejemplo los mesones D y el barión Λ_c^+ . La otra componente es denominada ‘convencional’, y se origina principalmente por la descomposición de piones y kaones a más bajas energías. La componente convencional es varios órdenes de magnitud mayor que la component prompt, sin embargo a energías a partir de los 10^6 GeV esta última componente domina la contribución al flujo. Múltiples mediciones muestran que los muones que llegan a la superficie de la tierra desde direcciones verticales se dan a razón de una partícula por cm^2 por minuto [55, 73, 120]. A nivel del mar, una forma muy común de medir los muones atmosféricos es mediante el uso de un tubo Geiger-Muller, en el que se ioniza el gas con el paso de una partícula cargada, produciendo una descarga eléctrica que crea un pulso. Para detectores subterráneos, el método más común de detección de muones es a través de radiación de Cherenkov, en donde los muones que viajan más rápido que la luz en un determinado medio, emiten radiación electromagnética que es detectada por tubos fotomultiplicadores.

Al igual que los muones, los neutrinos atmosféricos se producen principalmente en las desintegraciones de piones y kaones. Sin embargo, a energías más bajas también hay un componente importante de desintegración de muones. La cadena de desintegración completa de los piones que producen neutrinos tanto electrónicos como muónicos es

$$\pi^\pm \rightarrow \mu^\pm + \nu_\mu(\bar{\nu}_\mu) \rightarrow e^\pm + \nu_e(\bar{\nu}_e) + \bar{\nu}_\mu(\nu_\mu) + \nu_\mu(\bar{\nu}_\mu). \quad (4.16)$$

Esta cadena de decaimiento implicaría en principio que si todos los muones decaen, debería haber un factor 2 en la proporción de neutrinos muónicos a neutrinos electrónicos. Sin embargo, en 1994 el experimento Kamiokande encontró que esta

relación es aproximadamente igual a 1, dando lugar a la llamada anomalía de neutrinos atmosféricos [95, 132, 61, 47]. La explicación de esta observación se basa en la oscilación de los neutrinos. Partiendo de la fórmula general de oscilación de neutrinos, es posible demostrar que en el caso más simple de mezcla de dos familias, la probabilidad de oscilación es [113, 139, 128]

$$P(\nu_\alpha \rightarrow \nu_\beta) = \sin^2 2\theta \sin^2(\phi_{ij}), \quad (4.17)$$

con $\alpha \neq \beta$, y ϕ_{ij} la fase de oscilación de neutrinos

$$\phi_{ij} = 1.27 \frac{\Delta m^2 (\text{eV}^2) L (\text{km})}{E_\nu (\text{GeV})}. \quad (4.18)$$

En la ecuación 4.17 la amplitud de la oscilación depende del ángulo de mezcla y es máxima para $\theta = \pi/4$. La frecuencia de la oscilación está controlada por la diferencia de masas Δm^2 , y el radio entre la distancia de oscilación y la energía del neutrino L/E_ν . Para los neutrinos atmosféricos, la diferencia de masas relevante es del orden de $\Delta m^2 \sim 10^{-3} \text{ eV}^2$, y considerando una energía típica de neutrino de alrededor de 10 GeV, las oscilaciones son visibles para neutrinos propagándose a lo largo de cientos de kilómetros, viajando a través de la tierra desde debajo del horizonte. Super-Kamiokande estableció la existencia de oscilaciones de neutrinos con una mezcla casi máxima de neutrinos muónicos y neutrinos tau, al descubrir que el número de neutrinos muónicos ascendentes, es decir, los generados en el otro lado de la Tierra, son la mitad del número de neutrinos muónicos descendentes, lo que implica la desaparición de este tipo de neutrinos en el escenario anterior. El valor de la diferencia de masa de neutrinos atmosféricos al cuadrado, tiene un valor de medición actual de $|\Delta m_{23}^2| = (2,40^{+0,11}_{-0,12}) \times 10^{-3} \text{ eV}^2$, mientras que el ángulo de mezcla atmosférica es $\sin^2 \theta_{23} = 0.44 \pm 0.008$, ambas mediciones están dadas aquí asumiendo un orden normal de jerarquía de masas [182].

Los flujos de neutrinos se pueden calcular usando MCEq, y al igual que con los muones, MCEq permite estimar la contribución individual de cada progenitor al flujo de neutrinos. En el caso de los neutrinos tau, la contribución dominante es de los leptones tau en todo el rango de energía. Por otro lado, los neutrinos electrónicos y muónicos reciben una contribución dominante de los kaones en el rango de energía de $10^2 - 10^6 \text{ GeV}$ y, a energías más altas, la contribución dominante proviene de los mesones D, que dominan el componente prompt del flujo. La parte del espectro de baja energía requiere un tratamiento más delicado y no está incluido en MCEq. El flujo de neutrinos atmosféricos alcanza su punto máximo alrededor de los 100 MeV, donde los efectos geomagnéticos y la actividad solar son relevantes. Por una parte, el campo magnético terrestre impacta los rayos cósmicos fuera y dentro de la atmósfera. En el interior de la atmósfera desvía las partículas cargadas que se producen en los secundarios, mientras que en el exterior actúa como filtro excluyendo aquellas partículas de baja energía. El rango de energía relevante en el que el efecto produce un impacto importante depende de la ubicación y dirección, de tal forma que el flujo debe calcularse por separado para cada ubicación del detector. El impacto del efecto geomagnético sobre los rayos cósmicos y los flujos de neutrinos se ha medido desde hace años, siendo Super-Kamiokande una de las principales colaboraciones que ha realizado múltiples mediciones locales y simulaciones de este efecto [99]. Por otro lado, el sol emite un plasma magnetizado que los rayos cósmicos tienen que difundir para alcanzar la atmósfera terrestre y, como resultado, durante los periodos de alta actividad solar, la parte de baja energía del espectro CR se encuentra suprimido. El impacto total en el flujo de neutrinos se encuentra en torno al 10% – 20%, y depende de la ubicación geográfica del experimento [134].

4.2 Búsqueda de nuevas partículas producidas en la atmósfera

En esta sección presentaremos una discusión de los resultados principales encontrados en los distintos proyectos de investigación. Estudiaremos la producción y detección de partículas BSM en lluvias atmosféricas. Vamos a separar las búsquedas de estas partículas BSM entre aquellas que son de larga vida media, es decir, partículas que una vez producidas en la atmósfera pueden llegar a la superficie, penetrar bajo tierra y desintegrarse dentro de grandes detectores, y aquellas partículas que son estables y que pueden penetrar profundamente bajo tierra, produciendo una señal detectable al colisionar con un material objetivo.

4.2.1 Partículas de larga vida media: Leptones Neutros Pesados y Neutralinos

Leptones Neutros Pesados

Un leptón neutro pesado (HNL) es un fermión con números cuánticos de neutrinos estériles. Seguiremos la notación de las referencias [34, 69], y consideraremos un modelo mínimo que realiza masas de neutrinos basado en el Seesaw Tipo-I, el cual cuenta con al menos dos singletes de Majorana pesados N_j

$$\mathcal{L}_N \supset \sum_j i \bar{N}_j \not{\partial} N_j - \left(Y_{\alpha j} \bar{\ell}_\alpha \tilde{H} N_j + \frac{m_{N_j}}{2} \bar{N}_j N_j^c \right), \quad (4.19)$$

donde $\tilde{H} \equiv i\sigma_2 H^*$ es el complejo conjugado del campo de Higgs H , ℓ_α es un doblete leptónico del SM con sabor α , Y es una matriz de Yukawa, y m_{N_j} es la masa de Majorana del singlete N_j . Los HNLs se acoplan a los campos del SM de la misma manera que los neutrinos activos pero con un acoplamiento suprimido que es proporcional al ángulo de mezcla $U_{\alpha j} = m_{D\alpha} m_{N_j}^{-1}$, con $\alpha = e, \mu$, o τ . Los leptones neutros pesados pueden producirse en interacciones de corrientes cargadas (CC), mediadas por los bosones $W^\pm$, o en interacciones de corriente neutra (NC), mediadas por el bosón Z , en completa analogía con la producción de neutrinos. Además los HNL pueden decaer mediante los mismos tipos de interacción. Dado que los HNL no degenerados se producen y decaen de forma independiente sin oscilaciones entre ellos, podemos describir la fenomenología relevante con cuatro parámetros; la masa HNL y los tres ángulos de mezcla U_α . Sin embargo, al estudiar la producción y detección de HNLs, consideraremos un escenario en el que se asume que uno de los ángulos de mezcla es mucho mayor que los otros dos, por lo que el espacio de parámetros relevantes es solo uno de los acoples de mixing y la masa.

El principal mecanismo de producción de HNL en las lluvias atmosféricas es similar a los neutrinos, con contribuciones al flujo de HNL proveniente de las desintegraciones de mesones, las cuales pueden ser descritas mediante una forma simplificada de la ecuación 4.15

$$\frac{d\Phi_{\text{HNL}}}{dX}(E, X) = \sum_M \int_E^\infty dE_M \frac{dN_{M \rightarrow \text{HNL}}^{\text{dec}}}{dE} \frac{\Phi_M(E_M, X)}{\lambda_{\text{dec}, M}(E_M, X)}, \quad (4.20)$$

donde la suma corre sobre todas las contribuciones posibles de mesones que pueden decaer en la atmósfera a un HNL, dado un mixing U_α . La distribución $\frac{dN_{M \rightarrow \text{HNL}}^{\text{dec}}}{dE}$ incluye la fracción de HNL producidos en la desintegración de los mesones, así como también el ancho de decaimiento diferencial, que en general depende de los

parametros fundamentales U_α y m_N , así como de las energías de las partículas involucradas en el proceso. Se ha considerado la producción de HNL a partir de la desintegración de piones, kaones, mesones D , y del leptón τ , cuando el coupling U_τ es el dominante. Se han analizado dos escenarios diferentes, uno para masas por sobre el kaon, donde hemos tomado una dependencia genérica en la fracción de decaimiento y la vida media. Mientras que para masas por debajo del kaon, hemos considerado la dependencia completa y detallada de la fracción de decaimiento, y vida media a través de los parametros fundamentales U_α y m_N . Nótese que la ecuación 4.20 entrega información sobre la tasa diferencial de partículas producidas a una cierta altura, por columna de aire atravesado. Para encontrar el flujo en la superficie de la tierra, es necesario incluir la probabilidad de decaimiento de la partícula a lo largo de su propagación. La probabilidad de decaimiento es incluida pesando la tasa de producción con la exponencial $\exp(-l/L_{\text{dec}})$, que contiene el cociente entre la distancia recorrida l desde el punto de producción hasta la superficie, y el largo de decaimiento del HNL: $L_{\text{dec}} = \hbar/\Gamma_N$. Este largo de decaimiento contiene los factores relativistas, así como la dependencia con los parámetros fundamentales a través de la vida media τ_N .

Una vez que se conoce el espectro del flujo en el detector dados tanto la fracción de composición y la vida media, es posible estimar el número de eventos que se puede esperar en un detector subterráneo de gran volumen. El número total de HNL que puede decaer dentro de un detector durante una ventana de tiempo ΔT , con energías en el intervalo $[E, E + dE]$, y coseno del ángulo zenith en $[\cos \theta, \cos \theta + d\cos \theta]$, se puede estimar por

$$\frac{dN}{dE d\cos \theta} = \Delta T A_{\text{eff}}(E, \cos \theta) \frac{d\phi_N}{dE d\cos \theta}, \quad (4.21)$$

donde $\frac{d\phi_N}{dE d\cos \theta}$ es el flujo diferencial, mientras que $A_{\text{eff}}(E, \cos \theta)$ es el área efectiva que incluye la probabilidad de que se produzca un decaimiento del HNL dentro del detector. Esta área se puede estimar integrando la superficie del detector normal a la dirección del flujo, ponderada por la probabilidad de decaimiento del HNL dentro del detector. Detalles de la expresión analítica para A_{eff} puede ser encontrada en [34]. Para estimar el área de decaimiento efectiva suponemos que el volumen fiduciario del detector es un cilindro de radio R y altura H . Realizaremos una búsqueda de HNL atmosférico utilizando datos de dos grandes detectores de neutrinos subterráneos, IceCube (IC) y Super-Kamiokande (SK).

A continuación describiremos los conjuntos de datos utilizados para estudiar un posible exceso de desintegraciones HNL dentro de estos detectores de neutrinos. La estrategia consiste en que si hay un exceso de eventos con suficiente estadística, para los parámetros involucrados en la producción y descomposición de HNL, entonces estos parámetros pueden excluirse ya que implicaría que tanto IC o SK observarían un exceso de eventos atmosféricos por sobre el fondo esperado.

El Observatorio de neutrinos IceCube [7] es un detector de neutrinos de 1.0 km^3 en el Polo Sur, optimizado para detectar neutrinos con energías superiores a 100 GeV . Consideramos la muestra de datos correspondiente al análisis presentado en la referencia [5], para el cual las áreas efectivas para la interacción de neutrinos están disponibles públicamente en la página web de IceCube [1]. La estadística corresponde a 641 días de eventos de neutrinos atmosféricos y astrofísicos. Las muestras de datos reconstruidos se dividen en tracks y cascadas, tanto para neutrinos provenientes desde los hemisferios norte y sur. Se proporcionan áreas efectivas para cada

una de estas muestras, incluyendo información sobre el tipo de interacción y energía real del neutrino. A partir de estas áreas efectivas podemos inferir las eficiencias del detector dividiéndolas por las secciones transversales de interacción correspondientes. Suponemos además que los eventos de desintegración HNL se pueden medir con eficiencias similares a las interacciones de neutrinos de corrientes cargadas.

$$N^{\text{IC}}_i = \text{Br}(\text{HNL} \rightarrow e - \text{like}) \int d \cos \theta dE \epsilon \nu^{\alpha\beta}(E_i, E, \theta) \mathcal{P}_{\text{pass}}(E, \theta) \frac{dN(E, \cos \theta)}{dE d \cos \theta}, \quad (4.22)$$

donde tomamos $0.2 \leq \cos \theta \leq 1$, mientras que $\alpha = \text{cascade}$ y $\beta = \text{CC } e\text{-like}$. Ponderamos la señal por una fracción de decaimiento que incluye la descomposición del HNL a electrones, fotones, leptones τ , o lluvias hadrónicas en el estado final, que se observarían como eventos tipo-electron (o tipo cascada) en el detector. En la ecuación 4.22, incluimos las fracciones de paso del veto muónico $\mathcal{P}_{\text{pass}}(E, \theta)$, el cual reduce los eventos con una señal de muon dentro de una ventana de tiempo de $3 \mu\text{s}$. Este veto se ha aplicado a los de IC, por lo que debemos incluirlo para mantener la coherencia en el análisis. Dado que nos enfocamos en eventos similares a una cascada, las fracciones de paso más apropiadas para usar son las correspondientes ν_e , que se toman de Ref. [37].

Super-Kamiokande es un detector de Cherenkov significativamente más pequeño, pero tiene un umbral de energía más bajo donde el flujo atmosférico es mucho más alto. Solo consideraremos eventos tipo cascada, es decir, eventos con una topología tipo-electrón, o $e\text{-like}$, como se definió anteriormente. En este caso, hemos utilizado la muestra total de eventos $e\text{-like}$ de la referencia [8] con un ángulo cenital reconstruido en el rango $0 \leq \cos \theta \leq 1$. En este caso, si bien tenemos información sobre el ángulo cenital de la muestra del evento, no tenemos información disponible sobre la energía de los eventos. Por lo tanto, para SK agrupamos la señal de HNL solo en $\cos \theta$ e integramos sobre todas las energías de estas partículas para valores entre 1.0 GeV y 90 GeV, que es el rango de energías de neutrinos principales para eventos contenidos de tipo-electron. También asumimos que las eficiencias para la detección del decaimiento de HNL son similares a las de los eventos CC ν_e en el rango de energía que consideramos, además asumimos que la eficiencia es plana en este rango. El número de eventos en el i -ésimo bin en $\cos \theta$ se puede calcular como:

$$N^{\text{SK}}_i = \text{Br}(\text{HNL} \rightarrow e - \text{like}) \int_{\cos \theta_i^{\min}}^{\cos \theta_i^{\max}} d \cos \theta \int_{1.0 \text{ GeV}}^{90 \text{ GeV}} dE \epsilon_{\text{SK}} \frac{dN(E, \cos \theta)}{dE d \cos \theta}, \quad (4.23)$$

donde $\cos \theta_i^{\min}$ y $\cos \theta_i^{\max}$ son los límites inferior y superior del bin, y $\epsilon_{\text{SK}} = 0.75$. La región del espacio de parámetros que puede restringirse es obtenida a partir de la observación de un exceso de eventos sobre el fondo. La señal que buscamos sería un exceso de eventos, donde el vértice de interacción está contenido en el volumen fiduciario del detector. Suponemos una distribución de Poisson para los eventos, y consideramos como background los datos de neutrinos atmosféricos detectados en ambos experimentos. Aplicando un test de chi-cuadrado para HNL con masas por encima de la masa de kaón, encontramos que dado que la producción de HNL a partir de mesones D -mesones y leptones τ es muy pequeña, mientras que la producción y la desintegración están fuertemente correlacionadas, los límites que se pueden obtener en el plano $m_N - U_a$ son muy poco restrictivos. En virtud de esto,

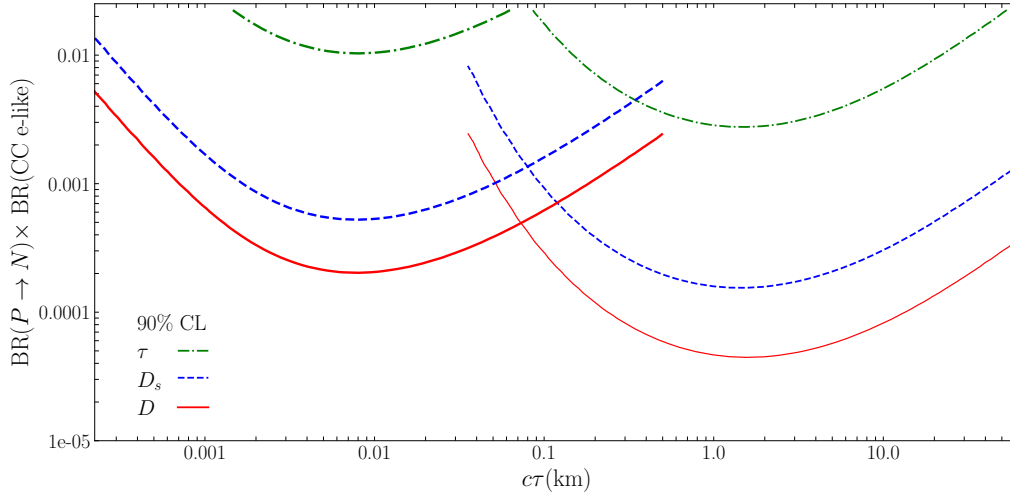


FIGURE 4.3: Límites para un HNL de masa $m_N = 1.0$ GeV, en Ice-Cube (curvas gruesas) y Super-Kamiokande (curvas delgadas). Las cotas se encuentran en el plano $BR(P \rightarrow N) \times BR(N \rightarrow CC\text{-}e\text{-like})$ vs $c\tau$, e incluyen diferentes tipos de líneas para las distintas partículas progenitoras $P = D$ (curvas solidas), D_s (curvas punteadas), y τ (curvas semi-punteadas).

consideramos un escenario más general donde tanto la producción como la desintegración son independientes, lo que significa que eliminamos la dependencia de los parámetros m_N, U_α en la producción y desintegración, y consideramos la fracción de ramificación total $Br_{\text{tot}} = Br(P \rightarrow \text{HNL} + e) \times Br(\text{HNL} \rightarrow e\text{-like})$, y el tiempo de vida τ_N , como los parámetros libres para restringir. De esta forma, nuestros resultados son generales y pueden traducirse a otros escenarios no mínimos donde la producción y el decaimiento no están necesariamente correlacionados. Nuestros resultados se muestran en la Fig. 4.3 para el HNL, asumiendo $m_N = 1.0$ GeV. Las líneas gruesas corresponden a los límites de IC, mientras que las líneas finas indican el resultado de SK.

Los dos experimentos observan eventos en regímenes de energía muy diferentes: $\mathcal{O}(\text{TeV})$ en el caso de IC, mientras que $\mathcal{O}(\text{GeV})$ en SK, y por lo tanto, el factor de impulso (que también depende de la masa del HNL) será muy diferente. Como resultado, sus sensibilidades son altamente complementarias y exploran diferentes regiones de $c\tau$. Para IC, el mejor límite se alcanza para $c\tau \sim \mathcal{O}(10^{-3} - 10^{-2})$ km, mientras que para SK los mejores límites se obtienen para $c\tau \sim \mathcal{O}(1 - 10)$ km.

Finalmente, discutiremos los resultados que se obtienen para HNL producido a partir de desintegraciones de piones y kaones. La producción de HNL a partir de piones y kaones es mucho mayor que la producción de mesones D. En este caso incluiremos un análisis con la correlación y dependencia explícita de producción y decaimiento a través de los parámetros fundamentales m_N, U_α . Además, dado que ya sabemos por el resultado encontrado en la figura 4.3 que la sensibilidad de IceCube es peor que la de Super-Kamiokande, el análisis se centrará en el conjunto de datos de este último. También notamos que en este caso usamos los eventos multi-GeV totalmente contenidos para los tipos $\bar{\nu}$ -like y e-like. Consideramos solo un U_α dominante para cada caso. Nuestros resultados se muestran en la figura 4.4 para $|U_e|^2$ y para $|U_\mu|^2$

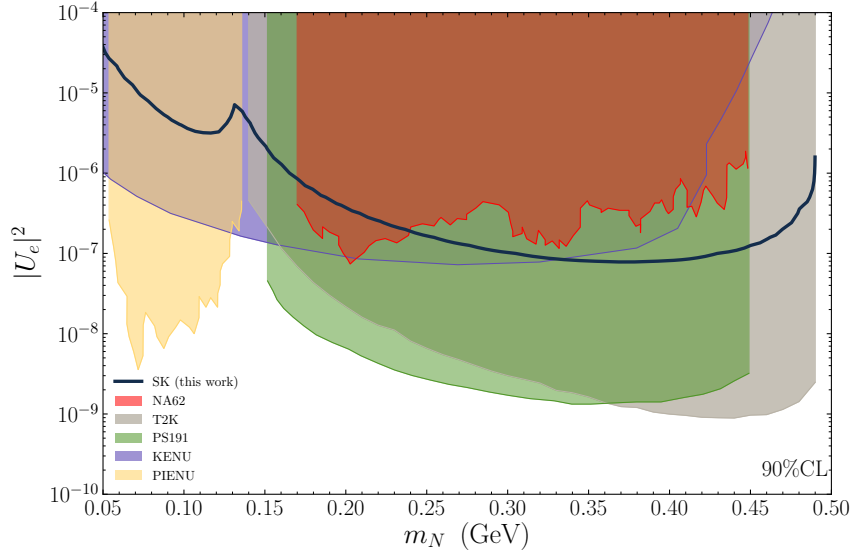


FIGURE 4.4: Límites al 90% de nivel de confianza para la búsqueda de HNL en Super-Kamiokande, proyectados en el plano $|U_e|^2$ vs m_N (con $|U_\mu|, |U_\tau| = 0$). Nuestros resultados (líneas negras continuas) se comparan con los límites obtenidos en T2K [9], NA62 [70], PS191 [50, 49], y PIENU [157]. La curva etiquetada como KENU fue obtenida en la referencia [53] usando mediciones de precisión en desintegraciones leptónicas de kaones.

en la figura 4.5, con un 90 % nivel de confianza. En el caso de $|U_e|^2$, la contribución de las desintegraciones de $\pi^\pm$ es claramente dominante para $m_N < 140$ MeV, como puede verse en el pico de sensibilidad alrededor de 0,1 GeV. Los límites derivados de nuestro análisis ya pueden establecer restricciones estrictas en la mezcla de HNL con los sectores de neutrinos de electrones y muones, entre 10^{-6} y 10^{-7} para m_N en el rango entre 150 MeV y 450 MeV. Nuestros límites también se comparan con los obtenidos de búsquedas de desintegración desplazada en PS191 [50, 49] y en el detector T2K [9], así como de búsquedas de picos en E949 [39], PIENU [157] y NA62 [70]. También mostramos el límite resultante derivado en Ref. [53] a partir de la medición de la descomposición del kaón en electrones o muones [145]. En el caso de $|U_e|^2$, los límites obtenidos de SK son comparables o incluso mejores que los límites análogos de las búsquedas de picos, mientras que no son competitivos con los de las búsquedas de decaimiento desplazado. En el caso de $|U_\mu|^2$, las búsquedas en E949 también arrojan mejores restricciones que nuestros límites de SK.

Neutralinos

Los neutralinos son sfermiones eléctricamente neutros. Como se comentó anteriormente, el MSSM es la extensión supersimétrica más simple del SM que es fenomenológicamente viable. En este modelo se la descomposición del protón está prohibida gracias a una simetría discreta Z_2 , denominada paridad R. En el MSSM hay cuatro neutralinos, siendo el más ligero la partícula supersimétrica más ligera (LSP), que proporciona un candidato interesante para la materia oscura fermiónica debido a su estabilidad y falta de interacciones electromagnéticas. Sin embargo, es posible construir otros modelos que rompen la paridad R (RPV), y que respetan los límites de la vida media del protón, siempre que se conserven el número bariónico o el número

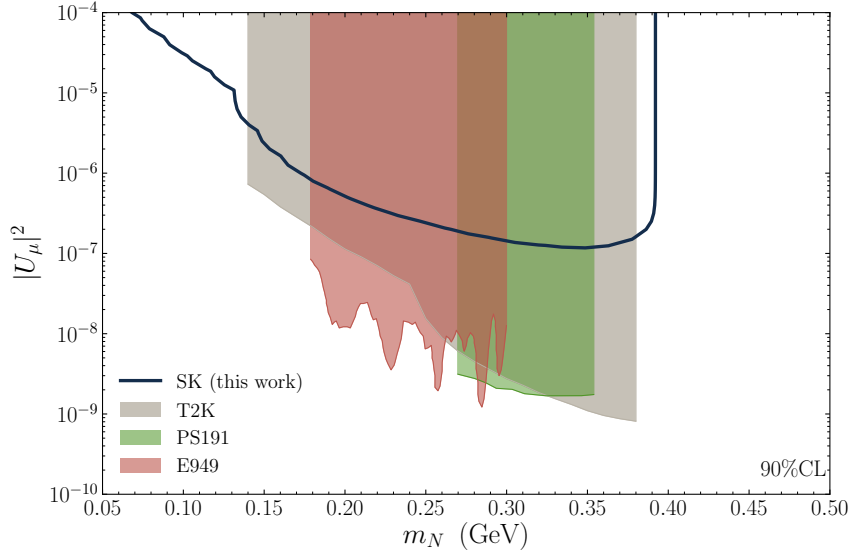


FIGURE 4.5: Límites al 90% de nivel de confianza para la búsqueda de HNL en Super-Kamiokande, proyectados en el plano $|U_\mu|^2$ vs m_N (con $|U_e|, |U_\tau| = 0$). Nuestros resultados (líneas negras continuas) se comparan con los límites obtenidos en T2K [9], E949 [39], y PS191 [50, 49].

leptónico. En tal caso, el neutralino LSP ya no es estable y puede decaer en partículas del SM. Puesto que los acoplamientos que violan la paridad R son pequeños, la vida media de estos neutralinos puede hacerse macroscópica, lo que convierte al neutralino $\tilde{\chi}_1^0$, en una partícula de larga vida (LLP). Sin imponer la paridad R, el lagrangiano más general que respeta las simetrías de gauge contiene el término [188]

$$\mathcal{L} \supset \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k^c, \quad (4.24)$$

donde los símbolos con sombrero denotan multipletes de gauge de supercampos que incluyen leptones ℓ_i y sleptons $\tilde{\ell}_i$ en $\hat{L}_i$, quarks de mano izquierda q_{jL} , squarks $\tilde{q}_{jL}$ en $\hat{Q}_j$, antiquarks de mano derecha $\bar{d}_{kR}$, antisquarks $\tilde{d}_{kR}^*$ en $\hat{D}_k^c$. Los índices familiares i, j y k van del 1 al 3, lo que implica 27 parámetros de acoplamiento trilineal independientes λ'_{ijk} . En este trabajo nos centramos en los acoplamientos trilineales λ'_{121} y λ'_{112} , que permiten producir neutralinos a partir de los decaimientos de mesones D y kaones, respectivamente. Estas desintegraciones siempre van acompañadas de un electrón, como se detalla en nuestros dos escenarios de referencia representados en la tabla 4.1. La introducción de parámetros RPV también permite que los neutralinos se desintegren en un mesón y un leptón. En nuestro primer escenario, estudiamos neutralinos con masas dentro del rango de la masa del kaón a la masa del mesón D , mientras que en el segundo escenario consideramos la masa del neutralino en el intervalo entre la masa del pión y la masa del kaon. Por lo tanto, para el segundo escenario, incluimos el trilineal λ'_{111} , ya que es el único acoplamiento asociado con los piones, y por lo tanto es necesario para hacer que el neutralino decaiga cuando su masa es menor que la masa del kaón.

La producción en la atmósfera de neutralinos es calculada usando la ecuación 4.20,

	RPV coupling	Production	Decay mode
B1	$\lambda'_{121}, \lambda'_{112}$	$D^\pm \xrightarrow{\lambda'_{121}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K_S^0 + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K^{*0} + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K^{(*)+} + e^-$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K_S^0 + \nu_e$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K^{*0} + \nu_e$
B2	$\lambda'_{112}, \lambda'_{111}$	$K^\pm \xrightarrow{\lambda'_{112}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^+ + e^-$ $\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^0 + \nu_e$

TABLE 4.1: Elección de parámetros RPV que definen a los escenarios B1, y B2. La conjugación CP de los procesos que se muestran en la tabla también contribuyen al ancho de desintegración de los neutralinos. Estos modos de desintegración son compatibles con señales tipo-electrón (o shower) en SK. Más detalles en el texto principal.

sin embargo hacemos notar que a diferencia del cálculo para los HNL, hemos agregado también la variación del flujo cuando se consideran otros modelos hadrónicos para la interacciones que involucran el scattering de estos, ya sea en colisiones hadrón-hadrón, hadrón-núcleo, o núcleo-núcleo. En MCEq, el modelo que viene configurado por defecto es el SYBILL-2.3 [88], el cual es el único que contiene los efectos relevantes para interacciones hadrónicas que involucran quarks charm, tales como los mesones D. Por otra parte, para la simulación de neutralinos provenientes de kaones, hemos considerado todos los modelos hadrónico que están actualizados con datos de calibración del LHC, los cuales incluyen los modelos QGSJET-II-04 [161], EPOS-LHC [166], y DPMJET-III [175]. Las diferencias que se pueden encontrar en estos modelos fenomenológicos surgen del tratamiento de las colisiones hadrónicas inelásticas, puesto que a altas energías se tiene la producción de partículas con momentos transversales bajos, donde la teoría standard de interacciones fuertes no es aplicable. Para evaluar el impacto de la elección de estos modelos hadrónicos en nuestros resultados finales, hemos considerado el modelo por defecto SYBILL-2.3 en comparación al modelo que presenta la mayor variación en el flujo, QGSJET-II-04, encontrando que a pesar de la diferencia en el flujo, los cambios en los límites en el espacio de parámetros son prácticamente imperceptibles.

El flujo en el detector puede ser estimado a partir de la tasa de producción en la atmósfera, incluyendo además la probabilidad de desintegración del neutralino durante su propagación, tal como se estimó para los HNL. La dependencia de la vida media es calculada a partir de los parámetros fundamentales, la masa y los acoplamientos trilineales, y similarmente para la fracción de desintegración de los mesones. Calculamos el número de eventos en Super-Kamiokande a partir de la ecuación 4.23. Para la señal, nos enfocamos en los evento tipo-electrón, originadas por la descomposición de neutralinos descrita en la tabla 4.1. En particular, para el caso de los mesones neutros en los estados finales, asumimos que decaen a una señal de tipo-electrón dentro del volumen del detector. La elección de parámetros del escenario B1 también permite K_L^0 como estado final para las desintegraciones de neutralino, pero no consideramos esto como parte de nuestra señal ya que esta partícula normalmente decaerá fuera del detector debido a su larga vida media. Consideramos como fondo los eventos de neutrinos atmosféricos clasificados como tipo-electrón,

proporcionada por la referencia [8]. Los resultados para el primer escenario que involucra producción de mesones D se muestran en la figura 4.6. Puesto que es posible determinar la dependencia de los anchos de decaimiento en la relación $\lambda'_{ijk}/m_{\tilde{f}}^2$, adoptamos dicha combinación como parámetro libre al presentar los resultados. Notamos que la cota más restrictiva que se puede encontrar en la literatura para éste escenario, proviene de la observación de la oscilación de mesones [57], la cuál involucra fijar una masa degenerada para los squarks. Como puede verse en el panel izquierdo, los límites impuestos por las oscilaciones de kaón para las masas degeneradas de sneutrino y squark, excluyen todo el espacio de parámetros dentro del alcance de sensibilidad proyectado en SHiP [185], nuestras restricciones resultantes del análisis para SK, y el límite del proceso del mono leptón Drell-Yan (DY) $pp \rightarrow \ell\bar{\nu}$ en el LHC (etiquetado como 'Colliders').

Por otro lado, dado que en Super-Kamiokande todos los productos de desintegración de neutralino enumerados en la tabla 4.1 son visibles como eventos tipo-electrón, es posible observar una señal incluso cuando el parámetro λ'_{112} se fija a cero, en cuyo caso solo hay kaones neutros en los estados finales. En tal caso, los límites de las oscilaciones de kaón no se aplican como se puede ver en la figura 4.6, panel derecho. En este caso, solo nos queda la restricción de colisionador del proceso dilepton DY $pp \rightarrow \ell^+\ell^-$ para el parámetro λ'_{121} [43]. Para este escenario encontramos que el límite obtenido de nuestro análisis con Super-Kamiokande es mejor que la restricción actual en el rango de masas correspondiente.

Las combinaciones de valores excluidas para las masas de neutralino y los acoplamien-

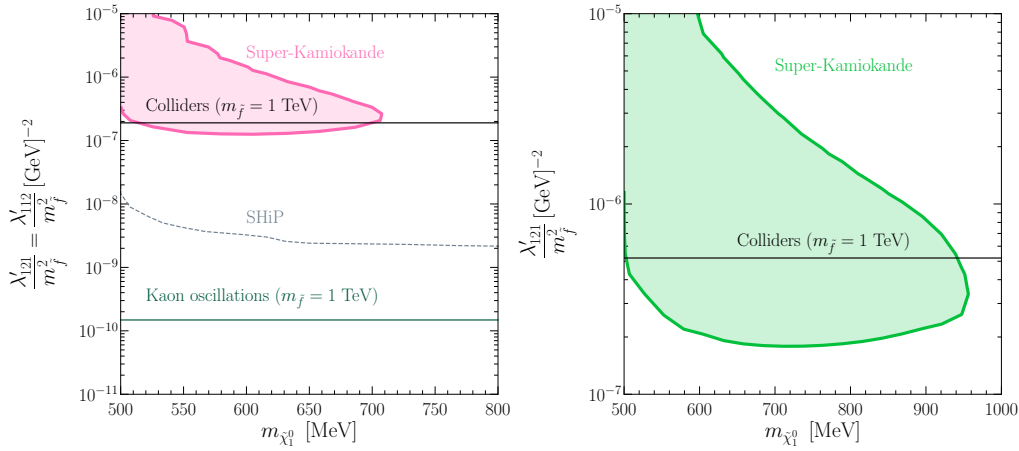


FIGURE 4.6: **Límites para el escenario B1.** Límites derivados al 90% de nivel de confianza en el espacio de parámetros definido por la masa del neutralino $m_{\tilde{\chi}_1^0}$ y los parámetros RPV trilineales λ'_{121} , λ'_{112} . En el panel izquierdo, el contorno rosa indica el espacio de parámetros excluido por SK cuando los dos acoplamientos RPV se establecen en el mismo valor, mientras que la línea discontinua muestra el alcance de sensibilidad de SHiP [185]. En el panel derecho, el contorno verde indica el espacio de parámetros excluido cuando $\lambda'_{121} \neq 0$ y $\lambda'_{112} = 0$. Las líneas sólidas indican restricciones de oscilaciones de kaón neutros [80] (curva verde) y procesos DY en el LHC (etiquetados como 'Colliders') [43] (curva negra), ambos límites son obtenidos cuando la masa de los sfermions es de 1 TeV (para más detalles, ver texto).

tos RPV nos permiten determinar la región del espacio de parámetros que se puede restringir en el plano definido por la fracción de ramificación total $\text{Br}_{\text{tot}} = \text{Br}(P \rightarrow \tilde{\chi}_1^0 + e) \times \text{Br}(\tilde{\chi}_1^0 \rightarrow e - \text{like})$, y la vida útil del neutralino. Los resultados se pueden ver en la figura 4.7. Dado que para el benchmark B1 estamos considerando los mismos canales de producción y decaimiento que en las referencias [185] y [75], podemos comparar la región actual que puede ser excluida por Super-Kamiokande, con el alcance de sensibilidad esperado para FASER y otros posibles experimentos futuros, incluidos CODEX-b, MATHUSLA y SHiP (panel izquierdo de la figura 4.7). En el caso en que los neutralinos se producen a partir de las desintegraciones de kaones en la atmósfera, no hay estudios sobre el alcance esperado de la sensibilidad en la próxima generación de detectores. Los resultados obtenidos en este caso demuestran la gran capacidad del detector de neutrinos Super-Kamiokande, para poner límites estrictos a los neutralinos, esto debido a su capacidad para probar partículas con vidas medias con una sensibilidad pico en torno a los 10 km y fracciones de ramificación totales del orden de 10^{-9} .

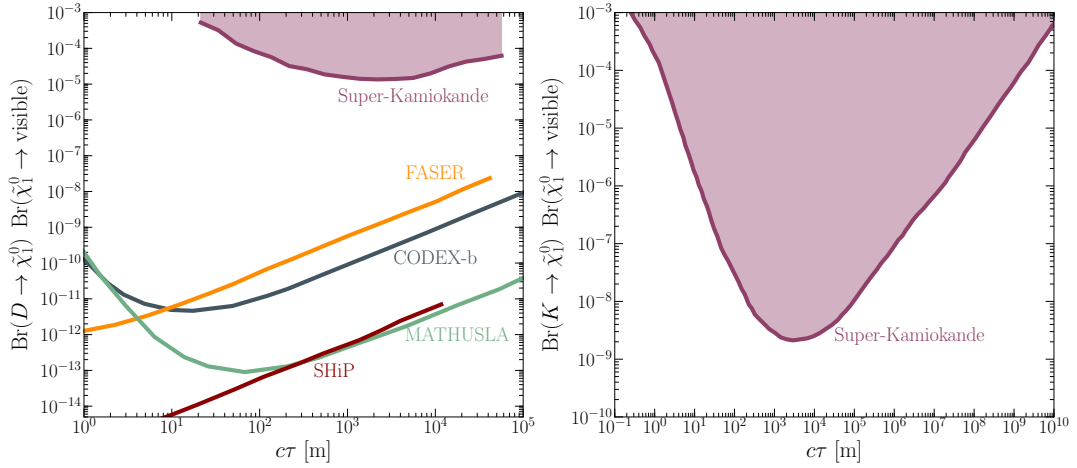


FIGURE 4.7: *Comparación de sensibilidad para experimentos explorando la frontera de la vida media.* Límites al 90% de nivel de confianza en el plano definido por la fracción de ramificación total, y la longitud propia de decaimiento del neutralino. En el panel izquierdo, el contorno morado muestra los límites de SK derivados en esta tesis, a partir del escenario B1. Las líneas de colores muestran el alcance de sensibilidad proyectado en experimentos futuros para el mismo escenario, derivados en las referencias [185, 75]. En el panel derecho, se muestra nuestro límite obtenido utilizando los datos de Super-Kamiokande para el escenario B2. Notamos que no existen otros estudios para este caso en la literatura.

4.2.2 Partículas estables: partículas milicargadas y materia oscura

Partículas Milicargadas

Como se mencionó en la introducción, una partícula milicargada o MCP es una partícula teórica con una pequeña carga eléctrica $Q = \epsilon e$, con $\epsilon \ll 1.0$ un número pequeño. Exploraremos las restricciones que se pueden imponer a la fracción de carga, cuando las MCPs se producen en la atmósfera y penetran bajo tierra dejando señales en detectores subterráneos. Haremos la suposición fenomenológica estándar con respecto a la naturaleza de la partícula milicargada, que tiene cierta masa y

un pequeño acoplamiento con el fotón a través de una pequeña carga eléctrica. Las MCPs se pueden producir a partir de la descomposición de los mesones producidos en las lluvias atmosféricas, incluidos los mesones pseudoescalares, como π^0 y η , y también mediante mesones vectoriales como ρ , ω , ϕ y J/Ψ . Otros mesones como η' también pueden contribuir al flujo de MCP, sin embargo, la contribución del resto de posibles mesones, o bien están suprimidas, o bien la tasa de producción en la atmósfera es irrelevante. Para masas mayores que $\frac{m_{J/\Psi}}{2}$ las MCPs pueden ser producidas en procesos de Drell-Yan, sin embargo se espera que la producción para este proceso sea muy pequeño, tal como se ha obtenido para otros modelos con partículas que se acoplan débilmente al SM. Simulamos la producción de MCPs resolviendo la ecuación de cascada bajo la consideración que las MCPs solo son producidas en la descomposición de los mesones, y que luego no pierden una cantidad relevante de energía al propagarse en la atmósfera. Hacemos notar además que dado que las MCPs son partículas estables, el flujo diferencial en energía puede ser directamente obtenido al integrar la ecuación 4.20 en la variable de columna de aire X . El flujo es dominado por la contribución de π_0 tal y como es de esperarse teniendo en cuenta la tasa de producción de este meson y la amplitud de decaimiento a fotones. Además hemos considerado la variación del flujo de MCP para cada mesón, cuando se cambia el modelo de interacción hadrónico y el modelo de rayos cósmicos, encontrando que puede haber una variación límite en torno al 60%. Los resultados presentes en este trabajo se darán para el modelo hadrónico SYBILL-2.3, y el modelo de rayos cósmicos de Hillas-Gaisser H3a, los cuales son los modelos más comunes para modelar flujos atmosféricos.

Para determinar el flujo de MCP en detectores subterráneos, es necesario considerar la pérdida de energía cuando estas partículas se propagan bajo la superficie de la tierra. Para ello hemos simulado la pérdida de energía utilizando una fórmula de tipo Bethe-Bloch, pero incluyendo una dependencia en los coeficientes de pérdida de energía que escala con ε^2 . Hemos tomado los parámetros de pérdida de energía para una sobrecarga de roca estándar con $a = 0.223 \text{ GeV/mwe}$, y $b = 4.64 \times 10^{-4} \text{ mwe}$, según se reportan en la referencia [143]. De esta forma, para cada MCP con una energía inicial, dirección incidente, y masa dada, se tendrá una probabilidad de que la partícula llegue con una cierta energía final luego de atravesar una cantidad de sobrecarga debajo de la superficie de la tierra. Dicha probabilidad dependerá entonces de los parámetros del modelo. Al hacer esta simulación, se encuentra que para valores relativamente grandes de la fracción de carga $\varepsilon^2 \sim 10^{-2}$, el flujo ascendente con $-1 \leq \cos \theta \leq 0$ se encuentra extremadamente atenuado. Por otro lado, para valores más pequeños $\varepsilon^2 \sim 10^{-5}$, se encuentra que debido a que las pérdidas de energía se suprimen en dos órdenes de magnitud, el flujo es mayormente independiente de la dirección incidente dada por $\cos \theta$.

Una vez que se ha simulado el flujo de MCP estimado en el detector, se considera que las partículas milicargadas pueden interactuar con el medio de éste y dispersar electrones, produciendo así una señal visible en diferentes tipos de detectores. Para este estudio hemos considerado detectores de neutrinos basados en radiación de Cherenkov, tales como SK, y detectores de neutrinos radiativos basados en centelleadores líquidos, tales como JUNO. Además, hemos considerado la posibilidad de observar una señal de electrones de dispersión en detectores de materia oscura como XENON1T. Sea $\frac{d\phi}{dE_\chi d\cos\theta}$ el espectro del flujo de MCPs en el detector, entonces la tasa de eventos provenientes de MCPs se puede obtener mediante una convolución de este flujo con la sección transversal diferencial de dispersión de electrones,

multiplicada por el número de partículas y la eficiencia de detección:

$$\frac{dN}{dE_r d\cos\theta} = n_e \epsilon(E_r) \int dE_\chi \frac{d\phi}{dE_\chi d\cos\theta} \frac{d\sigma}{dE_r}, \quad (4.25)$$

donde n_e es el número total de electrones en el detector y $\epsilon(E_r)$ es la eficiencia de reconstrucción (dependiente del detector) de estos eventos. Notamos que la sección eficaz diferencial es tal que en el régimen relativista donde $E_\chi \gg E_r, m_\chi, m_e$, escala con E_r^{-2} , de forma tal que la tasa de eventos en el detector es mayor para energías de umbrales más bajas. Sin embargo, el fondo suele también crecer para energías más bajas, siendo esto una limitación para poder maximizar el radio de la señal vs fondo. Los tipos de fondo para los experimentos analizados incluyen interacciones de neutrinos atmosféricos a través de corrientes cargadas y neutras, desintegración de muones atmosféricos, así como también electrones de baja energía producidos en desintegraciones radiactivas. Para los tres experimentos considerados en este estudio, hemos tomado los fondos reportados en la referencia [46] para SK, los fondos radiactivos de la referencia [28] para JUNO, y el fondo reportado en [32] para XENON1T. Para construir nuestro test estadístico, calculamos, para cada uno de los experimentos considerados, el número de eventos de señal esperados en un bin dado de energía de retroceso de electrones. Obtenemos restricciones realizando un análisis con una función de verosimilitud. En la figura 4.8, se muestra las restricciones obtenidas al 90% de nivel de confianza en el plano que incluye la masa del MCP y la fracción de carga ϵ , en comparación con los límites existentes del experimento SLAC-mQ [171], ArgoNeuT [14], milliQan [42], y las cotas de colisionadores [72]. Como se muestra en la figura 4.8, el límite de corriente más fuerte para 100 MeV $\lesssim m_\chi \lesssim 500$ MeV proviene de nuestro análisis usando los datos de Super-Kamiokande. JUNO puede llegar a ser competitivo puesto que es capaz de medir energías de retroceso menores. Por otro lado, aunque XENON1T (curva gris) puede alcanzar retrocesos de electrones en la escala keV, su pequeño volumen de exposición limita sus capacidades en relación con SK o JUNO.

La tasa de eventos de MCPs crece cuando la energía de retroceso de los electrones visibles E_r es más pequeña. Sin embargo, tanto en el caso de Super-Kamiokande como de JUNO, el fondo también crece a energías más bajas. Especialmente para JUNO, el fondo se vuelve prohibitivo para $E_r \lesssim 12$ MeV. Es posible reducir, o eliminar este fondo requiriendo múltiples dispersiones dentro de una pequeña ventana de tiempo, puesto que cada uno de los fondos radiactivos está asociado con una vida media relativamente larga y, por lo tanto, la posibilidad de dos emisiones de desintegración en una ventana temporal de $\sim 10^{-6}$ s, es muy pequeña. En consecuencia, este tipo de análisis requiere que un MCP que penetra bajo tierra sea capaz de depositar energía al menos dos veces al atravesar el detector. Asumiendo que un MCP dado viaja a través de una región con densidad de electrones n_e y longitud $L_{\text{det.}}$, podemos estimar la probabilidad de que el MCP interactúe dos o más veces a lo largo de su propagación por

$$P_{n \geq 2} = 1 - e^{-\frac{L_{\text{det.}}}{\lambda_S}} \left(1 + \frac{L_{\text{det.}}}{\lambda_S} \right), \quad (4.26)$$

donde $\lambda_H = 1/(n_e \sigma_S)$ es el camino libre medio para una dispersión suave, con electrones de retroceso de una energía demasiado baja para ser útil en búsquedas de señales de una sola colisión, pero lo suficientemente enérgicas para ser detectadas y utilizadas en una búsqueda de dispersión múltiple. Incluyendo esta probabilidad

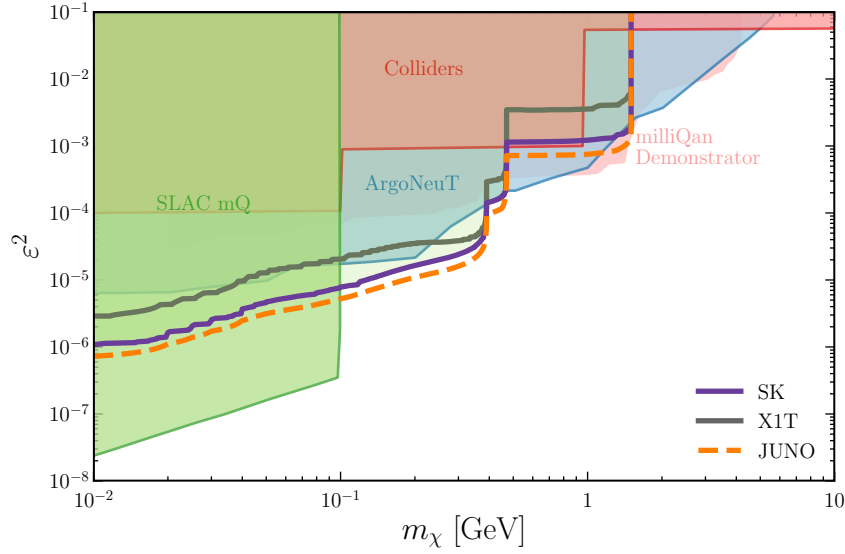


FIGURE 4.8: Restricciones en el espacio de parámetros de partículas milicargadas en SLAC mQ [171], ArgoNeuT [14], milliQan [42], y colisionadores [72], en comparación con los límites obtenidos al 90% de nivel de confianza en Super-Kamiokande (curva púrpura), JUNO (proyección futura en la curva naranja), y XENON1T (curva gris).

de múltiples eventos de dispersión en nuestro análisis, proyectamos la sensibilidad considerando diferentes retrocesos de umbral. Los resultados se muestran en la figura 4.9. Asumimos que los fondos en JUNO son lo suficientemente pequeños como para que 10 eventos de señal sean estadísticamente significativos para una exposición de 170 kt-año, y no usamos ninguna información espectral sobre la energía de los eventos de electrones, ya que esperamos que la resolución de la energía a baja escala sea difícil de percibir. Asumimos tres escenarios diferentes sobre el umbral mínimo de energía de los electrones que JUNO puede detectar; 1 MeV (curva de puntos y rayas), 100 keV (curva discontinua) y 10 keV (curva de puntos). Notar que si se pudiese alcanzar un umbral de 10 keV en JUNO, podemos esperar que la estrategia de impactos múltiples supere con creces las búsquedas de MCP atmosféricas de dispersión única.

Materia Oscura

Para la producción de materia oscura en la atmósfera, consideramos una realización mínima dentro de un modelo con un mediador escalar. Dicho mediador es hadrófilo y se acompla a un sabor específico, tal como se señala en la referencia [44, 45]. El Lagrangiano efectivo que describe este modelo incluye el siguiente término

$$\mathcal{L} \supset i\bar{\chi}(\not{D} - m_\chi)\chi + \frac{1}{2}\partial_\mu S\partial^\mu S - \frac{1}{2}m_S^2 S^2 - \left(g_\chi S\bar{\chi}_L\chi_R + g_u S\bar{u}_L u_R + h.c.\right), \quad (4.27)$$

donde S es un mediador escalar singlete bajo el SM, y de masa m_S , χ es la materia oscura representada por un fermión de Dirac con masa m_χ , la cual se encuentra estabilizada por una simetría Z_2 . Por otro lado g_u es el acoplamiento para el escalar

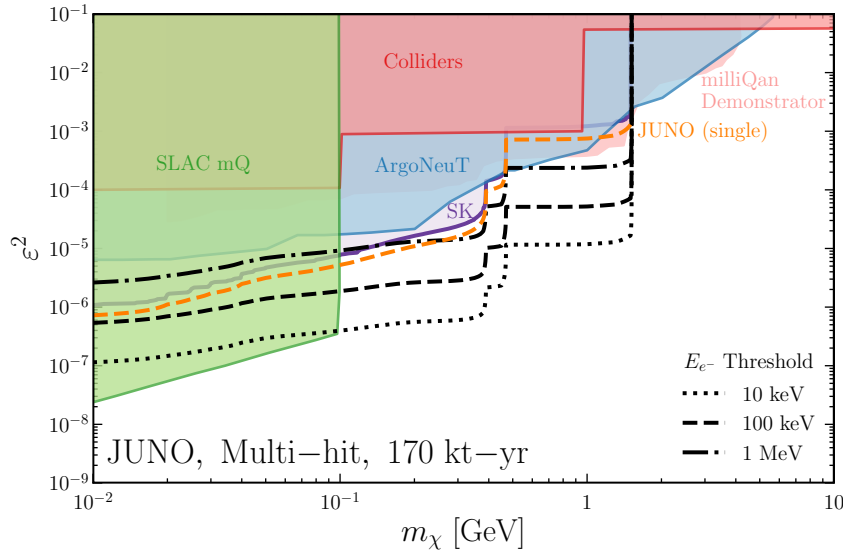


FIGURE 4.9: Líneas negras: sensibilidad esperada al 90% CL para el espacio de parámetros de MCPs, utilizando la estrategia de impacto múltiple en JUNO, tomando una exposición de 170 kt-año. Se considera tres suposiciones diferentes sobre la energía de retroceso de electrones mínima observable: 10 keV (puntos), 100 keV (curva semi-discontinua), y 1 MeV (curva de puntos y rayas). Las restricciones existentes (incluido nuestro análisis de SK en púrpura), se muestran como regiones rellenas, y el análisis de un solo scattering en JUNO de la figura 4.8 se muestra como una línea naranja discontinua.

con el quark up, mientras que g_χ es el acoplamiento del escalar con la materia oscura. En principio este modelo tiene cuatro parámetros libres, sin embargo, vamos a tomar la suposición estándar en donde el escalar está fuertemente acoplado a la materia oscura con $g_\chi = 1.0$, y que hay una relación constante entre la masa del escalar y la materia oscura dada por $m_S = 3m_\chi$. El mediador escalar S se puede producir abundantemente en las desintegraciones de mesones en lluvias atmosféricas. Consideramos la producción de este escalar a partir de la descomposición de dos cuerpos de η , η' y kaones. Usamos MCEq para extraer las tasas de producción de estos mesones, y calculamos la tasa de producción del escalar resolviendo la ecuación de cascada para un set de parámetros m_S , g_u fijos. Una vez producido el escalar, se espera que este se desintegre rápidamente a materia oscura. La tasa de producción de materia oscura se puede obtener integrando la distribución de energía en el decaimiento instantáneo del escalar, mediante una convolución con la tasa de producción de este. Se ha asumido que la DM se produce en el mismo punto espacial (altura) que el mediador escalar. Si el acoplamiento entre la materia oscura y el escalar no fuese lo suficientemente fuerte, el mediador podría convertirse en una partícula de larga vida y se suprimiría la producción de DM. Al igual como se calculó el flujo en la superficie de la tierra para MCPs, se ha integrado la tasa de producción en la columna de aire atravesada, y se ha obtenido el flujo de materia oscura proveniente de cada mesón. Para masas menores encontramos que el flujo está dominado por mesones η , a pesar que la producción de kaones en la atmósfera es mayor, mientras que para masas mayores el flujo es dominado por la contribución de η' .

Una vez se conoce el flujo en la superficie de la tierra, es necesario estimar el flujo

en el detector. La búsqueda de una señal de materia oscura en esta investigación se centrará en experimentos de detección directa, específicamente en XENON1T. A medida que la materia oscura atraviesa la tierra para llegar a este detector subterráneo, esta pierde energía mediante colisiones con núcleos en la corteza terrestre. La pérdida de energía por unidad de profundidad z es descrita por

$$\frac{dT_\chi}{dz} = - \sum_N n_N \int_0^{T_r^{\max}} T_r \frac{d\sigma_{\chi N}}{dT_r} dT_r, \quad (4.28)$$

donde T_r es la energía disipada por DM en una colisión con un núcleo N , mientras que n_N es la densidad numérica, y $\frac{d\sigma_{\chi N}}{dT_r}$ es la sección eficaz diferencial de interacción entre la materia oscura y el núcleo. Se estima la atenuación del flujo de materia oscura resolviendo numéricamente la ecuación 4.28, para diferentes valores de los parámetros $m_S g_u$. Esta ecuación tiene en cuenta los efectos relevantes del factor de forma nuclear, y hacemos hincapié en que los efectos del factor de forma nuclear en la atenuación pueden afectar significativamente el espacio de parámetros de materia oscura al que son sensibles los experimentos de detección directa de DM, tal y como veremos en los resultados de sensibilidad. Una vez que se tiene la predicción del flujo en el detector, es posible simular la tasa de eventos diferenciales, para un detector con N_t nucleones, mediante

$$\frac{dN}{dT_r} = N_t \int_{E_\chi^{\min}} dE_\chi \epsilon(T_r) \frac{d\sigma_{\chi N}}{dT_r} \frac{d\Phi_\chi}{dT_\chi}, \quad (4.29)$$

donde $\epsilon(T_r)$ es la eficiencia de detección tomada de la referencia [31]. A pesar que nuestro estudio se centra en el experimento XENON1T, hemos considerado diferentes materiales para detectar a la materia oscura incluyendo Hidrógeno, Oxígeno y Argón. La razón por la cual estudiamos esto, es porque analizamos las posibilidades de detección que tendrían los experimentos de neutrinos subterráneos como SK. Hemos encontrado que la tasa de eventos en este tipo de detectores no es tan competitiva puesto que el valor umbral de detección ocurre para energías mucho mayores, en torno al MeV, donde la señal cae rápidamente.

Finalmente, establecemos límites para el acoplamiento con el quark up g_u en XENON1T, mediante la simulación del número de eventos que se espere en el rango de energía de retroceso con un umbral de 5 keV, y un valor de retroceso máximo de 40,9 keV. Siguiendo Ref. [31], establecemos un límite al exigir que se observen más de 3,56 eventos para una exposición de 278,8 días de datos. El límite en el plano $g_u - m_S$ se muestra en la figura 4.10. A modo de comparación, también mostramos las restricciones provenientes de los experimentos E787/E949, y MiniBooNE [45]. También mostramos las proyecciones de alcance de sensibilidad para la futura actualización de XENON1T, el experimento DARWIN basado en xenón [4], que se espera que tenga un tamaño aproximadamente 50 veces mayor con una masa fiduciaria de ~ 50 toneladas. Los efectos de implementar un tratamiento numérico exacto de atenuación que también incluye el factor de forma nuclear se muestran en la figura 4.10 mediante la línea etiquetada como “absorción”, por encima de la cual el flujo de DM se atenuaría un tratamiento aproximado que no incluya el factor de forma fuese utilizado en su lugar. Esto destaca la importancia de la atenuación del flujo para identificar y explorar el espacio de parámetros de la materia oscura óptimo para experimentos de DM directos. Estos límites se traducen en un pico de sensibilidad para la sección de choque independiente de spin de la materia oscura en torno a los 10^{-33} cm^2 , para masas m_χ alrededor de los 100 MeV. Si bien los límites obtenidos en este estudio no superan los que ya existen en la literatura, constituyen otra prueba independiente del espacio de parámetros para este tipo de modelos.

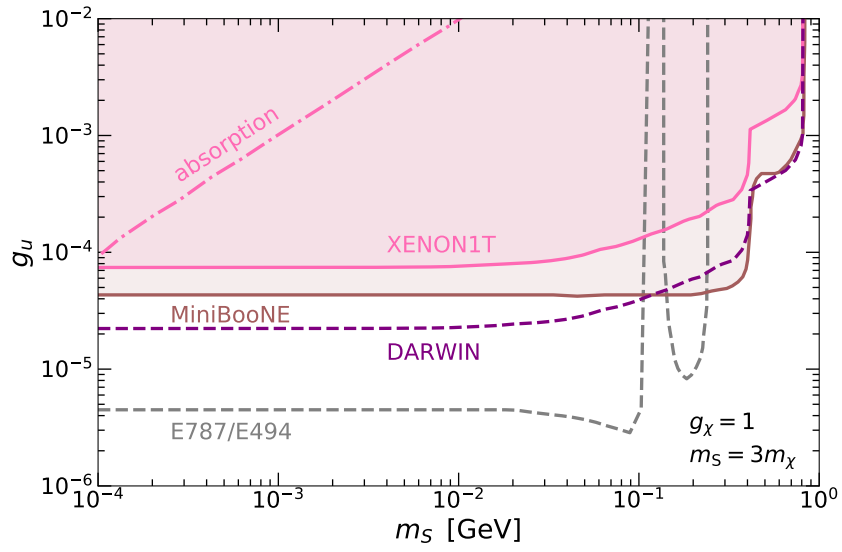


FIGURE 4.10: La región rosa muestra los límites obtenidos usando los datos de XENON1T, al 90% de nivel de confianza, en el plano definido por el acoplamiento del mediador con el quark tipo-up g_u , y la masa m_S . En este estudio se ha fijado $g_\chi = 1$, y $m_S = 3m_\chi$. La curva punteada morada es una proyección para DARWIN, mientras que las curvas de MiniBooNe y E787/E494 han sido obtenidas en la referencia [45]. Ver texto principal para más detalles.

Chapter 5

Part II: Published Articles

Here we present the published papers included in this thesis.

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Searches for atmospheric long-lived particles

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ABSTRACT: Long-lived particles are predicted in extensions of the Standard Model that involve relatively light but very weakly interacting sectors. In this paper we consider the possibility that some of these particles are produced in atmospheric cosmic ray showers, and their decay intercepted by neutrino detectors such as IceCube or Super-Kamiokande. We present the methodology and evaluate the sensitivity of these searches in various scenarios, including extensions with heavy neutral leptons in models of massive neutrinos, models with an extra $U(1)$ gauge symmetry, and a combination of both in a $U(1)_{B-L}$ model. Our results are shown as a function of the production rate and the lifetime of the corresponding long-lived particles.

KEYWORDS: Beyond Standard Model, Neutrino Physics, Solar and Atmospheric Neutrinos

ARXIV EPRINT: [1910.12839](https://arxiv.org/abs/1910.12839)

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1 Introduction

The Standard Model (SM) of particle physics has demonstrated to be an extremely predictive theory, which however cannot be complete. In particular, it lacks a mechanism to generate neutrino masses, dark matter, and the baryon asymmetry of the Universe that we observe today. While LHC upgrades and future colliders will keep pushing the energy frontier forward in order to search for new physics at high energies, it is important to keep in mind that the physics beyond the Standard Model (BSM) may very well involve weakly interacting particles at low scales instead, which in minimal extensions of the SM are expected to be relatively long-lived. In fact, the search for such long-lived particles (LLP) in

colliders and beam-dump experiments has become a field of intense activity in recent years and several experiments have been proposed to conduct dedicated searches [1–4].

In this work, we study the potential of atmospheric neutrino detectors to search for these particles since they can also be copiously produced in atmospheric showers. In order to do this, we compute the expected flux of exotic LLPs being produced in proton-nucleus collisions in the atmosphere, either through proton bremsstrahlung or as a product of meson decays. After traveling typical distances of tens of kilometers through the atmosphere, such LLPs may subsequently decay within the detector volume. The signal would therefore be an excess of events where the interaction vertex is contained in the fiducial volume of the detector. These are a priori indistinguishable from an atmospheric neutrino interactions, which therefore constitute an irreducible background to this search.

The experimental setup considered in this work has two main advantages with respect to searches in laboratory experiments. First, the center-of-mass energy in proton-nucleus collisions in the atmosphere will typically be much higher than at beam-dump experiments, allowing to produce an abundant flux of heavy mesons and tau leptons. Secondly, the size of these huge detectors results in larger decay volumes and could be optimal in the searches for particles with long lifetimes, of about $\mathcal{O}(10)$ km in the laboratory frame. In this work, we will be considering two experimental setups: Super-Kamiokande, most sensitive to events with energies below 100 GeV; and Icecube, sensitive to events with energies at the TeV scale and above. Therefore, the results between the two will be largely complementary as they will be sensitive to different values of the lifetime of the particle in its rest frame, τ . As we will see, the best sensitivities will be achieved for $\tau/m \sim \mathcal{O}(1)$ km/GeV for Super-Kamiokande and $\tau/m \sim \mathcal{O}(10)$ m/GeV for Icecube.

In this work, for concreteness, we will focus on two minimal and theoretically well-motivated extensions of the SM: (1) a scenario with heavy neutral leptons (HNL), that could be responsible for generating neutrino masses [5–8] and possibly the baryon asymmetry of the Universe [9, 10], and (2) a model with an extra U(1) symmetry which, after being broken, leads to a massive dark photon and a portal to the dark sector [11, 12]. In the first case we focus on LLPs in the mass range between the Kaon and D-meson masses, where existing laboratory bounds are weaker [13–16], while in the latter we will consider dark photons below the proton mass. The search for lighter, MeV range, atmospheric sterile neutrinos has been considered before in Super-Kamiokande [17, 18] and IceCube [19]. Our analysis significantly improves the methodology of these earlier studies.

In the most minimal models with just HNL or dark photons, both production and decay of the LLP are highly correlated and, thus, longer lifetimes also mean lower production rates. However, in non-minimal extensions this may no longer be the case, as different mechanisms may be involved in production and decay processes. An example of this framework is provided by a $B - L$ model with HNL [20]: in this case, production of the HNL could take place through the $B - L$ gauge interaction [21, 22], while the decay may occur via its mixing with the SM neutrinos. Such scenario will also be discussed in our work. Our results will be presented from a model-independent perspective, considering a more general BSM scenario where the production and decay rates could be decoupled, and we will discuss the interpretation of our results in the three scenarios outlined above.

The paper is organized as follows. In section 2 we discuss the production of LLPs and describe the method to compute LLP fluxes from two main sources: meson decays (two- and three body) and proton bremsstrahlung. In section 3 we present the general procedure to compute LLP signals in a neutrino detector which requires defining detector effective areas for decaying events, detector efficiencies, backgrounds and the data samples that will be used to extract limits on LLPs. In section 4 we focus on three BSM scenarios and use the results of the previous sections to compute the required fluxes. In sections 5 we present the sensitivities to LLPs of searches in IceCube and Super-Kamiokande, and our conclusions are drawn in section 6.

2 LLP production in atmospheric showers

The flux of any SM particle produced in the atmosphere is usually given as a function of the slant depth, X , related to the integral of the density in the direction of the particle from the top of the atmosphere to the production point at distance ℓ [23]:

$$X(\ell, \theta) = \int_{\ell}^{\ell_{\max}} \rho[h(l, \theta)] dl, \quad (2.1)$$

where $\rho(h)$ is the atmospheric density at height h , and θ the zenith angle defined by the trajectory of the particle. The distance ℓ is related to the height h and to the zenith angle θ as:

$$h(\ell, \theta) = \sqrt{R_{\oplus}^2 + 2\ell R_{\oplus} \cos \theta + \ell^2} - R_{\oplus}, \quad (2.2)$$

where $R_{\oplus} \simeq 6370$ km is the Earth radius. In our calculations, we consider $h(\ell_{\max}, \theta) = 80$ km as the maximum vertical height of the atmosphere.

In our calculations, we have used the Matrix Cascade Equation (MCEq) Monte Carlo software [24, 25] to compute the fluxes for the parent mesons and protons in the atmosphere, with the SYBILL-2.3 hadronic interaction model [26], the Hillas-Gaisser cosmic-ray model [27] and the NRLMSISE-00 atmospheric model [28]. With this procedure, meson and proton fluxes are obtained as a function of X , E and $\cos \theta$ (see appendix B for details). At this point, it is worth mentioning that these fluxes are subject to significant theoretical uncertainties related to the choice of cosmic-ray and hadronic interaction models, see e.g. figure 9 in ref. [25] or ref. [29], which can be of the order of 20–25%. Additional uncertainties may arise from seasonal and geomagnetic effects. As discussed in ref. [30], while seasonal effects are negligible at SK, they could be potentially large for the fluxes observed at IceCube. However, they tend to cancel out when you take data over a whole year (the impact of the seasonal effect for IceCube can be seen e.g. in ref. [31]). On the other hand, geomagnetic effects mostly affect the azimuthal distribution and horizontal events [32]. Since we only study time-averaged zenith distributions, their effects in the azimuthal distributions are averaged out. For SK, the three-dimensional effects change the height distribution of atmospheric neutrinos predominantly in the horizontal direction rather than the vertical direction where our signal dominates; see figures 18 and 19 in ref. [33]. In the case of the IceCube analysis effects on the horizontal events are negligible as geomagnetic effects are only relevant for neutrino energies below 20 GeV (see figure 11 in ref. [33]).

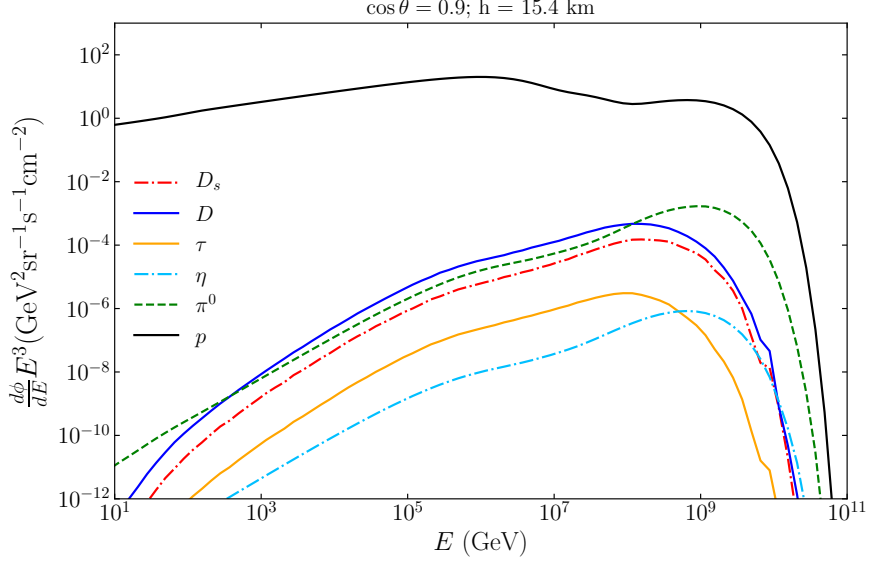


Figure 1. Differential fluxes of LLP parent particles as a function of the energy at fixed height in the atmosphere, $h = 15.4$ km.

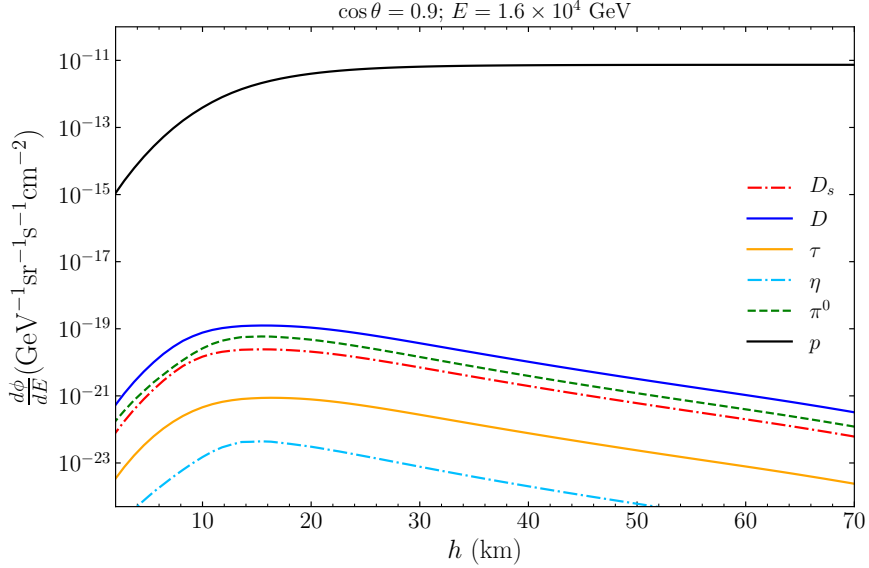


Figure 2. Differential fluxes of LLP parent particles as a function of the height at fixed energy, $E = 1.56 \cdot 10^4$ GeV.

We will assume that the boost is large enough $E \gg m_P$ so that the zenith angle of the LLP is that of the parent particle and the integration over the parent zenith angle becomes trivial. In the following we will consider models where parent particles of LLP are protons and mesons. The differential fluxes of these particles as a function of the energy are shown in figure 1 at a fixed height, and as a function of the height at fixed energy in figure 2.

In the rest of this section we discuss separately the different production mechanisms of LLPs, from SM mesons and τ lepton decays, as well as from proton bremsstrahlung. In the

case of production via meson decays, we will further distinguish between two- and three-body decays, as the outgoing LLP energy distributions will be different in the two cases.

2.1 Production from decays of SM particles

Let us first consider that a LLP, A , is produced in the decay of a SM meson P . Once the flux of the parent meson has been computed, the production profile of the LLP in the decay $P \rightarrow AY$ or $P \rightarrow AYZ$ is given by [23]

$$\frac{d\Pi_A}{dE d\cos\theta d\ell} = \sum_{ch} \int_{E_P^{\min}}^{E_P^{\max}} dE_P \frac{1}{\gamma_P \beta_P c \tau_P} \frac{d\Phi_P(E_P, \cos\theta)}{dE_P d\cos\theta} \frac{dn_{ch}(E_P, E)}{dE} \quad (2.3)$$

where the sum goes over all channels (labeled as ch) contributing to the LLP production from P decays. Here, dn_{ch}/dE is the fraction of decaying parents that produce an A with energy in the bin $[E, E + dE]$ in the decay channel ch , while γ_P , β_P and τ_P are the boost, velocity and lifetime in the rest frame of the parent particle P .

Both the distribution dn_{ch}/dE and the integration limits for E_P ($E_P^{\min}, E_P^{\max}$) depend on the decay kinematics and, in particular, they depend on whether the LLP is produced in a two- or in a three-body decay.

2.1.1 Two-body decays

In two-body decays $P \rightarrow AY$, the differential distribution dn/dE is flat in E :

$$\frac{dn(P \rightarrow AY; E_P, E)}{dE} \simeq \frac{1}{\Gamma_P} \frac{d\Gamma(P \rightarrow AY)}{dE} = \frac{\text{Br}(P \rightarrow AY)}{\Gamma(P \rightarrow AY)} \frac{d\Gamma(P \rightarrow AY)}{dE}, \quad (2.4)$$

with

$$\frac{1}{\Gamma(P \rightarrow AY)} \frac{d\Gamma(P \rightarrow AY)}{dE} = \frac{1}{p_P \sqrt{\lambda(1, y_A^2, y_Y^2)}}, \quad (2.5)$$

where $y_i \equiv \frac{m_i}{m_P}$ and

$$\lambda(a, b, c) \equiv a^2 + b^2 + c^2 - 2ab - 2ac - 2bc. \quad (2.6)$$

The kinematical limits for the energy are:

$$\gamma_P E_{\max} - \frac{p_P \sqrt{\lambda(1, y_A^2, y_Y^2)}}{2} \leq E \leq \gamma_P E_{\max} + \frac{p_P \sqrt{\lambda(1, y_A^2, y_Y^2)}}{2}, \quad (2.7)$$

where $E_{\max}$ is the energy of the LLP A when the decay takes place in the meson's rest frame, that is:

$$E_{\max} \equiv \frac{m_P}{2} (1 + y_A^2 - y_Y^2), \quad (2.8)$$

and $\gamma_P = E_P/m_P$. Alternatively if E is fixed, the kinematical limits for $E_P \gg m_P$ are:

$$\frac{m_P E}{E_{\max} + \frac{m_P}{2} \sqrt{\lambda(1, y_A^2, y_Y^2)}} \leq E_P \leq \frac{m_P E}{E_{\max} - \frac{m_P}{2} \sqrt{\lambda(1, y_A^2, y_Y^2)}}. \quad (2.9)$$

2.1.2 Three-body decays

In the case of three-body decay, $P \rightarrow AY Z$, we have:

$$\text{Br}(P \rightarrow AY Z)^{-1} \frac{dn(P \rightarrow AY Z; E_P, E)}{dE} = \frac{1}{\Gamma(P \rightarrow AY Z)} \frac{d\Gamma(P \rightarrow AY Z)}{dE}. \quad (2.10)$$

The distribution is no longer flat and the kinematical limits, neglecting masses of Y and Z , are:

$$\gamma_P \left(E'_{\max} - \sqrt{E'^2_{\max} - m_A^2} \right) \leq E \leq \gamma_P \left(E'_{\max} + \sqrt{E'^2_{\max} - m_A^2} \right), \quad (2.11)$$

with

$$E'_{\max} \equiv \frac{m_A^2 + m_P^2}{2m_P}. \quad (2.12)$$

Fixing E , the kinematical limits for E_P are therefore:

$$\frac{m_P E}{\left(E'_{\max} + \sqrt{E'^2_{\max} - m_A^2} \right)} \leq E_P \leq \frac{m_P E}{\left(E'_{\max} - \sqrt{E'^2_{\max} - m_A^2} \right)}. \quad (2.13)$$

2.2 Bremsstrahlung

An alternative production mechanism for vectors and scalars is through bremsstrahlung, in the scattering of SM particles with air nuclei. Let us consider the case of a dark photon V of mass m_V emitted by a proton p . In analogy to eq. (2.3), the production profile of the dark photon can be written in this case as [23]:

$$\frac{d\Pi_V}{dE d\cos\theta dX} = \int_E^\infty dE_p \frac{1}{\lambda_p(E_p)} \frac{d\Phi_p(E_p, \cos\theta, X)}{dE_p d\cos\theta} \frac{dn(p\mathcal{N} \rightarrow VY; E_p, E)}{dE}, \quad (2.14)$$

where Φ_p is the proton flux and λ_p is the interaction thickness of protons in air, i.e. the average amount of atmosphere (in g/cm²) traversed between successive collisions with air nuclei $\mathcal{N}$.

It is common to parametrize the coupling of the dark photon to the proton as the QED coupling multiplied by a small number ϵ . In this case, the differential distribution dn/dE takes the form [34–36]:

$$\frac{dn(p\mathcal{N} \rightarrow VY; E_p, E)}{dE} = \frac{1}{E_p} \int_0^{p_\perp^{\max}} dp_\perp^2 \frac{\sigma_{p\mathcal{N}}(2m_p(E_p - E))}{\sigma_{p\mathcal{N}}(2m_p E)} w(z, p_\perp^2), \quad (2.15)$$

where $|p_\perp^{\max}| = 1 \text{ GeV}$ and

$$w(z, p_\perp^2) \equiv \frac{\epsilon^2 \alpha_{\text{QED}}}{2\pi H} \left[\frac{1 + (1 - z)^2}{z} - 2z(1 - z) \left(\frac{2m_p^2 + m_V^2}{H} - z^2 \frac{2m_p^4}{H^2} \right) \right. \\ \left. + 2z(1 - z)(z + (1 - z)^2) \frac{m_p^2 m_V^2}{H^2} + 2z(1 - z)^2 \frac{m_V^4}{H^2} \right]. \quad (2.16)$$

Here, m_p is the mass of the proton and

$$H(z, p_\perp^2) \equiv p_\perp^2 + (1 - z)m_V^2 + z^2 m_p^2, \quad (2.17)$$

with $z \equiv p_V/p_p$, where p_V and p_p are the dark photon and proton momenta, respectively.

3 Expected number of LLP decays in atmospheric neutrino detectors

Once the production profile of the LLP has been computed, the flux of particles that arrive at the detector Φ_A is simply obtained integrating over all LLPs produced at different distances ℓ from the detector, weighted by their corresponding survival probabilities, as

$$\frac{d\Phi_A}{dE_A d\cos\theta} = \int_0^{\ell_{\max}} d\ell \frac{d\Pi_A}{dE_A d\cos\theta d\ell} e^{-\frac{\ell}{\ell_{\text{decay}}}}, \quad (3.1)$$

where $d\Pi_A$ is the flux of LLP produced at a distance ℓ of the detector within the interval $[\ell, \ell + d\ell]$, computed as outlined in section 2, and ℓ_{decay} is the decay length of the LLP in the lab frame, related to its lifetime (τ_A) as

$$\ell_{\text{decay}} = \gamma_A \beta_A c \tau_A \simeq \frac{E_A}{m_A} c \tau_A, \quad (3.2)$$

where β_A, γ_A are the boost parameters of the LLP, and we assume $E_A \gg m_A$. We will assume that, once the particle is produced, since it is very weakly coupled it does not interact any further in the atmosphere and only decays. Also, note that the upper limit of the integral in eq. (3.1) is in practice a function of θ (see eq. (2.2)): $\ell_{\max} \equiv \ell(h_{\max}, \theta)$ and $h_{\max} \simeq 80 \text{ km}$ is the maximum height of the atmosphere where cosmic showers are produced.

Using the flux in eq. (3.1), the number of decays inside the detector within a given time window ΔT , for LLPs with energies and trajectories in the intervals $[E_A, E_A + dE_A]$, $[\cos\theta, \cos\theta + d\cos\theta]$, can be computed as

$$\frac{dN}{dE_A d\cos\theta} = \Delta T A_{\text{decay}}^{\text{eff}}(E_A, \cos\theta) \frac{d\Phi_A}{dE_A d\cos\theta}, \quad (3.3)$$

where A^{eff} is an effective area which accounts for the probability that a decay takes place inside the detector. This area can be estimated integrating the surface of the detector normal to the flux direction, weighted by the decay probability of the LLP inside the detector:

$$A_{\text{decay}}^{\text{eff}}(E_A, \cos\theta) = \int dS_{\perp} \left\{ 1 - \exp\left(-\frac{\Delta\ell_{\text{det}}(\cos\theta)}{\ell_{\text{decay}}(E_A)}\right) \right\}. \quad (3.4)$$

Here, $\Delta\ell_{\text{det}}$ is the length of the segment of the LLP trajectory that cuts into the detector. Its calculation, which just depends on the zenith angle defining the trajectory of the LLP and on the geometry of the detector, is outlined in appendix A. Figure 3 shows the effective decay areas for the IceCube and Super-Kamiokande detectors as a function of the lifetime of the LLP in the lab frame.

3.1 Detector efficiencies and datasets used

The number of events will also depend on the detector efficiencies and reconstruction effects. In this work, two atmospheric neutrino detectors have been considered: IceCube (IC) [37] and Super-Kamiokande (SK) [38]. Since they observe events in two very different energy regimes, we expect their results to be complementary and to probe different regions in parameter space. In this section we describe the assumptions and data sets used for each of the two detectors separately.

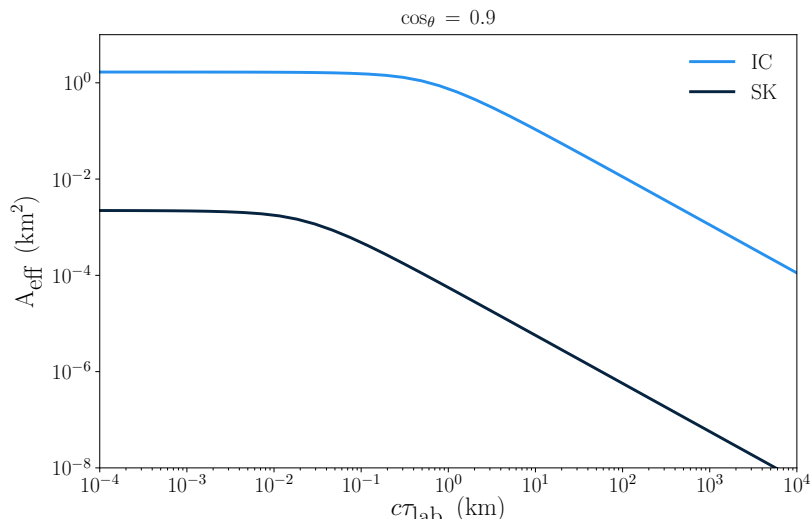


Figure 3. Effective decay area for the IceCube detector ($R_{\text{IC}} = \frac{1}{\sqrt{\pi}}$ km and $H_{\text{IC}} = 1$ km) and Super-Kamiokande ($H_{\text{SK}} = 0.04$ km and $R_{\text{SK}} = 0.02$ km) as a function of the decay length in the laboratory frame for down-going events ($\cos \theta = 0.9$).

3.1.1 IceCube

The IceCube Neutrino Observatory [37] is a $\sim 1 \text{ km}^3$ neutrino detector at the Geographic South Pole, optimized for detecting neutrinos with energies above 100 GeV. For the effective area, eq. (A.1), we consider a simplified cylindrical geometry with $H = 1$ km and volume 1 km^3 . The effective area as a function of the decay length in the laboratory frame is shown in figure 3.

We consider the data sample corresponding to the analysis presented in ref. [39], for which effective areas for interacting neutrino events are publicly available in the IceCube collaboration webpage [40]. The statistics corresponds to 641 days of contained events, and the reconstructed data samples are divided into tracks and cascade events, in the northern and southern hemispheres, and are binned in reconstructed energy. Effective areas for each of these samples are provided for each particle, interaction type and true neutrino energy. From these effective areas we can infer the detector efficiencies by dividing by the corresponding interaction cross sections. We therefore estimate the detector efficiencies as

$$\epsilon_{\nu}^{\alpha\beta}(E_{\text{rec}}^{\nu}, E_{\text{true}}^{\nu}, \theta_{\text{true}}) \simeq \frac{A_{\text{effIC}}^{\alpha\beta}(E_{\text{rec}}^{\nu}, E_{\text{true}}^{\nu}, \theta_{\text{true}})}{\rho_{\text{ice}} V_{\text{IC}} N_A \sigma_{\beta}^{\nu}(E_{\text{true}}^{\nu})} \quad (3.5)$$

where α selects the reconstructed sample: cascade/track and north/south, β selects the true neutrino interaction type, CC-e, NC, etc. In the denominator we have the total cross section in the IceCube detector σ_{β}^{ν} , which is estimated in terms of the nucleon-neutrino cross section, times the number of nucleons in the detector: Avogadro number N_A , times the ice density ($\rho_{\text{ice}} \simeq 0.92 \text{ gr cm}^{-3}$) times the fiducial volume V_{IC} . The required neutrino-nucleon cross sections are extracted from ref. [41].

We will make the assumption that our LLP decay events have the same efficiencies as charged-current (CC) neutrino interactions with a similar set of observed particles in

the final state. We also assume that the relationship between the energy of the decaying particle and the deposited energy of all particles produced in the detector is similar to the relationship between true and reconstructed neutrino energies in CC interactions. We believe that this is reasonable as long as all particles produced in the LLP decay are visible in the detector. More concretely we will consider only LLP visible decays to electrons and charged hadrons, mainly because the expected neutrino background is smaller, they are more likely to be contained, and their energy can be better determined. We assume that LLP decays of this type have an efficiency similar to those of CC e -like neutrino interactions:

$$\epsilon_{\text{LLP}}^{\alpha \text{ } e\text{-like, hadronic}}(E_{\text{rec}}, \theta_{\text{rec}}; E_A, \theta) \simeq \epsilon_{\nu}^{\alpha \text{ CC-}e}(E_{\text{rec}}, \theta_{\text{rec}}; E_A, \theta). \quad (3.6)$$

In our IC calculations, we use the efficiencies corresponding to $\alpha = \text{cascades}$. Therefore, we weight the number of signal events by the branching ratio into *e-like events*, in order to consider all decay channels which have electrons, photons, τ -leptons or hadronic showers in the final state, which would be observed as cascade-like events.

An important point is that these effective areas *do not* include the muon veto applied in the event pre-selection. This veto reduces events with a muon signal within a $3\mu\text{s}$ time window of the neutrino signal. This veto is designed to optimize the search for astrophysical neutrinos, as it reduces significantly the overwhelming background from atmospheric muons, and the atmospheric neutrino background in southern sky events [42]. At sufficiently high energies it is quite likely that muons from the same atmospheric shower that produce the LLPs will also reach the detector within this large time window, and therefore it also reduces significantly our signal. Since this veto has been applied to the background and the data, we need to impose it also on our signal to extract bounds, however it is quite likely that an optimized LLP search can be carried out, where the time window for a coincident muon track is reduced significantly, since the LLP has an atmospheric origin and an accompanying muon is not unexpected. The veto is expected to affect similarly our signal and the prompt neutrino signal, which is the dominant background, so it does not improve the signal-to-noise ratio of the LLP search. To implement this veto on our signal we use the passing fractions of the muon veto, $\mathcal{P}_{\text{pass}}(E_A, \theta)$ obtained in ref. [43], which depend on the true energy and zenith of the LLP. Given that we are concentrating on cascade events, and in most LLP production there are no associated muons, the most appropriate passing fractions to use are those corresponding prompt ν_e .

The number of signal events in the i -th energy bin at IC can be finally obtained as

$$N_i^{\text{IC}} = \text{Br}(A \rightarrow e\text{-like}) \int d\cos\theta \int dE_A \epsilon_{\nu}^{\alpha\beta}(E_i, \theta_{\text{rec}}; E_A, \theta) \mathcal{P}_{\text{pass}}^{\text{prompt-}\nu_e}(E_A, \theta) \frac{dN}{dE_A d\cos\theta}, \quad (3.7)$$

where $0.2 \leq \cos\theta_{\text{rec}} \leq 1$ corresponds to the southern sky, while $\alpha = \text{cascade}$ and $\beta = \text{CC } e\text{-like}$.

It might be useful to also include DeepCore atmospheric data. However, the detector volume in this case is much smaller than for IceCube, and the energy range considered for atmospheric neutrino analyses for DeepCore is comparable to that of the multi-GeV data in SK. Therefore, we expect similar results for DeepCore as those obtained for SK. It may

however be interesting to search for a signal using atmospheric neutrinos with energies between 100 GeV and 1 TeV, which may be sensitive to a different range of values of $c\tau$. This is unfortunately not possible using currently available public data.

3.1.2 Super-Kamiokande

Super-Kamiokande is a significantly smaller water Cherenkov detector, but it has a lower energy threshold where the atmospheric flux is much higher and therefore might have a comparable or even better sensitivity than IceCube to the scenarios considered. In the case of SK, in the computation of the effective decay areas (eq. (A.1)) we have assumed a cylindrical geometry with $H = 0.04$ km and $R = 0.02$ km. The result is shown in figure 3.

We will again only consider electron-like events, that is, events with an e-like topology as defined above. We have used in this case the fully-contained multi GeV e-like sample from ref. [44] (including single and multi-ring) with reconstructed zenith angle $0 \leq \cos \theta \leq 1$. We extract data as well as neutrino background from figure 5 of [44]. In this case, while we do have information regarding the zenith angle of the event sample (see figure 5 in ref. [44]) we do not have available information regarding the energy of the events. Therefore, for SK we bin the data only in $\cos \theta$ and integrate over all neutrino energies between 1 GeV and 90 GeV, which is the range of parent neutrino energies for multi-GeV e-like contained events (see figure 6 of ref. [44]). We think this is conservative, as SK may be sensitive to events outside this range. We also assume that the efficiencies for the decay of the LLP are similar to those of electron neutrino CC events in the multi-GeV range. Our only assumption regarding detection efficiencies for Super-Kamiokande is that they are flat as a function of energy, which we think is a reasonable assumption for fully contained events. We note that migration matrices for SK are not publicly available.

The number of events in the i -th bin in $\cos \theta$ can therefore be computed as:

$$N_i^{\text{SK}} = \text{Br}(A, e\text{-like}) \int_{\cos \theta_i^{\min}}^{\cos \theta_i^{\max}} d \cos \theta \int_{1 \text{ GeV}}^{90 \text{ GeV}} dE_A \epsilon^{SK} \frac{dN}{dE_A d \cos \theta}, \quad (3.8)$$

where $\cos \theta_i^{\min}$ and $\cos \theta_i^{\max}$ are the lower and upper limits of the bin. In this case, we have assumed a flat detection efficiency $\epsilon^{SK} = 0.75$, in line with the values quoted in ref. [44] for the multi-GeV ν_e event sample.

4 Atmospheric LLP in selected scenarios

We now consider three simple BSM scenarios of very weakly interacting sectors that can lead to LLPs. In the first two examples, both production and decay are controlled by the same couplings, in such a way that the requirements of long enough lifetimes and large enough production go in opposite directions. Instead in the third example we consider a scenario where it is possible to decouple production from decay.

4.1 Heavy neutral leptons

The existence of heavy neutral leptons (HNL) is a generic prediction of Type-I seesaw models of neutrino masses [5–8]. The model is the simple extension of the SM with at least two heavy Majorana singlets N_j . In the basis where the Majorana mass terms are diagonal, the Lagrangian reads:

$$\mathcal{L}_N = \mathcal{L}_{\text{SM}} + \sum_j i \bar{N}_j \gamma^\mu \partial_\mu N_j - \left(Y_{\alpha j} \bar{L}_\alpha \tilde{\Phi} N_j + \frac{m_{N_j}}{2} \bar{N}_j N_j^c \right), \quad (4.1)$$

where $\tilde{\Phi} \equiv i\sigma_2 \Phi^*$ is the complex conjugate of the Higgs field Φ , L_α is the SM lepton doublet with flavor α , Y_j is a Yukawa coupling and m_{N_j} is the Majorana mass of the singlet N_j , which in principle is a free parameter of the model. After spontaneous symmetry breaking, the heavy Majorana states mix with the standard neutrinos resulting in a spectrum of almost standard light states (with masses $m_\nu \propto \frac{(Yv)^2}{m_N}$) which correspond to the light neutrinos, and almost singlet heavy states corresponding to the HNL (with masses $\propto m_N$). In the following, the full leptonic mixing matrix will be denoted as U and the mixing between the charged leptons and the heavy states is given by $|U_{\alpha j}|^2$, which naively scales like $\left(\frac{Yv}{m_N}\right)^2$. It is through this mixing that the heavy singlets could be produced either through CC or neutral-current (NC) processes and also how they would decay back to SM particles.

A priori, the Majorana masses of the HNL are free parameters of the theory and therefore have to be probed experimentally. The GeV range is interesting from the theoretical point of view, since it has been shown that models with HNL at the GeV scale could explain neutrino masses and mixing parameters, as well as the matter-antimatter asymmetry of the Universe [9, 10] without conflicting cosmological bounds. The search for GeV scale HNL is a very active field and several future experiments such as SHIP [45], FCCee [46], or DUNE [47] have the potential to improve significantly over present bounds [13, 14]. On the phenomenological front, new avenues to constrain these models are being proposed and/or further explored using current and past data [15, 16, 48–52].

Here we consider instead the production of HNL in atmospheric showers, where their dominant production mechanism is through meson decays. Depending on their mass, the leading production channel is π , K , $D_{(s)}$ -meson decays (or in the decays of even heavier resonances, such as B mesons). Particularly interesting is the mass range slightly above the kaon mass, where existing laboratory bounds are weaker [13–16]. We have therefore considered the production via D - and D_s -meson decays, as well as from τ decays if one considers the possibility that the production is dominated by the poorly constrained coupling to taus. The production in the lower-mass range from kaon and pion decays will be considered elsewhere. For simplicity, hereafter we use the notation $U_\alpha \equiv U_{\alpha j}$ ($\alpha = e, \mu, \tau$) to refer to the elements of the leptonic mixing matrix that control both the production and decay of the heavy state N .

For $D_{(s)}$ meson decays, and for HNL between 0.5 and 1.5 GeV the dominant decay is two-body: $D^\pm \rightarrow N l_\alpha^\pm$, where the flavour of the lepton $\alpha = e, \mu$ (note that, if $m_N \geq m_K$, this decay with $\alpha = \tau$ is kinematically forbidden). The N fluxes are obtained from

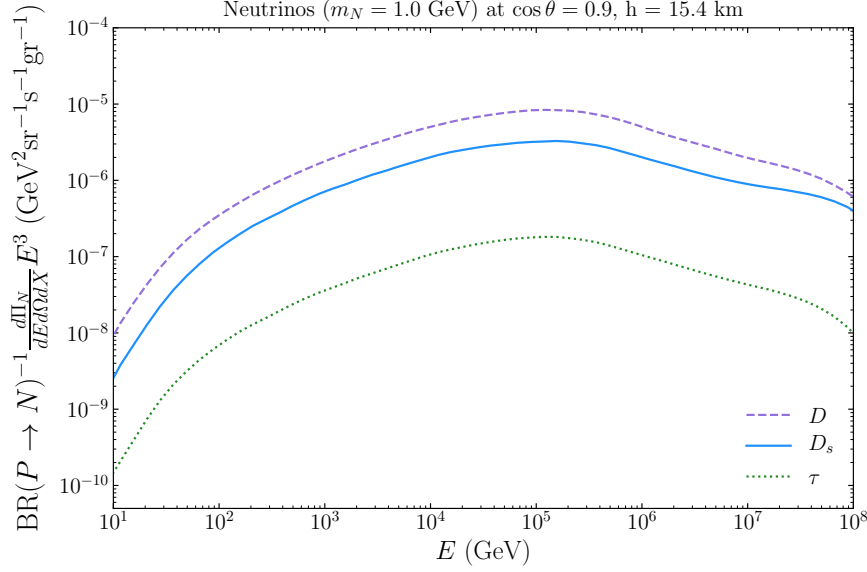


Figure 4. Production profile of HNL in the atmosphere, shown for D, D_s and τ decays separately, in units of the corresponding branching ratios and for $m_N = 1$ GeV. The profile is shown at approximately 15 km from the Earth's surface, since this is where the parent meson fluxes are expected to be maximal, see figure 2, and for $\cos \theta = 0.9$ (we expect the signal to be largest as $\cos \theta \rightarrow 1$).

eqs. (2.3), (2.4) and (2.5). The production profile of N from D and D_s decays are shown in figure 4 in units of the corresponding branching ratios into this channel.

In the case of taus, the decays to produce N are necessarily three-body. Assuming that the dominant mixing is to τ leptons and the mass of the HNL is in the GeV range, the dominant decay channels are $\tau^\pm \rightarrow N\nu_\alpha l_\alpha^\pm$, where $l_\alpha \equiv e, \mu$. In this case it is more convenient to perform the change of variables $E \rightarrow z \equiv E/\gamma_\tau$, where $\gamma_\tau = E_\tau/m_\tau$ is the boost factor of the τ . Equation (2.10) then becomes

$$\text{Br}(\tau \rightarrow N\nu_\alpha l_\alpha)^{-1} \frac{dn(\tau \rightarrow N\nu_\alpha l_\alpha; E, E_\tau)}{dE} = \frac{1}{\gamma_\tau \Gamma(\tau \rightarrow N\nu_\alpha l_\alpha)} \frac{d\Gamma(\tau \rightarrow N\nu_\alpha l_\alpha)}{dz}. \quad (4.2)$$

The differential decay width in the rest frame of the τ lepton can be found in ref. [53], or in appendix B of ref. [54]. Neglecting all lepton masses except those of the τ and N , after boosting to the laboratory frame we get¹

$$\begin{aligned} & \frac{1}{\Gamma(\tau \rightarrow N\nu_\alpha l_\alpha)} \frac{d\Gamma(\tau \rightarrow N\nu_\alpha l_\alpha)}{dz} \\ &= \frac{1}{72z^3 X_\tau} (m_\tau - z) \left[m_N^6 (4m_\tau^2 - 5m_\tau z - 5z^2) - 9m_N^4 m_\tau^2 z (m_\tau - 3z) \right. \\ & \quad \left. + 9m_N^2 m_\tau^2 z^3 (z - 3m_\tau) + m_\tau^3 z^3 (5m_\tau^2 + 5m_\tau z - 4z^2) \right] \\ & \quad + \mathcal{O}(1 - \beta_\tau), \end{aligned} \quad (4.3)$$

where

$$X_\tau \equiv \frac{-m_N^8 + 8m_N^6 m_\tau^2 - 24m_N^4 m_\tau^4 \log\left(\frac{m_N}{m_\tau}\right) - 8m_N^2 m_\tau^6 + m_\tau^8}{24m_\tau}. \quad (4.4)$$

¹Here we assume $\gamma_\tau m_N \geq E_{\text{max}}$, where E_{max} can be found in eq. (2.12).

The contribution to the flux from τ decays is also shown in figure 4 in units of the $\tau \rightarrow N l_\alpha \nu_\alpha$ branching ratio.

The partial decay widths of D , D_s and τ to N can be obtained from refs. [53], [54], and [55]. For the D and D_s mesons, the total branching ratio for two-body decays with a HNL in the final state is obtained adding the contributions from $D_{(s)} \rightarrow Ne$ and $D_{(s)} \rightarrow N\mu$. Assuming that the partial decay width of the parent meson to N is very small compared to its SM decay width, the branching ratio reads

$$\text{Br}(D_{(s)} \rightarrow N) \simeq \frac{G_F^2 f_{D_{(s)}}^2 |V_{D_{(s)}}|^2 m_{D_{(s)}}^3}{8\pi \Gamma_{D_{(s)}}} \sum_{\alpha=e,\mu} |U_\alpha|^2 [y_N^2 + y_{l_\alpha}^2 - (y_N^2 - y_{l_\alpha}^2)^2] \sqrt{\lambda(1, y_N^2, y_{l_\alpha}^2)}, \quad (4.5)$$

where $\Gamma_{D_{(s)}}$ is the total decay width of the $D_{(s)}$ meson (taken from ref. [56]), $y_i \equiv m_i/m_{D_{(s)}}$ and λ is defined in eq. (2.6). In eq. (4.5) G_F is the Fermi constant, $f_{D_{(s)}}$ is the decay constant of the parent meson and $V_{D_{(s)}}$ is the mixing matrix element in the CKM matrix that participates in the decay vertex. The values used for the meson masses and decay constants ($f_D = 212 \text{ MeV}$; $f_{D_s} = 249 \text{ MeV}$) have been taken from ref. [56]. For the CKM matrix elements, we use $V_D \equiv V_{cd} = 0.22$ and $V_{D_s} \equiv V_{cs} = 0.995$. In the case of decays from taus, we assume that the τ branching ratio into N is dominated by the value of U_τ since it is less constrained by other experiments. In this case, we get:

$$\begin{aligned} \text{Br}(\tau^\pm \rightarrow N) &= \sum_{\alpha=e,\mu} \text{Br}(\tau^\pm \rightarrow N l_\alpha^\pm \nu_\alpha) \\ &= 2 \frac{|U_\tau|^2}{192\pi^3} \frac{G_F^2 m_\tau^5}{\Gamma_\tau} [1 - 8y_N^2 - 24y_N^4 \log(y_N) + 8y_N^6 - y_N^8], \end{aligned} \quad (4.6)$$

where, again in this case, for the total value of the width of the τ we use the SM value of Γ_τ (i.e. ignoring the small contribution of the HNL).

The expected number of decays inside the volume of the detector will also depend on the lifetime of the N , as outlined in section 3. The total width is computed adding the partial widths into two-body decay channels (that is, into charged mesons and charged leptons, or neutral mesons and neutrinos), the three-body decay channels into charged leptons and neutrinos, and the three-body invisible decay into three light neutrinos. The partial widths for these decays can be found in several references [13, 47, 54, 57] although they do not all agree. We have re-computed the two-body decay widths into mesons and leptons in the effective Lagrangian at low energies, and found good agreement with the results from the recent review in ref. [57]. The total width can be written as

$$\Gamma_{\text{tot}} = \sum_{\alpha=e,\mu,\tau} \Gamma_\alpha, \quad (4.7)$$

where the dependence on the mixing matrix elements goes as $\Gamma_\alpha \propto |U_\alpha|^2$. In figure 5 we show the ratios $|U_\alpha|^{-2} \Gamma_\alpha$ as a function of the N mass. From these, it is easy to reconstruct the total width and lifetime for any given values of the mixing matrix elements $|U_\alpha|^2$.

As already outlined, in this model both the production and decay rates (that is, the lifetime of the LLP) are controlled by the same set of parameters and therefore are highly

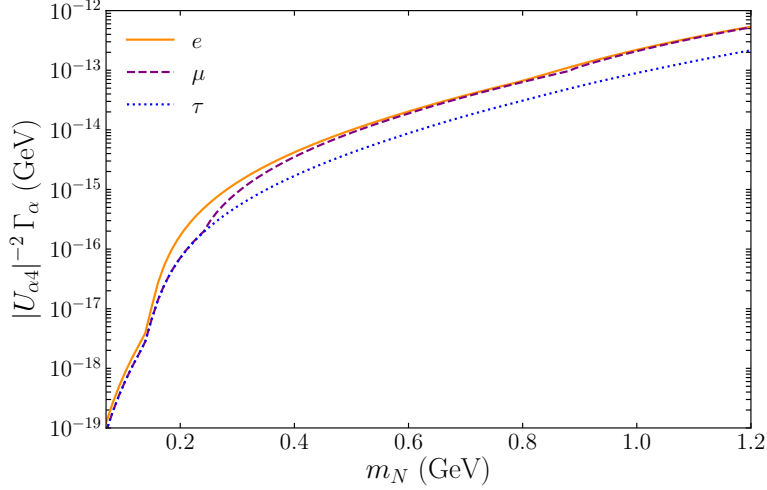


Figure 5. Ratios $|U_\alpha|^{-2}\Gamma_\alpha$ for $\alpha = e, \mu, \tau$ as a function of m_N .

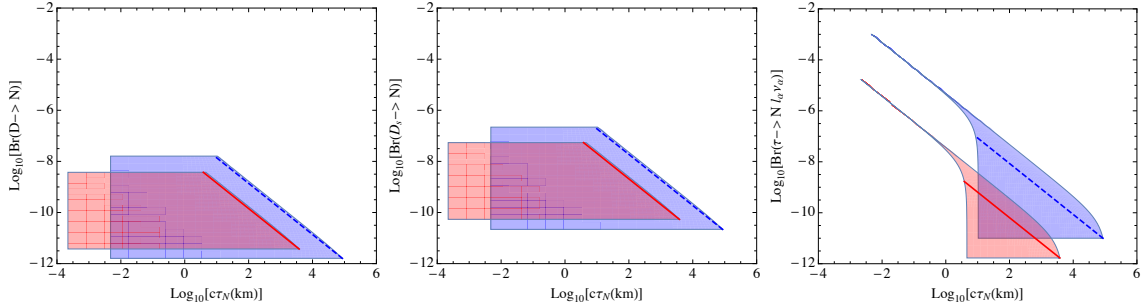


Figure 6. Allowed ranges for $\text{Br}(D \rightarrow N)$ (left panel) and $\text{Br}(D_s \rightarrow N)$ (middle panel) and $\text{Br}(\tau^\pm \rightarrow N l_\alpha^\pm \nu_\alpha)$ (right panel) versus $c\tau_N$, for $m_N = 0.5 \text{ GeV}$ (dashed line, blue region) and 1 GeV (solid line, red region). The lines correspond to $10^{-10} \leq |U_e|^2 = |U_\mu|^2 = |U_\tau|^2 \leq 10^{-6}(10^{-7})$ for the lighter (heavier) mass. The ranges correspond to $10^{-10} \leq |U_e|^2 = |U_\mu|^2 \leq 10^{-6}(10^{-7})$ and $10^{-10} \leq |U_\tau|^2 \leq 10^{-2}(10^{-3})$.

correlated. Figures 6 show the regions allowed on the $\text{Br}(P \rightarrow N)$ vs $c\tau_N$ plane (for $P = D, D_s, \tau$) allowing the mixing matrix elements to vary within the presently allowed regions [13], for two values of the HNL mass. As can be seen from both figures, in this model it is possible to achieve values of the lifetime in the right ballpark needed to obtain a signal in neutrino detectors ($c\tau \sim \mathcal{O}(10)$ km, after boosting to the lab frame), although at the price of very small production branching ratios.

Finally, the number of events will also be proportional to the branching ratio of the decays leading to a cascade-like signature (at IC), or leading to e-like events (at SK). In practice, this amounts to adding the branching ratios into decay channels with either electrons, taus or hadronic resonances in the final state:

$$\begin{aligned} \text{Br}(\text{CCe-like}) = & \sum_M \sum_{\alpha=e} \text{Br}(N \rightarrow M^- l_\alpha^+) + \sum_M \text{Br}(N \rightarrow M^0 \nu) \\ & + \sum_{\alpha=e, \tau} \text{Br}(N \rightarrow \nu l_\alpha^- l_\alpha^+) + \sum_{\alpha, \beta=e, \tau} \text{Br}(N \rightarrow \nu l_\alpha^- l_\beta^+), \end{aligned} \quad (4.8)$$

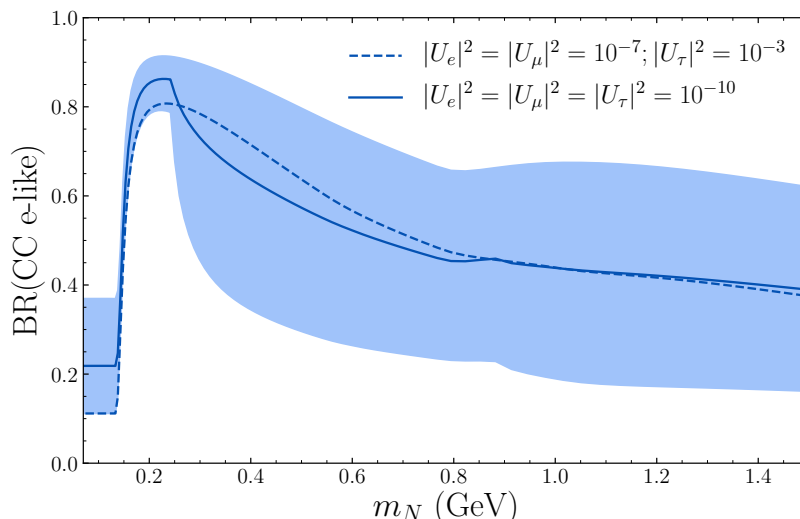


Figure 7. Value of the branching ratio of the HNL into e-like final states, defined in eq. (4.8), as a function of its mass. The lines correspond to fixed values of the mixing with the light neutrinos as indicated by the labels, while the shaded areas show the range of branching ratios accessible by varying independently the mixing matrix elements within the following limits: $10^{-10} < |U_e|^2 < 10^{-7}$, $10^{-10} < |U_\mu|^2 < 10^{-7}$ and $10^{-10} < |U_\tau|^2 < 10^{-4}$. For illustration purposes, we have chosen the same ranges for the mixing matrix elements as in figure 6. However, note that for specific mass ranges some of these values are excluded by laboratory constraints [58–61] (see also refs. [13, 16]).

where the processes involving light neutrinos include the diagrams for all active neutrino flavors, and the sum over M includes all charged or neutral pseudoscalar and vector mesons below the mass of the N ($\pi^\pm, K^\pm, \rho^\pm, \pi^0, \eta, \eta', \rho^0$, etc). The result is shown in figure 7 as a function of the mass of the HNL.

4.2 Dark photons

Extensions of the Standard Model with an extra secluded U(1) gauge boson are ubiquitous in the BSM landscape. We first consider a model where the dark photon V is coupled to the visible sector via kinetic mixing [12]:

$$\mathcal{L}_V = \mathcal{L}_{\text{SM}} - \frac{1}{4}V_{\mu\nu}V^{\mu\nu} + \frac{\epsilon}{2}B_{\mu\nu}V^{\mu\nu} - \frac{1}{2}m_V^2V_\mu V^\mu, \quad (4.9)$$

where $B_{\mu\nu}$ is the hypercharge field strength tensor and $V_{\mu\nu}$ is analogously defined for the dark U(1), which is assumed spontaneously broken. The Lagrangian in eq. (4.9) implies that the dark photon couples universally to all charged particles, like the SM photon, although with a coupling that is reduced by the factor ϵ . An enormous amount of work has been done in recent years to derive bounds on this scenario from beam dump experiments, e^+e^- and pp colliders, neutrino scattering and other intensity frontier experiments. A recent summary of present bounds can be found in [62].

The mechanisms leading to the production of dark photons in the atmosphere are the same as in a proton beam dump, which has been extensively studied in the literature, see e.g., refs. [34, 63, 64]. For $m_V \leq m_{\pi_0}(m_{\eta'})$, the dominant production channel is the

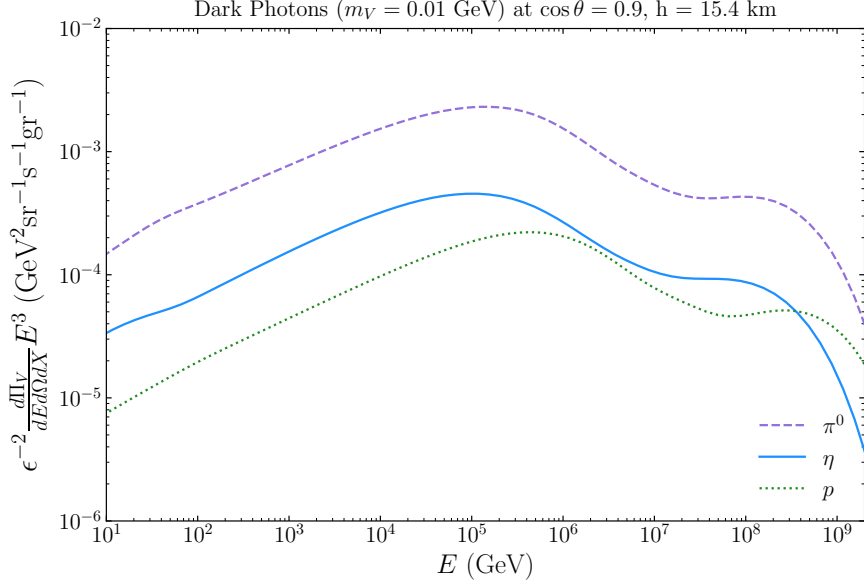


Figure 8. Production profile of dark photons produced from π_0, η decay and proton bremsstrahlung, for $m_V = 10^{-2}$ GeV. The profile is shown at approximately 15 km from the Earth’s surface, since this is where the parent meson fluxes are expected to be maximal, see figure 2, and for $\cos \theta = 0.9$ since we expect the signal to be largest as $\cos \theta \rightarrow 1$.

two-body decay $\pi_0(\eta') \rightarrow \gamma V$. For heavier masses and $m_V \leq m_p$, it can be produced via bremsstrahlung. The flux of V from neutral meson decays can be obtained from eqs. (2.3), (2.4) and (2.5), with the production branching ratios [65]:

$$\text{Br}(P \rightarrow V\gamma) \simeq \epsilon^2 \left(1 - \frac{m_V^2}{m_P^2}\right) \text{Br}(P \rightarrow \gamma\gamma), \quad (4.10)$$

for $P = \pi_0, \eta$. The flux from bremsstrahlung is obtained from eqs. (2.14) and (2.15). The production profile of dark photons from these processes in units of ϵ^{-2} are shown in figure 8.

The next ingredient we need in order to compute the expected number of decays in the detector is the lifetime of the dark photon. The partial decay widths into leptonic and hadronic channels can be found, for example, in refs. [34, 63, 64]:

$$\Gamma(V \rightarrow l^+ l^-) = \frac{1}{3} \epsilon^2 \alpha_{\text{QED}} m_V \sqrt{1 - \frac{4m_l^2}{m_V^2}} \left(1 + \frac{2m_l^2}{m_V^2}\right), \quad (4.11)$$

$$\Gamma(V \rightarrow \text{hadrons}) = \frac{1}{3} \epsilon^2 \alpha_{\text{QED}} m_V \frac{\sigma(e^+ e^- \rightarrow \text{hadrons})}{\sigma(e^+ e^- \rightarrow \mu^+ \mu^-)}, \quad (4.12)$$

where l stands for a SM charged lepton and $\alpha_{\text{QED}} \equiv e^2/(4\pi)$ stands for the SM fine structure constant. Adding up all possible decay channels below the hadronic threshold, and neglecting the electron mass, we get:

$$c\tau(\text{km}) \simeq \frac{8.1 \cdot 10^{-17} \text{ GeV}}{\epsilon^2} \frac{1}{m_V} \left[1 + \sqrt{1 - 4 \frac{m_\mu^2}{m_V^2}} \left(1 + 2 \frac{m_\mu^2}{m_V^2}\right)\right]^{-1}, \quad 2m_\mu \leq m_V \leq 2m_\pi. \quad (4.13)$$

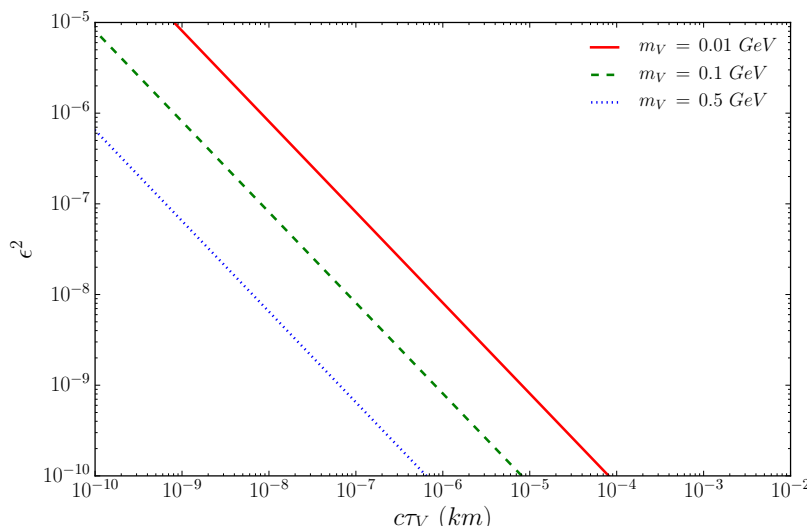


Figure 9. Correlation between the value of ϵ and the lifetime of the dark photon in its rest frame, $c\tau_V$, for three values of m_V : 0.01 GeV (solid), 0.1 GeV (dashed) and 0.5 GeV (dotted).

Conversely, above the hadronic production threshold we can approximate the lifetime by the fitted formula:

$$c\tau(\text{km}) \approx \frac{3 \cdot 10^{-17} \text{ GeV}}{\epsilon^2} \frac{1}{m_V} \left[1 + 2.54 \frac{m_V}{\text{GeV}} - 4.76 \left(\frac{m_V}{\text{GeV}} \right)^2 \right], \quad 2m_\pi \leq m_V \leq m_\rho. \quad (4.14)$$

Again in this case, both the production rate and the decay length depend on just two parameters: ϵ and m_V . Figure 9 shows the correlation of ϵ and the lifetime of the dark photon, for three different values of m_V in the range considered in this work.

Finally, in our sensitivity calculations we focus on decays of the LLP leading to an electron-like signal in the detectors (i.e. decays with electrons or hadrons in the final state). Therefore, in order to compute the number of events in this sample we need the corresponding branching ratio. The result is shown in figure 10 as a function of the mass of the dark photon.

4.3 Heavy neutral leptons and $U(1)_{B-L}$ gauge symmetry

We have finally considered an extension of the SM that includes both HNL and a new $U(1)$ gauge symmetry associated to the $B - L$ number. Models with HNL and extra $U(1)$ have been recently discussed in the context of the MiniBoone anomaly [51, 66–68], as well as in the context of explaining the matter-antimatter asymmetry or/and the observed dark matter abundance in the Universe [69]. The relevant terms in the Lagrangian are

$$\mathcal{L} = \mathcal{L}_N - \frac{1}{4} V_{\mu\nu} V^{\mu\nu} - \frac{1}{2} m_V^2 V_\mu V^\mu + g_{B-L} \left(\sum_f Q_{B-L}^f \bar{f} \gamma^\mu f V_\mu - \bar{N} \gamma^\mu N V_\mu \right), \quad (4.15)$$

where g_{B-L} is the gauge coupling of the dark photon, and the sum runs over all standard model fermions, which are charged under this new symmetry ($Q_{B-L}^f = 1/3$ for quarks, and

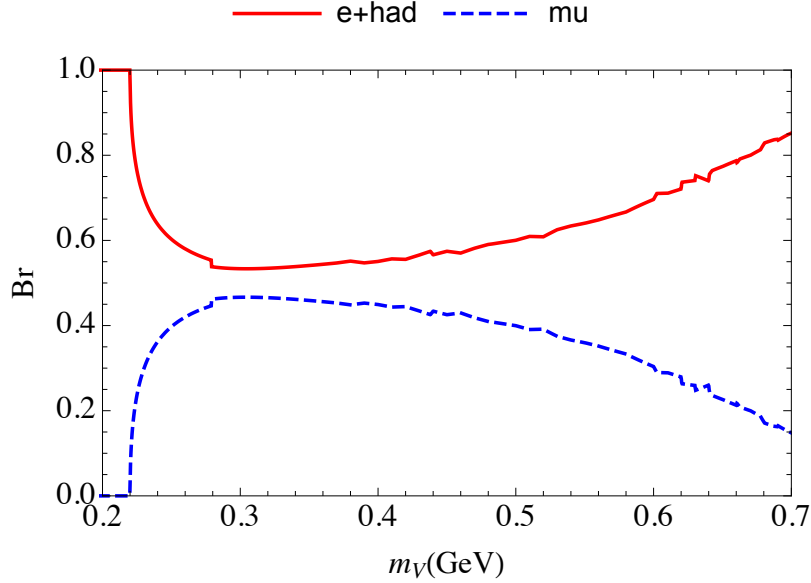


Figure 10. Branching ratios of the dark photon to electron/hadronic or muon decay channels as a function of the mass.

$Q_{B-L}^f = -1$ for leptons). In this model, the origin of the B-L gauge boson mass typically requires a more complex hidden sector. However, our approach is phenomenological and therefore we will not address the possible origin of the mass term which is beyond the scope of this work. Finally, note that in principle a kinetic mixing term could also be added to eq. (4.15). If the kinetic mixing is large, the phenomenological consequences of the model reduce to the dark photon case, which was already discussed in section 4.2. Here we focus on the opposite limit instead, when the kinetic mixing term is negligible and the coupling dominates, since it leads to a different phenomenology. While a kinetic mixing term is always generated at loop level, this is expected to be subleading.

An interesting feature of this model is that the production of the N is dominated by the decays of the dark photon $V \rightarrow NN$, provided $m_N \leq m_V/2$, while its decay is controlled by the mixing with the light neutrinos. We might expect in this case to have an enhanced production if g_{B-L} is not too small, while the N might be very long-lived as shown in figure 6. Present upper limits [62, 70, 71] on g_{B-L}^2 are around $\sim 10^{-7}$ in the mass range of interest, around $m_V \sim \mathcal{O}(\text{GeV})$. From figure 6 we see that $c\tau_N$ in the range $[10^{-2}, 10^5]$ km for $m_V \sim 0.5 \text{ GeV}$ are allowed.

Neglecting the decay of the dark photon, the flux of V can be obtained as in section 4.2 simply replacing [62]:

$$\alpha_{\text{QED}}\epsilon^2 \rightarrow \alpha_{B-L} \equiv \frac{g_{B-L}^2}{4\pi}. \quad (4.16)$$

However, in the $B-L$ model the dark photon decays very promptly. Its lifetime (for $2m_N \leq m_V$, and $2m_\mu \leq m_V \leq m_\omega$) is approximately given by

$$c\tau_V(\text{km}) \simeq \frac{0.74 \cdot 10^{-17} \text{ GeV}}{g_{B-L}^2} \frac{1}{m_V} \left[\frac{5}{2} + \sqrt{1 - 4 \frac{m_\mu^2}{m_V^2}} \left(1 + 2 \frac{m_\mu^2}{m_V^2} \right) + \frac{1}{2} \left(1 - \frac{4m_N^2}{m_V^2} \right)^{3/2} \right]^{-1}. \quad (4.17)$$

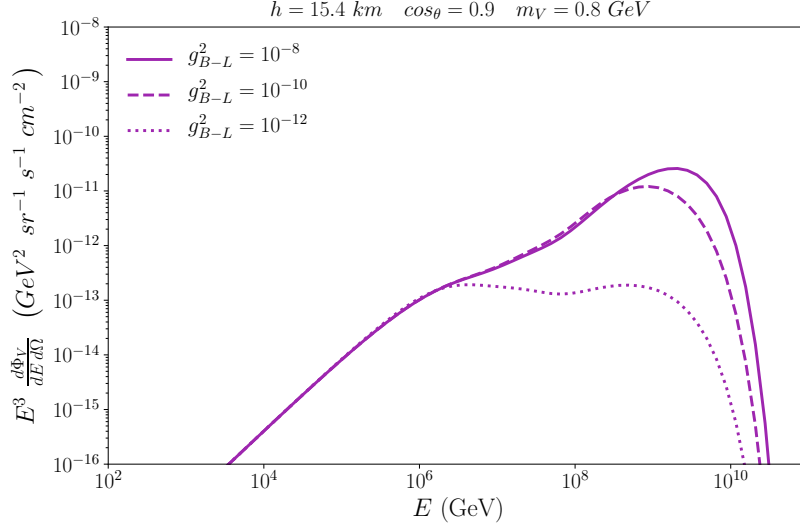


Figure 11. V fluxes in the $B - L$ model with $m_V = 0.8 \text{ GeV}$ and for three values of g_{B-L}^2 .

Therefore, decay effects should be accounted for in the calculation of the flux, following the same arguments as in appendix B. It is easy to show that, if we denote by $\tilde{\phi}_V$ the dark photon flux obtained without accounting for the decay, then the solution to the cascade equation including decay effects is given by:

$$\phi_V(\ell) = \int_{\ell}^{\infty} d\ell' e^{-\frac{\ell' - \ell}{\ell_{\text{decay}}}} \frac{d\tilde{\phi}_V}{d\ell'}, \quad (4.18)$$

where $\ell_{\text{decay}} = \gamma_V \beta_V c \tau_V$. The flux of dark photons taking into account the decay is shown in figure 11 for a three different values of g_{B-L} in the experimentally allowed range. Interestingly, while at high energies there is clearly a suppression of the dark photon flux, which depends on the value of g_{B-L} , at low energies the production remains practically independent from the value of g_{B-L} . This can be understood as follows. In this part of the spectrum the lifetime of the dark photon is very short and the exponential in eq. (4.18) goes to zero very rapidly, except for values of ℓ' such that $\Delta\ell = \ell' - \ell \propto c\tau_V$. Thus, at low energies the resulting flux will take the form

$$\phi_V \propto c\tau_V \frac{d\tilde{\phi}_V}{d\ell}, \quad (4.19)$$

and since $c\tau_V \propto g_{B-L}^{-2}$ this effectively renders the result independent of g_{B-L} at low energies.

From these fluxes, we can get the production profile of N as in eq. (2.3), where now the parent particle P is the dark photon instead of SM mesons. The decay $V \rightarrow NN$ is two-body and therefore the decay distribution is the same as in eqs. (2.4) and (2.5). Neglecting the electron mass, its partial decay width reads:

$$\Gamma(V \rightarrow NN) \simeq \frac{1}{2} \left(1 - \frac{4m_N^2}{m_V^2} \right)^{3/2} \Gamma(V \rightarrow e^+ e^-). \quad (4.20)$$

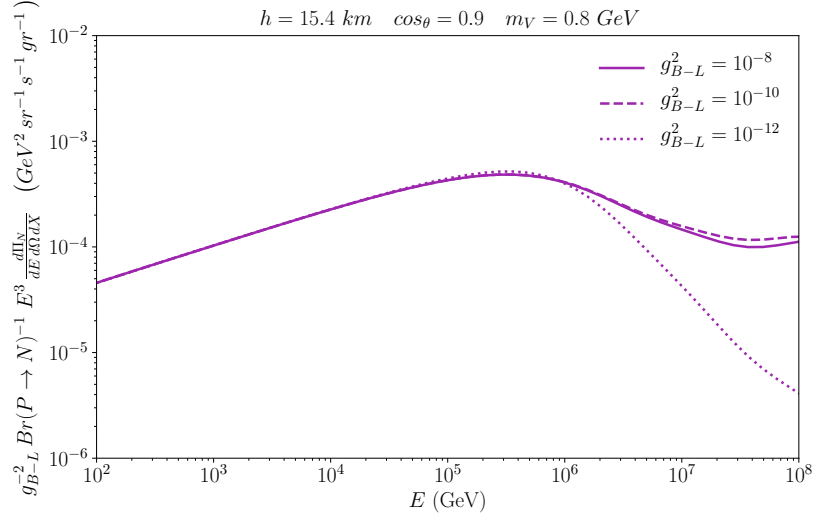


Figure 12. N production from V decays in the $B-L$ model with $m_V = 0.8$ GeV, $m_N = 0.35$ GeV and for three values of g_{B-L}^2 . Note that because two N are produced in each decay we define $\text{Br}(V \rightarrow N) = 2\text{Br}(V \rightarrow NN)$.

Finally, the branching ratio in this channel (in the mass region of interest) can be well approximated by:

$$\text{Br}(V \rightarrow NN) \simeq \frac{1}{2} \left(1 - \frac{4m_N^2}{m_V^2}\right)^{3/2} \left[\frac{5}{2} + \sqrt{1 - 4\frac{m_\mu^2}{m_V^2}} \left(1 + 2\frac{m_\mu^2}{m_V^2}\right) + \frac{1}{2} \left(1 - \frac{4m_N^2}{m_V^2}\right)^{3/2} \right]^{-1}, \quad (4.21)$$

since the decays into mesons are suppressed.

The resulting production profile of N from dark photon decays is shown in figure 12. Since the production and decay processes are effectively decoupled, the decay length of the N (which is controlled by the mixing) can now take a wide range of values (as shown in figure 6) without affecting the production rates.

5 Results

In this section we present our numerical results for the sensitivity to the scenarios described in section 4. Using the data sets described in section 3, a χ^2 fit to the data is performed. A Poissonian χ^2 function has been used, defined as:

$$\chi^2(\text{LLP}) = 2 \sum_{i,\alpha} \left(N_i^\alpha + B_i^\alpha - n_i^\alpha + n_i^\alpha \log \left(\frac{n_i^\alpha}{N_i^\alpha + B_i^\alpha} \right) \right), \quad (5.1)$$

where the sum runs over the energy (angular) bins for the IC (SK) detector analyses, and we use $\alpha = \text{cascade(e-like)}$ for IC (SK). Here, n_i^α stands for the data observed in each bin while N_i^α is the predicted number of signal events and B_i^α is the background prediction, which includes the predicted number of atmospheric neutrino events in the SM (plus the contribution from the fitted astrophysical neutrino flux, in the case of IC).

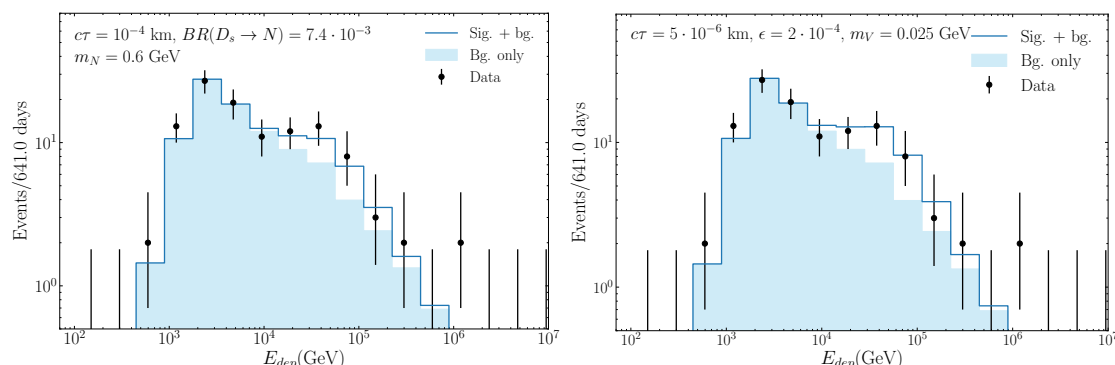


Figure 13. Expected number of decays for the IC Medium Energy Starting Events (MESE) sample, as a function of the deposited energy in the detector and for the down-going sample only ($\cos \theta > 0$), for two of the scenarios considered in this work. The shaded histograms show the background prediction, including a fitted distribution to the astrophysical neutrino signal, the cosmic muon background, and atmospheric neutrino events. The black dots correspond to the observed data with error bars, as in ref. [39]. In the left panel, the signal has been computed for a HNL with a mass $m_N = 0.6$ GeV, a $c\tau_N = 10^{-4}$ km and $\text{Br}(D_s \rightarrow N) \text{Br}(N \rightarrow \text{CC e-like}) = 7.4 \cdot 10^{-3}$. In the right panel, the signal has been computed for a dark photon with a mass $m_V = 25$ MeV for $\epsilon = 2 \cdot 10^{-4}$ and an (uncorrelated) lifetime $c\tau_V = 5 \cdot 10^{-6}$ km.

A comment regarding the treatment of the astrophysical neutrino background for IC is in order, since the distribution of astrophysical neutrinos is obtained from a fit to the same data that we use in order to derive a limit to LLP models. Although this may seem inconsistent, it should be noted that due to the nature of the decay and the spectra of the parent mesons in the atmosphere (which obey a power law), at IC we expect most of our sensitivity to come from events observed at energies below 10 TeV, where the contribution from the astrophysical neutrinos is subdominant. Moreover, from the comparison between the left and right panels in figure 8 in ref. [39], it seems that the fit to the astrophysical neutrino background is mostly driven by the data from the northern sky, whereas we expect LLP decays to contribute mostly to the southern sky data set, which is the only data we use. Therefore, although a more detailed analysis by the experimental collaboration would be necessary to do this analysis properly, we believe the results would not be very different. Of course, given the very characteristic zenith angle dependence of the LLP signal, a greater sensitivity is expected if 2D binned data (using both energy and angular information) were to be used instead.

For illustration, figure 13 shows the expected number of signal events for HNL (left panel) and dark photons (right panel), for an assumed value of the production branching ratio of the HNL and its lifetime. The shaded histograms show the background prediction, while the solid lines show the expected signal plus background event rates per bin. For reference, the black points show the observed data, as given in ref. [39].

The zenith angle distribution for the neutrino background and LLP is signal events in Super-Kamiokande is shown in figure 14. As expected the signal is peaked at small zenith angles, that is, most of the decaying LLP are expected to come from trajectories entering the detector from above.

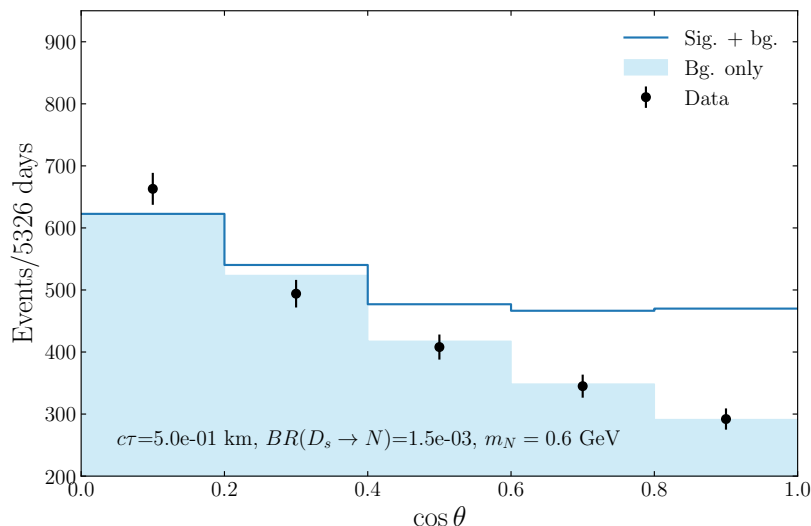


Figure 14. Expected number of decays in SK, as a function of the zenith angle. The shaded histograms show the atmospheric background prediction, while the solid histograms show the signal plus neutrino background. The signal in this case has been computed for a HNL with a mass $m_N = 0.6$ GeV, a $c\tau_N = 0.5$ km and $\text{Br}(D_s \rightarrow N)\text{Br}(N \rightarrow \text{CCe} - \text{like}) = 1.5 \cdot 10^{-3}$. The data points, extracted from ref. [44], are shown with statistical error bars for reference.

Before moving on to the discussion of our results, we should note that no systematic uncertainties have been included in our calculations, since this is an extremely challenging task using only publicly available information. Eventually, a detailed study including a proper implementation of systematic uncertainties and detector reconstruction effects should be performed by the experimental collaborations themselves.

5.1 Exclusion limits for IceCube and Super-Kamiokande

Following the procedure described above, we proceed to derive exclusion limits for the HNL and dark photon LLPs, for both IC and SK. In doing so, we consider that the data is chi-squared distributed with n degrees of freedom (d.o.f.), n being the number of bins in the data (that is, 5 bins in $\cos \theta$ in the case of SK, as opposed to 15 energy bins in the case of IC). Our exclusion limits thus indicate the region in parameter space where the χ^2 value exceeds the corresponding one at 90% CL, regardless of where the best-fit point lies in the parameter space. Our results are shown in figure 15 for the HNL, assuming $m_N = 1$ GeV, and in figure 16 for the dark photons, for three different values of m_V as indicated in the legend. In both figures, thick lines correspond to the IC limits, while thin lines indicate the SK result.

For both IC and SK, the best limits are reached when the decay length in the lab frame is $c\tau_{\text{lab}} \sim \mathcal{O}(10)$ km, as expected from naive arguments. However, the two experiments observe events in very different energy regimes: $\mathcal{O}(\text{TeV})$ in the case of IC, while $\mathcal{O}(\text{GeV})$ in SK, and therefore the boost factor (which also depends on the mass of the LLP) will be very different. As a result, their sensitivities are highly complementary and explore different regions of $c\tau$. For IC, the best limit is reached for $c\tau \sim \mathcal{O}(10^{-3} - 10^{-2})$ km in

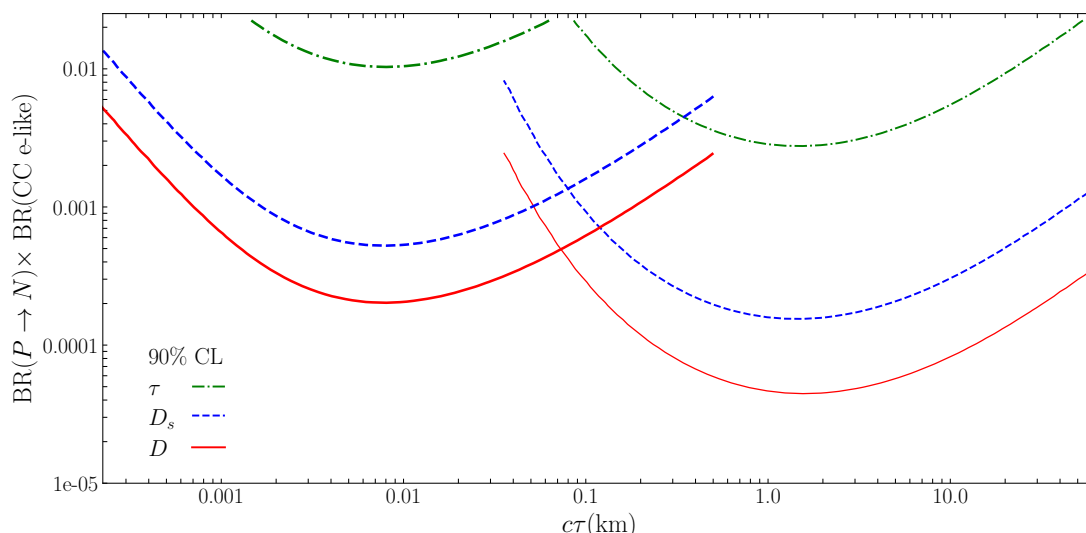


Figure 15. Limits on HNL of $m_V = 1$ GeV from IceCube (thick lines) and SuperKamiokande (thin lines) on the $BR(P \rightarrow N) \times BR(N \rightarrow CC\text{-e like})$ vs $c\tau$ plane including production from the parent particles $P = D$ (solid), D_s (dashed) and τ (dash-dotted).

the case of LLP with masses $m \sim 0.5 - 1$ GeV, while for SK the best limits are obtained for $c\tau \sim \mathcal{O}(1 - 10)$ km.

In the case of the HNL (figure 15), the solid red, dashed blue, and dot-dashed green lines show our results assuming that the LLP is mainly produced from D , D_s or τ decays. Although our results are shown for $m_N = 1$ GeV, the regions do not change significantly for other values of m_N between 0.5 and 1.5 GeV. As can be seen from this figure, the limits obtained for the two experiments are very similar, although slightly better for SK. However, in both cases the limits cannot probe the allowed regions of parameter space once we take into account the correlation between production and decay, as shown in figures 6. Therefore we conclude that, in the minimal scenario described in section 4.1, these limits are not competitive (although this may not be the case in non-minimal BSM scenarios where the production and decay may be uncorrelated).

The exclusion limits for the dark photon scenario (section 4.2) are shown in figure 16. In this case our results include all production mechanisms kinematically available for the production of the dark photon: π^0 decays, η decays, and p bremsstrahlung. The solid red, dashed green and dotted blue lines indicate the results obtained for three different values of the dark photon mass, as indicated in the legend. Again in this case, as for the HNL scenario, we find that this search cannot probe the region of parameter space shown in figure 9 for correlated production and decay rates. However, the limits might be useful in the context of more complex models where production and decay are uncorrelated. In particular, it is noticeable that the limits for SK are three orders of magnitude better than for IC. This is because, in the case of light mesons and protons, the flux follows a harder power law than for $D_{(s)}$ mesons and taus, and therefore reducing the energy of the events detected leads to huge enhancement in the expected number of events.

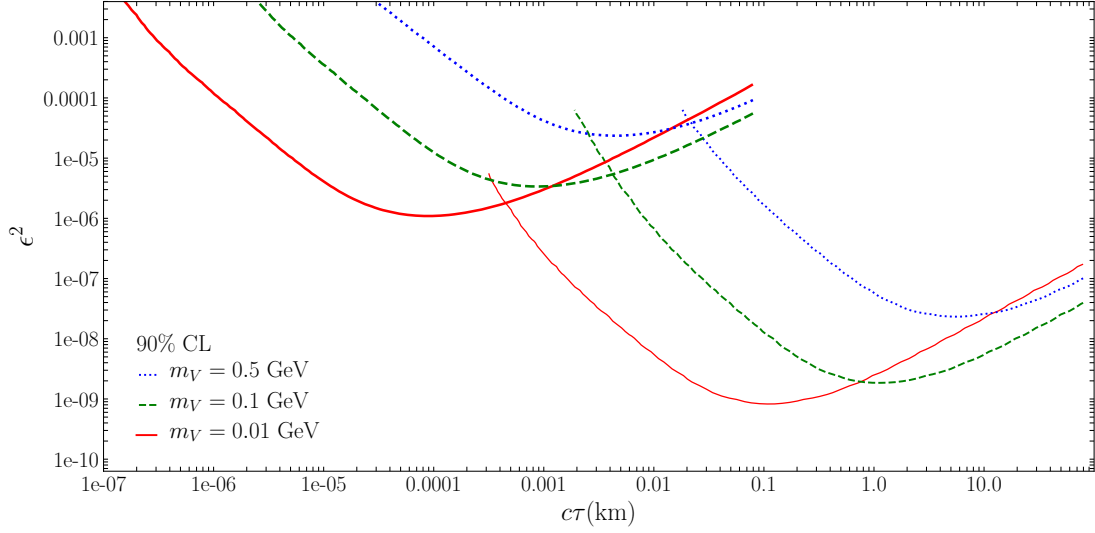


Figure 16. Limits on dark photons decays from IceCube (thick lines) and Super-Kamiokande (thin lines) on the ϵ^2 vs $c\tau_V$ plane including production from π_0, η decay and bremsstrahlung for $m_V = 0.01$ GeV (solid), 0.1 GeV (dashed) and 0.5 GeV (dotted).

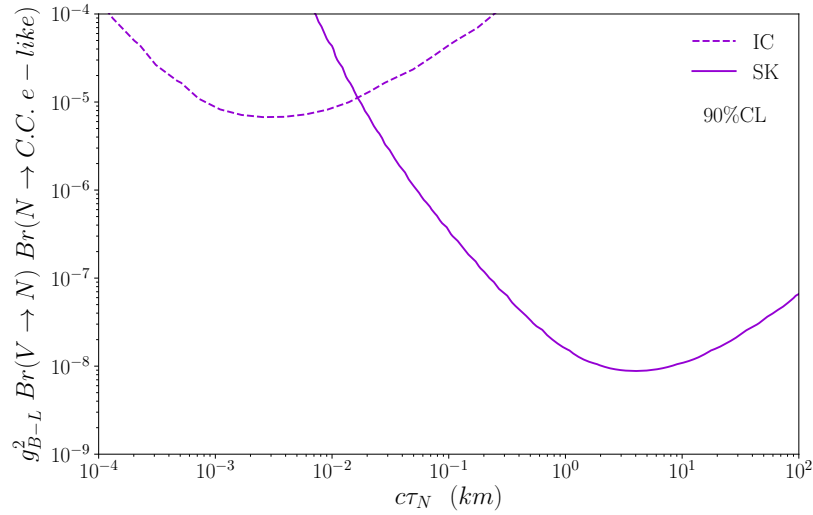


Figure 17. Limits on dark photons decays from IceCube (thick lines) and Super-Kamiokande (thin lines) on the $g_{B-L}^2 \text{Br}(V \rightarrow N) \text{Br}(N \rightarrow CCe - like)$ vs $c\tau_N$ for $m_V = 0.8$ GeV and $m_N = 0.35$ GeV, where $\text{Br}(V \rightarrow N) = 2\text{Br}(V \rightarrow NN)$.

On the other hand, in the $B - L$ model it is possible to uncorrelate production and decay since the LLP is not the dark photon, but the HNL to which it decays. Therefore, production rates in this model are controlled by the $B - L$ gauge coupling, while the decay depends on the mixing with the light neutrinos. The limits for the $B - L$ model on the plane $g_{B-L}^2 \text{Br}(V \rightarrow N) \text{Br}(N \rightarrow CCe - like)$ vs. $c\tau_N$ are shown in figure 17, for $m_V = 0.8$ GeV and $m_N = 0.35$ GeV. As in the case of the dark photon, the limits for SK are three orders of magnitude better since the production in this case is controlled by

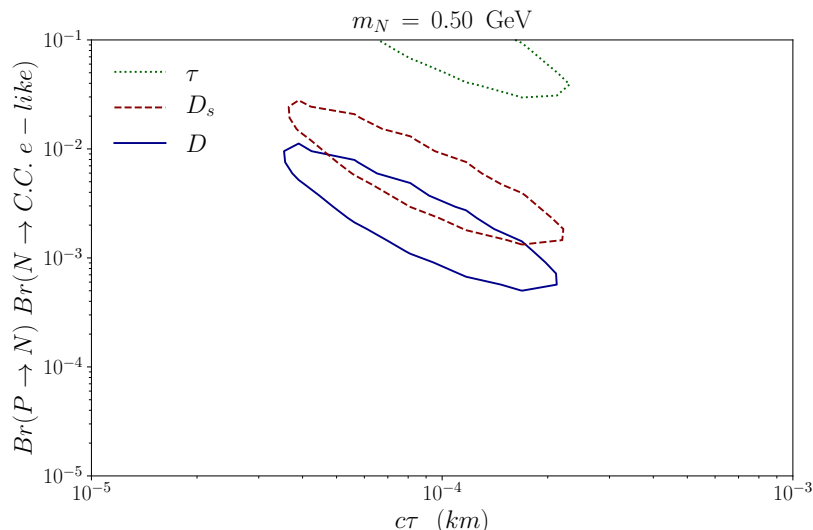


Figure 18. 1σ best-fit regions for HNL decays in IceCube assuming uncorrelated $\text{Br}(P \rightarrow N)$ vs $c\tau$ for $m_N = 0.5 \text{ GeV}$ and considering separately the production from three parent particles $P = D, D_s, \tau$.

p bremsstrahlung. While IC limits fall short of competitiveness, assuming large enough branching ratios $\text{Br}(V \rightarrow N)\text{Br}(\text{CCe} - \text{like})$ our results show that SK would be able to reach values of $g_{B-L}^2 \leq 10^{-7}$, which are comparable to BaBar limits in this region [62, 70, 71]. At the same time, the values of $c\tau_N \sim \mathcal{O}(1) \text{ km}$ needed for the HNL are allowed by present constraints, see figure 6.

5.2 IceCube preferred regions

An interesting observation is that, in the case of IC, the minimization of the χ^2 on these two-dimensional planes gives a minimum of the χ^2 that lies at a non-zero value of the LLP production rate. This is the result from a small excess in the data at around 30 TeV, as can be seen by naked eye in figure 13. Although the source of the excess is unclear, and it might be related to an underestimation of the μ background or an unidentified source of systematic uncertainties, in the remainder of this section we explore possible explanations in the LLP models discussed in section 4.

Figure 18 shows the 1σ regions for 2 d.o.f., i.e. the region $\Delta\chi^2 = \chi^2 - \chi_{\min}^2 \leq 2.3$, where $\chi_{\min}^2$ is the global minimum, for a HNL produced from D, D_s and τ decays. On the other hand, $\chi^2(\text{no LLP}) - \chi_{\min}^2 \simeq 7$. Although our results are shown for $m_N = 0.5 \text{ GeV}$, the regions do not change significantly for other values of m_N between 0.5 and 1.5 GeV. For the dark photon scenario, figure 19 shows the allowed 1σ regions (2 d.o.f.) for three values of the dark photon mass. Unfortunately the minimal models discussed in sections 4.1 and 4.2 cannot live on those regions, according to figures 6 (for the HNL) and figure 9 (for the dark photon).

Regarding the $B - L$ model with HNL, as in the previous cases we get a best fit to IC data away from zero and, in particular, for $\text{Br}(V \rightarrow N)g_{B-L}^2 \sim 10^{-6} - 10^{-5}$ and $c\tau_N \sim 10^{-4} \text{ km}$. Unfortunately, the required value of g_{B-L} is excluded by collider con-

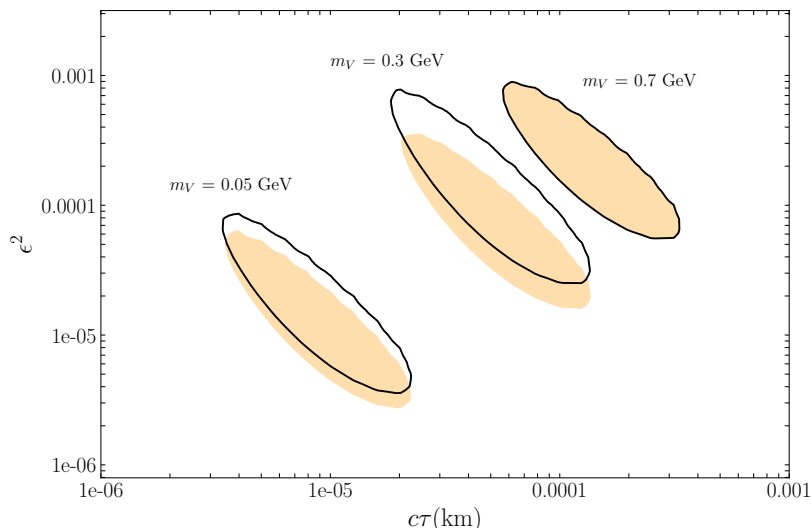


Figure 19. 1σ best fit regions on dark photons decays in IceCube assuming uncorrelated ϵ^2 vs $c\tau$. The black lines correspond to considering only the signal from π_0 decay for $m_V = 0.05$ GeV, only η decay for $m_V = 0.3$ GeV and only bremsstrahlung for the heavier mass.

straints [62, 70, 71] in this dark photon mass range and, at the same time, such as small $c\tau_N$ for a HNL at $m_N = 0.35$ GeV is at least one order of magnitude too small, assuming that the mixing is below the present constraints from direct searches of HNL [13].

Before finalizing this section it should be stressed, however, that our IC analysis has been performed using just the energy information available on the data release of ref. [39], and that in the case of decay of an LLP the zenith angle distribution of the signal events could differ significantly from the distributions observed for the data (see figure 9 of [39]). More stringent constraints might be attainable if information on the zenith angle distribution of the events were included in the analysis, which is not possible with the publicly available information.

6 Conclusions

Very weakly interacting particles beyond the Standard Model (SM) might be light enough to be produced in accelerators, but easy to miss in standard BSM search analyses if they are very long-lived. Many running accelerator experiments, both fixed-target and colliders, are actively exploring new strategies to improve the sensitivity to such exotic signals, and several future experiments are being proposed to target specifically such BSM scenarios.

In this paper we have presented the first detailed analysis of the sensitivity of large neutrino detectors, such as Super-Kamiokande and IceCube, to putative long-lived particles (LLP) that might be produced in atmospheric showers from the decay of standard model mesons and/or bremsstrahlung. We have presented the methodology to evaluate LLP fluxes from these processes, and the procedure to quantify the expected number of signal events. We have considered in particular two minimal and theoretically well-motivated scenarios where the LLP are either heavy neutral leptons produced primarily

in meson decays, or dark photons produced from bremsstrahlung or/and radiative π_0 and η decays. We have evaluated bounds using Super-Kamiokande and IceCube data from refs. [39, 44], as a function of the branching ratios (or production rates) of $\text{SM} \rightarrow \text{LLPs}$, and the LLP lifetime. The best sensitivity is obtained for lifetimes of $\mathcal{O}(10^{-2} \times \frac{m_{\text{LLP}}}{1\text{GeV}})$ km in Icecube and $\mathcal{O}(10 \times \frac{m_{\text{LLP}}}{1\text{GeV}})$ km for Super-Kamiokande. However, in the minimal models considered in this work production and decay are usually strongly correlated since both are controlled by the same coupling. In this case, we find that atmospheric searches do not lead to competitive bounds. On the other hand, our bounds might be complementary to other searches in more complex models where production and decay are uncorrelated. As an example for such scenario we considered an extended model with a $B-L$ gauge boson that couples to the HNL. In this case, the production is controlled by the gauge coupling, g_{B-L} , while the LLP is the HNL whose lifetime is controlled by its mixing with the light neutrinos. We have also presented the sensitivity to this scenario.

Interestingly, in the case of Icecube the addition of an LLP to the SM provides a better fit to the observed data than the prediction obtained in the SM (with a $\Delta\chi^2 \sim 7$), due to a small excess observed at around 30 TeV. This is obtained, for example, for a HNL produced in D meson decays with a $\text{Br}(D \rightarrow X) \sim 2 \cdot 10^{-3}$ and a lifetime $c\tau \sim 10^{-4}$. However, we find that the values of the parameters needed in order to fit the excess lie outside of the allowed regions of parameter space for the three models considered.

Finally, we stress again that no systematic uncertainties have been included in our calculations. A more detailed study, including a realistic implementation of systematic and reconstruction errors by the experimental collaborations, would be mandatory to validate our results.

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A Computation of effective decay areas

Assuming the fiducial volume of the detector is a cylinder of radius R and height H , it is a simple geometrical exercise to obtain $\Delta\ell_{\text{det}}$, and the effective area:

$$A_{\text{decay}}^{\text{eff}}(E_A, \cos\theta) = |\cos\theta|A_1(E_A, \cos\theta) + |\sin\theta|A_2(E_A, \cos\theta), \quad (\text{A.1})$$

where A_1 correspond to the flux entering from the top:

$$A_1(E_A, \cos \theta) = \int_0^R r dr \int_0^{2\pi} d\phi \left(1 - e^{-\frac{\Delta \ell_{\text{det}}^{(1)}(\theta, r, \phi)}{c\tau_{\text{lab}}(E_A)}} \right), \quad (\text{A.2})$$

with

$$\Delta \ell_{\text{det}}^{(1)}(\theta, r, \phi) \equiv \text{Min} \left[\frac{H}{|\cos \theta|}, \frac{R\sqrt{1 - r^2/R^2 \sin^2 \phi} + r \cos \phi}{|\sin \theta|} \right], \quad (\text{A.3})$$

and A_2 is that entering laterally:

$$A_2(E_A, \cos \theta) = R \int_0^H dx \int_{-\pi/2}^{\pi/2} d\phi \left(1 - e^{-\frac{\Delta \ell_{\text{det}}^{(2)}(\theta, x, \phi)}{c\tau_{\text{lab}}(E_A)}} \right), \quad (\text{A.4})$$

with

$$\Delta \ell_{\text{det}}^{(2)}(\theta, x, \phi) \equiv \text{Min} \left[\frac{H - x}{|\cos \theta|}, \frac{2R \cos \phi}{|\sin \theta|} \right]. \quad (\text{A.5})$$

B Computation of parent particle fluxes in the atmosphere

In all our calculations, the SM parent particle fluxes in the atmosphere have been computed using the MCEq software [24, 25], with the SYBILL-2.3 hadronic interaction model [26], the Hillas-Gaisser cosmic-ray model [27] and the NRLMSISE-00 atmospheric model [28]. Obtaining the parent particle fluxes is however not straightforward, because only the flux for protons and unstable particles that are relatively long-lived compared to their interaction length can be directly extracted using MCEq. A possible strategy to circumvent this is to manually switch off the decay of the parent particle and extract the flux, which then needs to be corrected to account for the decay.²

To simplify notation, we denote the differential flux of the parent particle P by ϕ_P :

$$\phi_P \equiv \frac{d\Phi_P}{dE}. \quad (\text{B.1})$$

Let us assume that $\tilde{\phi}_P$ is the differential flux of parent particle P *assuming* that P is stable.³ Assuming that P is directly produced in nucleon interactions, this flux satisfies a cascade equation of the form [23]

$$\frac{d\tilde{\phi}_P}{dX} = -\frac{\tilde{\phi}_P}{\lambda_P} + Z_{NP} \frac{\phi_N}{\lambda_N} + Z_{PP} \frac{\tilde{\phi}_P}{\lambda_P}, \quad (\text{B.2})$$

where ϕ_N is the flux of nucleons, and the spectrum-weighted momenta are defined as

$$Z_{kh} = \int_E^\infty dE_P \frac{\phi_k(E_P)}{\phi_k(E)} \frac{\lambda_k(E)}{\lambda_k(E_P)} \frac{dn(k\mathcal{N} \rightarrow hY; E_P, E)}{dE}. \quad (\text{B.3})$$

²We thank A. Fedynitch for providing guidance on this point.

³In MCEq, this can be easily extracted by manually switching off the decay of the particle P (using the advanced feature settings).

Here, λ_k is the particle interaction length of hadron k and $\frac{dn(k\mathcal{N}\rightarrow hY)}{dE}$ is the number of hadrons h produced with energy between E and $E+dE$, in the scattering of hadron k with energy E_P on a nucleus $\mathcal{N}$.

The second and third terms in eq. (B.2) are the production and regeneration terms, respectively, while the first term is the absorption term. Had the particle decayed, the correct cascade equation should have also included a decay term:

$$\frac{d\phi_P}{dX} = -\frac{\phi_P}{\rho d_P} - \frac{\phi_P}{\lambda_P} + Z_{NP}\frac{\phi_N}{\lambda_N} + Z_{PP}\frac{\phi_P}{\lambda_P}, \quad (\text{B.4})$$

where d_P is the decay length of the parent particle in the laboratory frame and ρ is the column density at the point with slant depth X . It is easy to see that from $\tilde{\phi}_P$ we can get the solution to eq. (B.4) as

$$\phi_P = \int_0^X dX' e^{-(X-X')\left(\frac{1}{\rho d_P} + \frac{1}{\lambda_P}\right)} \left(\frac{d\tilde{\phi}_P(X')}{dX} + \frac{\tilde{\phi}_P(X')}{\Lambda_P} \right), \quad (\text{B.5})$$

with

$$\Lambda_P \equiv \frac{\lambda_P}{1 - Z_{PP}}. \quad (\text{B.6})$$

To the best of our knowledge, although MCEq allows to obtain the values of λ_P it does not allow to extract the values of Z_{PP} directly. Therefore, we assume $Z_{\tau\tau} = 0$ while for the rest of the mesons considered we assume Feynmann scaling (that is, $Z_{MM} = 0.3$) since, according to figure 4 in ref. [23], this is a reasonable approximation for charged pions and kaons. In any case, we do not expect our results to change significantly if these assumptions are modified.

With this procedure we get a reasonable agreement with the output of MCEq available at large energies, where decay can be ignored. We do a small rescaling of our approximate result by matching the two curves in the overlapping region.

Finally, in the case of D and D_s mesons both the interaction and regeneration terms in eq. (B.4) can be neglected. In this case, the meson flux may be computed directly from the flux of protons in the atmosphere, following ref. [23] (eqs. (25–27) in that reference):

$$\phi_M^{\text{low}} \simeq Z_{NM} \frac{\rho d_M}{\lambda_N} \phi_N(X), \quad (\text{B.7})$$

where $\phi_N(X)$ and the yields, Z_{NM} , are extracted from MCEq. We have also checked that this method gives a good agreement with our fluxes, in the case of D and D_s mesons.

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New constraints on heavy neutral leptons from Super-Kamiokande data

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Abstract Heavy neutral leptons are predicted in many extensions of the Standard Model with massive neutrinos. If kinematically accessible, they can be copiously produced from kaon and pion decays in atmospheric showers, and subsequently decay inside large neutrino detectors. We perform a search for these long-lived particles using Super-Kamiokande multi-GeV neutrino data and derive stringent limits on the mixing with electron, muon and tau neutrinos as a function of the long-lived particle mass. We also present the limits on the branching ratio versus lifetime plane, which are helpful in determining the constraints in non-minimal models where the heavy neutral leptons have new interactions with the Standard Model.

1 Introduction

There are compelling reasons to believe that neutrino masses are the first manifestation of a new physics (NP) scale, which can be identified with the mass of the heavy mediator(s) that generates neutrino masses. Under this assumption, neutrino masses and mixings provide information on a combination of the mediator mass and couplings. Although an upper limit of the new physics scale can be derived by requiring the new couplings to be perturbative, generically, no lower bound results from this constraint. As it is well known, a very high NP scale leads to a hierarchy problem [1,2], and to vacuum instability issues [3,4], both of which can be avoided if the NP scale is not much higher than the electro-weak scale. The possibility that the neutrino mass mediators are light enough to be produced in accelerators and in atmospheric showers is therefore worth exploring.

The most popular extension of the Standard Model (SM) realizing neutrino masses is the Type-I seesaw model [5–8], with at least two heavy Majorana singlets N_j :

$$\mathcal{L}_N \supset \sum_j i \bar{N}_j \not{\partial} N_j - \left(Y_{\alpha j} \bar{L}_\alpha \tilde{\Phi} N_j + \frac{m_{N_j}}{2} \bar{N}_j N_j^c \right) + \text{h.c.}, \quad (1)$$

where $\tilde{\Phi} \equiv i\sigma_2 \Phi^*$ is the complex conjugate of the Higgs field Φ , L_α is the SM lepton doublet with flavor α , Y is a generic Yukawa matrix, and m_{N_j} is the Majorana mass of the singlet N_j . After spontaneous electro-weak symmetry breaking, the heavy Majorana states mix with the standard neutrinos resulting in a spectrum of three light states with masses $m_\nu \propto m_N^{-1} (Yv)^2$, and two or more heavy neutral leptons (HNL) with masses $\propto m_N$. In this model, all massive neutrino states are admixtures of the standard neutrinos and the singlet states, as dictated by the mixing matrix $U_{\alpha j}$ (which diagonalizes the mass Lagrangian of the whole system). The phenomenology of the HNL depends crucially on their mass and their mixing to the charged leptons. In fact, it is through this mixing that the heavy singlets can be produced either through charged-current (CC) or neutral-current (NC) processes and also how they decay back to SM particles. For simplicity, in this work we will adopt the simplified notation $|U_\alpha| \equiv U_{\alpha j}$, assuming that only one of these states is kinematically accessible for our purposes.

HNLs have been extensively searched for in laboratory experiments, using mainly two types of signatures: either displaced vertices from the decay of the HNL, or through precise determination of the decay product kinematics in meson decays (see, e.g., Refs. [9–12] for reviews of available constraints in the MeV–GeV mass range, or Refs. [13–23] for future prospects to improve over current bounds). In Ref. [24] we studied in detail the search for long-lived particles produced in the atmosphere, which would then decay inside large-volume neutrino detectors. We derived bounds on HNL for masses above the kaon mass, where laboratory limits are weaker. We found, however, that the limits from atmospheric searches are not competitive with laboratory constraints in this mass range.

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In this letter we focus on the lighter mass region instead, where the HNL can be produced in kaon and pion decays for which the atmospheric flux is significantly larger. The most stringent bounds for HNL below the kaon mass come from peak searches [25, 26] and displaced HNL decay vertex searches [27–29]. However, for masses below the kaon and pion mass the HNL becomes very long-lived: although the value of its lifetime in the rest frame (τ) depends strongly on its mass and on its mixing with the light states, in the minimal model described above it ranges between $c\tau \sim (10^{-4} - 50) \times |U_\alpha|^{-2}$ (km), for m_N between 40 MeV and 400 MeV. This makes atmospheric neutrino detectors a well-suited setup to search for their decay products.

In this work, we use the framework developed in Ref. [24] to extract limits from the multi-GeV muon and electron neutrino data samples observed at Super-Kamiokande (SK). The search of atmospheric HNL in a similar mass range has been considered before for SK [30, 31] and IceCube [32]. Our analysis significantly improves the methodology of these earlier studies. In addition, extensions of the minimal model of Eq. (1) that decouple production and decay have been considered in recent works, particularly in relation with the LSND/MiniBoone anomaly [33–38]. In Refs. [39–42] an extension that includes a dipole interaction of the heavy singlets provides a new radiative decay channel, which dominates the HNL decay and significantly reduces its lifetime. Therefore, we will present our limits not only in the context of the minimal HNL model of Eq. (1) (that is, on the plane $|U_\alpha|$ versus m_N), but also on the plane $\text{Br}(K/\pi \rightarrow N)$ versus $c\tau$, which is phenomenologically motivated. As we will see, this can be useful in order to constrain non-minimal scenarios with uncorrelated production and decay, such as for example the dipole extension of Ref. [42].

2 HNL production mechanisms

The leading mechanism for HNL production is through the decays of mesons produced in the atmosphere. The computation is explained in detail in Ref. [24], and here we summarize it for convenience. Defining ℓ as the distance traveled between the production point of the HNL to its point of entry in the detector, the HNL production profile (in a differential of distance $d\ell$) can be obtained as:

$$\frac{d\Pi_N^{\text{decay}}}{dE d\cos\theta d\ell} = \sum_P \int_{E_P^{\min}}^{E_P^{\max}} \frac{dE_P}{\gamma_P \beta_P c \tau_P} \frac{d\Phi_P}{dE_P d\cos\theta} \frac{dn}{dE}, \quad (2)$$

where dn/dE stands for the differential distribution of the HNL energies while Φ_P is the parent meson flux and, in this work, $P = K^\pm, \pi^\pm$. We have used the Matrix Cascade Equation (MCEq) Monte Carlo software [43, 44] to compute the fluxes for the parent mesons in the atmosphere,

with the¹ SYBILL-2.3 hadronic interaction model [45], the Hillas-Gaisser cosmic-ray model [46] and the NRLMSISE-00 atmospheric model [47]. In Eq. (2), β_P and γ_P are the parent boost factors while τ_P is its lifetime in the rest frame. For the dominant two-body decay $K^\pm, \pi^\pm \rightarrow N l_\alpha^\pm$ (denoted generally as $P \rightarrow NY$), the differential distribution dn/dE reads

$$\begin{aligned} \frac{dn(E_P, E)}{dE} &= \frac{1}{\Gamma_P} \frac{d\Gamma(P \rightarrow NY)}{dE} \\ &= \frac{\text{Br}(P \rightarrow NY)}{\Gamma(P \rightarrow NY)} \frac{d\Gamma(P \rightarrow NY)}{dE}, \end{aligned} \quad (3)$$

where Br stands for branching ratio, and

$$\frac{1}{\Gamma(P \rightarrow NY)} \frac{d\Gamma(P \rightarrow NY)}{dE} = \frac{1}{p_P \sqrt{\lambda(1, y_N^2, y_Y^2)}}, \quad (4)$$

with $y_i \equiv \frac{m_i}{m_P}$ and

$$\lambda(a, b, c) \equiv a^2 + b^2 + c^2 - 2ab - 2ac - 2bc. \quad (5)$$

Finally, the kinematical limits for $E_P \gg m_P$ are:

$$E_P^{\text{max(min)}} \equiv \frac{2E}{(1 + y_A^2 - y_Y^2) \mp \sqrt{\lambda(1, y_A^2, y_Y^2)}}. \quad (6)$$

Figure 1 shows a representative example of the HNL production from kaon and pion decays in the atmosphere, and compares it to the result obtained for heavier parent mesons and τ leptons in our previous work [24]. As can be seen from this figure, the production profile grows by several orders of magnitude when the mass of the HNL allows for it to be produced from the decays of lighter resonances, due to their more abundant fluxes in the atmosphere.

A second contribution to the flux comes from the HNL production in NC scattering of standard atmospheric neutrinos as they pass through the Earth's matter (for instance in $\nu \mathcal{N} \rightarrow NX$, where $\mathcal{N}$ stands for a nucleon and X is a hadronic shower). This contribution can be estimated as follows. Assuming a flux of standard neutrinos ϕ_ν , the HNL produced at a distance r from SK is:

$$\frac{d\Pi_N^{\text{int}}}{d\cos\theta dE dr} = \int_{E_\nu^{\min}}^{E_\nu^{\max}} dE_\nu \frac{d\phi_\nu(r)}{d\cos\theta dE_\nu} \frac{d\sigma_{\nu \rightarrow N}(E_\nu)}{dE} N_A \rho_\oplus, \quad (7)$$

where $\sigma_{\nu \rightarrow N}$ is the HNL production cross section (which can be estimated as the standard ν NC scattering cross section multiplied by the corresponding mixing $|U_\alpha|^2$), while N_A is the Avogadro number and $\rho_\oplus$ is the Earth density at distance

¹ While significant variations are expected for the prompt neutrino flux if the cosmic ray or hadronic interaction models are changed (see e.g. Ref. [43]), the conventional neutrino contribution is understood at the $\mathcal{O}(10 - 20\%)$ level. We expect similar variations in our results, if these assumptions are modified.

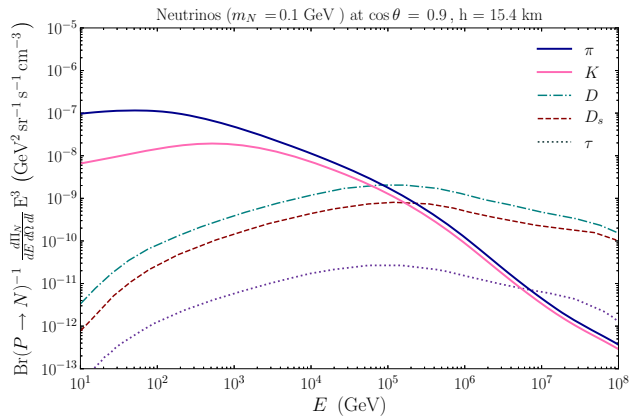


Fig. 1 HNL production rate (for $m_N = 0.1$ GeV) per energy, solid angle and path length, at a height of 15.4 km and $\cos \theta = 0.9$, for different parent particles P . In this work, we consider only π and K decays. For comparison we also show the corresponding result for two-body decays of heavier mesons (D , D_s) and leptonic (three-body) τ decays from Ref. [24]. For all lines shown here, effects due to the mass of the charged lepton produced together with the HNL have been neglected for simplicity

r . This assumes that both standard neutrinos and the resulting HNL are highly boosted and therefore the zenith angle of the HNL is roughly the same as that of the incoming neutrino. The upper and lower limits in the integral correspond to the kinematically allowed ranges, where a SM neutrino scattering on a nucleon at rest can produce a HNL with energy E .

3 HNL decays in Super-Kamiokande

The flux that arrives to the detector for HNL produced in meson decays, Φ_N^{decay} , is obtained integrating over all LLPs produced at different distances ℓ , weighted by their corresponding survival probabilities, as

$$\frac{d\Phi_N^{\text{decay}}}{d\cos\theta dE} = \int_0^{\ell_{\text{max}}} d\ell \frac{d\Pi_N^{\text{decay}}}{d\cos\theta dE d\ell} e^{-\ell/L_{\text{decay}}}, \quad (8)$$

where L_{decay} is the decay length of the HNL in the laboratory frame. The maximum distance $\ell_{\text{max}} \equiv \ell(h_{\text{max}}, \theta)$ is a simple function of the maximum height of the atmosphere where cosmic showers are produced, $h_{\text{max}} \simeq 80$ km, and the zenith angle. Analogously, the flux entering the detector for HNL produced in the interaction of SM neutrinos in the Earth, Φ_N^{int} is:

$$\frac{d\Phi_N^{\text{int}}}{d\cos\theta dE} = \int_0^{R_{\text{max}}(\theta)} dr \frac{d\Pi_N}{d\cos\theta dE dr} e^{-r/L_{\text{decay}}}, \quad (9)$$

where $R_{\text{max}}(\theta) = 2R_{\oplus} \cos \theta$ is the maximum distance traveled through the Earth for trajectories with a zenith angle θ , $R_{\oplus}$ being the Earth's radius.

The total number of HNL decays inside the detector, within a given time window ΔT , energies in the interval $[E, E + dE]$ and zenith $[\cos \theta, \cos \theta + d\cos \theta]$, can be computed as

$$\frac{dN}{dEd\cos\theta} = \Delta T A_{\text{decay}}^{\text{eff}}(E, \cos\theta) \frac{d\Phi_N}{dEd\cos\theta}, \quad (10)$$

where $\Phi_N \equiv \Phi_N^{\text{int}} + \Phi_N^{\text{dec}}$, and A^{eff} is an effective area which accounts for the probability that a decay takes place inside the detector. This area can be estimated integrating the surface of the detector normal to the flux direction, weighted by the decay probability of the N inside the detector:

$$A_{\text{decay}}^{\text{eff}}(E, \cos\theta) = \int dS_{\perp} \left[1 - e^{-\frac{\Delta\ell_{\text{det}}(\cos\theta)}{L_{\text{decay}}(E)}} \right]. \quad (11)$$

Here $\Delta\ell_{\text{det}}$ is the length of the segment of the HNL trajectory that cuts into the detector (for explicit expressions we refer the reader to Ref. [24]). A cylindrical geometry for SK with height of 40 m and radius of 20 m is assumed.

The two contributions to the total number of events (coming from meson decays and from SM neutrino interactions in the Earth) have a very different angular dependence: while the flux from decays is expected to be larger from above, those from interactions come obviously from below. However, we have checked that the final contribution to the number of events coming from HNL produced in meson decays is several orders of magnitude larger than the one obtained from interactions of SM neutrinos on the Earth. The ratio between the two contributions decreases for larger $c\tau$ but, within the range $c\tau \in [1, 10^4]$ km, it lies within the range $[10^6, 10^2]$. We can therefore safely neglect the contribution from neutrino interactions in the Earth and, in the rest of this work, we will only consider HNL production from meson decays.

In order to derive our limits, we use the data samples as well as the expected neutrino background prediction from Fig. 5 of Ref. [48]. We use both the μ - and e -like fully-contained multi-GeV events, adding the single- and multi-ring samples together (labeled as “multi-GeV” and “multi-ring” in Fig. 5 in Ref. [48], respectively). In the case of e -like events, we add both the ν_e and $\bar{\nu}_e$ samples as well. In Ref. [48] data are binned in zenith angle, while information on the energy of the events is not publicly available. Therefore, in this work the data is binned only in $\cos \theta$. As for the neutrino energies considered, we integrate over all energies between 1 GeV and 90 GeV, as this is the range corresponding to the fully-contained sample (see Fig. 6 in Ref. [48]). We think this is conservative, as SK may be sensitive to events outside this range.

The number of events observed in a given sample will also depend on detector efficiencies and reconstruction effects, which should be included in the form of migration matrices giving the relationship between true and reconstructed vari-

ables. Such information is however not publicly available for SK. Thus, here we make the simplifying assumption that the efficiencies are independent on the neutrino energy and, in particular, we take a flat detection efficiency $\epsilon^\alpha = 0.75$ both for μ and e -like, in line with the values quoted in Ref. [48] for the multi-GeV ν_e event sample. While a priori some loss of efficiency could be expected at high energies (mainly due to a reduction in the containment of the events), we think that the impact on our results would be small since we expect our sensitivities to come mostly from the events at low energies. This is because the heavy neutrino flux produced in the atmosphere will follow a very steep power law, peaked at low energies as shown in Fig. 1, where we expect the SK efficiencies to be best. We also assume that the angular reconstruction is much better than the width of the bins in zenith angle, so migration between different bins in $\cos\theta$ can be neglected. If this assumption were to be relaxed, our sensitivities would probably be worsened. This is so because in the case of very long-lived particles (as is the case for HNL with masses below 500 MeV, and weakly coupled to the SM) their very long lifetimes lead to a higher sensitivity for specific angular bins, where the distance traveled by the HNL is comparable to (or smaller than) its lifetime in the laboratory frame. While here we have used only publicly available information, a detailed evaluation by the experimental collaboration is needed to validate our results. Finally, we should also point out that we expect an increase in sensitivity if the analysis were to be carried out using both energy and angular information. This is beyond the scope of this work.

The number of events in the i -th bin in $\cos\theta$ in a given sample can therefore be computed as:

$$N_i^{\text{SK}, \alpha\text{-like}} = \text{Br}(N \rightarrow \alpha\text{-like}) \times \int_{1 \text{ GeV}}^{90 \text{ GeV}} dE \epsilon^\alpha(E) \int_{\cos\theta_i^{\min}}^{\cos\theta_i^{\max}} d\cos\theta \frac{dN}{dE d\cos\theta}, \quad (12)$$

where $\cos\theta_i^{\min}$ and $\cos\theta_i^{\max}$ are the lower and upper limits of the bin. Here, $\text{Br}(\alpha\text{-like})$ stands for the total branching ratio for all decay channels including muons, electrons or photons in the final state, depending on the sample (α) considered. For example, in the case of μ -like we consider only those decay channels including one or more muons in the final state. In the case of e -like events we require that no muons are present, but we also include decay channels with photons as these are easily mis-identified with electrons at SK (such as $N \rightarrow \nu\pi^0$, since the π^0 decays promptly to two photons).

4 Results

In this section we derive limits on HNL production by performing a χ^2 fit to the SK data. A Poissonian χ^2 function

has been used:

$$\chi^2 = 2 \sum_{i,\alpha} N_i^\alpha + B_i^\alpha - n_i^\alpha + n_i^\alpha \log \left(\frac{n_i^\alpha}{N_i^\alpha + B_i^\alpha} \right), \quad (13)$$

where the sum runs over the angular bins, and $\alpha = \{e\text{-like}, \mu\text{-like}\}$. Here, n_i^α stands for the data observed in each bin while N_i^α is the predicted number of signal events and B_i^α is the background prediction, which includes the predicted number of atmospheric neutrino events in the SM. In our χ^2 calculations, we consider separately the e - and μ -like samples and we add their two contributions to the total χ^2 . Therefore, our limits will be derived using 20 degrees of freedom, corresponding to the total number of bins in $\cos\theta$. In our calculations we find, however, that the sensitivity is largely dominated by the e -like contribution since the size of the branching ratio $\text{Br}(e\text{-like})$ is much larger than $\text{Br}(\mu\text{-like})$ in the mass range considered.

First we show in Fig. 2 the results on the plane $\text{Br}(K/\pi \rightarrow N) \times \text{Br}(N \rightarrow \text{visible})$, where $\text{Br}(N \rightarrow \text{visible})$ accounts for the probability that the HNL decays visibly in the detector. Our results are presented as a function of $c\tau$, assuming no correlation between the production and decay mechanisms for the HNL. It is interesting to note that SK can outperform the powerful displaced-decay search limits from beam dumps such as those in Refs. [27–29], for models where the decay does not involve two charged tracks, since all laboratory searches request this condition to reduce background contamination. In the extended model of Ref. [42] this is precisely the case, since the HNL decay is dominated by the radiative $N \rightarrow \nu\gamma$ decay via the dipole interaction, therefore the stringent limits from PS191 [27,28] and the recent T2K limits [29] do not apply. The shaded purple region in Fig. 2 shows the range where the MiniBooNE anomaly could be explained, for a HNL with $m_N = 260$ MeV, extracted from Ref. [42]. In this case, the HNL would only be produced in kaon decays and, unfortunately, our limits from $K^\pm$ decay (pink lines) fall short to probe this region. However, SK would be sensitive to non-minimal models with HNL produced in $\pi^\pm$ decays with a similar value of the production BR and lifetime and, obviously, larger neutrino experiments such as DeepCore or Hyper-Kamiokande could significantly improve over these constraints.

In addition, the shaded light blue region in Fig. 2 shows the expectation for the minimal model outlined in Eq. (1), for $m_N = 250$ MeV and mixing matrix elements in the range $|U_{e/\mu}|^2 \in [10^{-8}, 10^{-10}]$, $|U_\tau|^2 \in [10^{-10}, 10^{-4}]$ (in agreement with current constraints from Ref. [9]). As readily seen from this figure, relevant constraints are expected in this case. In view of this result, next we derive constraints on the minimal scenario, assuming only one non-

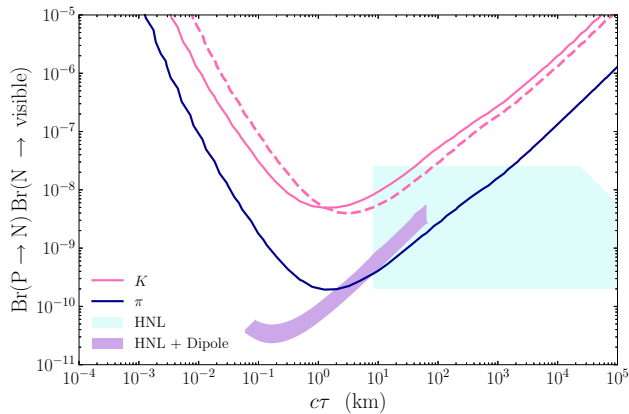


Fig. 2 Limits at 90% confidence level (CL), on the plane $\text{Br}(K/\pi \rightarrow N) \times \text{Br}(N \rightarrow \text{visible})$ versus $c\tau$ for $m_N = 0.1$ GeV (solid) and $m_N = 0.25$ GeV (dashed), for HNL produced in π (dark blue) or K (pink) decays. Our limits have been obtained assuming uncorrelated production and decay mechanisms for the HNL. The shaded light blue region corresponds to the (correlated) values obtained for the minimal model, varying the mixings within their presently allowed constraints as explained in the text. The purple area corresponds to the 1σ allowed region where a HNL with a dipole interaction could explain the MiniBooNE anomaly, according to Ref. [42], for $m_N = 260$ MeV

vanishing U_α at a time. Our results are shown in Fig. 3 for $|U_e|^2$ (left panel) and for $|U_\mu|^2$ (right panel), at 90% confidence level (CL). In the case of $|U_e|^2$ the contribution from $\pi^\pm$ decays is clearly dominant for $m_N < 140$ MeV, as can be seen from the peak in sensitivity at around 0.1 GeV. We find that the limits derived from our simplified analysis is already able to set tight constraints on the mixing of HNL with the electron and muon neutrino sectors, between 10^{-6} and 10^{-7} for m_N in the range between 150 MeV and 450 MeV. Our limits are also compared with those obtained from displaced decay searches at PS191 [27,28] and at the T2K near detector [29], as well as from peak

searches in E949 [25], PIENU [49], and NA62 [26]. We also show the resulting bound derived in Ref. [11] from the measurement of the kaon decays into electrons or muons [50]. In the case of $|U_e|^2$, the limits obtained from SK are comparable or even better than analogous limits from peak searches, while they are not competitive with those from displaced decay searches. In the case of $|U_\mu|^2$, peak searches in E949 also yield better constraints than our limits from SK data.

Finally, while the HNL cannot be produced via $|U_\tau|$ in K or π decays (because it is not possible to produce the HNL together with a τ lepton in this case), competitive limits can still be derived on U_τ if we allow for non-vanishing $|U_\alpha|$, $\alpha = e, \mu$, even if these are well below present bounds from laboratory experiments. The reason is that, in this case, the HNL could be produced via the mixing U_e or U_μ , and a large U_τ can induce a significant decrease in its lifetime of the HNL while allowing for a significant branching ratio into e -like or μ -like events through NC-mediated decays. Therefore, in Fig. 4 we show the sensitivity to $|U_\tau|$, as a function of m_N , for fixed $U_e = 10^{-8}$ or $|U_\mu| = 10^{-8}$ (which are both below the best present upper bounds). Our limits obtained in this way are already much better than existing direct constraints on $|U_\tau|$ from CHARM [51], which however have been obtained assuming vanishing values for $|U_e|$ and $|U_\mu|$. A similar exercise can be done for PS191, which would probably lead to better limits on $|U_\tau|$ than the SK results, for the same assumed values of $|U_e|$ and $|U_\mu|$.

5 Conclusions

In this letter, we have used the latest publicly available SK data to derive strong constraints on HNL production from

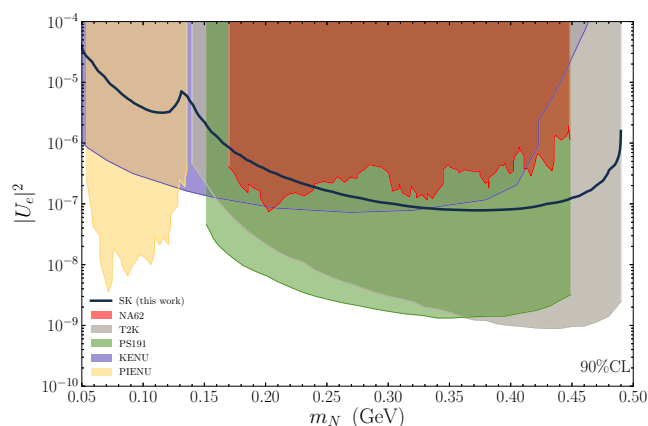
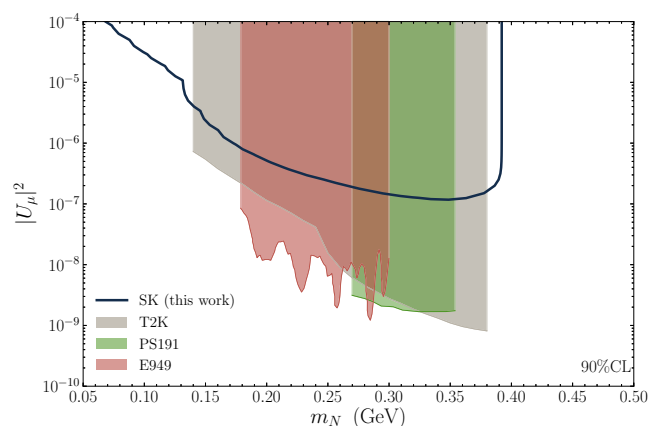


Fig. 3 SK constraints on the minimal HNL scenario at 90% CL, projected onto the plane $|U_e|^2$ vs m_N (for $U_{\mu,\tau} = 0$). Our results (solid black lines) are compared to corresponding limits obtained for the T2K



near detector [29], NA62 [26], E949 [25], PS191 [27,28], and PIENU [49]. The line labeled as KENU was derived in Ref. [11] using precision measurements of leptonic decay channels of the kaon [50]

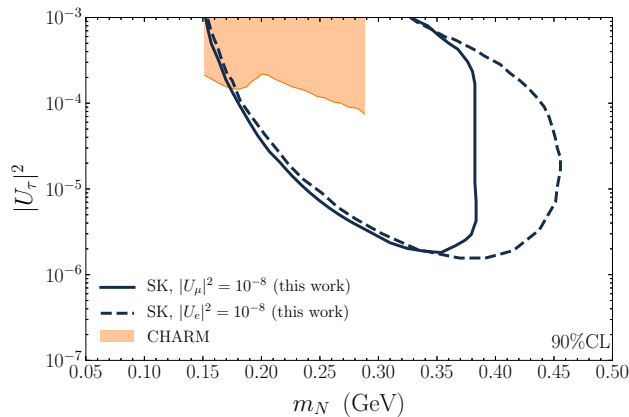


Fig. 4 SK constraints on the minimal HNL scenario at 90% CL, projected onto the plane $|U_\tau|^2$ vs m_N , for $|U_e|^2 = 10^{-8}$ and $|U_\mu|^2 = 10^{-8}$. Our results (black lines) are compared to the limits from CHARM [51] (which however have been obtained under the assumption $|U_e|^2 = |U_\mu|^2 = 0$)

kaon and pion decays in the atmosphere. Using a χ^2 analysis, and binning our events only in $\cos\theta$, we find that SK data is able to provide strong constraints on the minimal HNL scenario for masses between 150 MeV and 400 MeV. It is therefore expected that a more detailed analysis performed by the collaboration may be able to significantly improve over our results. We have also shown our limits in the Br vs $c\tau$ plane, which is applicable to a wider range of NP models where the HNL interact with the SM not only via mixing but also through other interactions (such as, for instance, a dipole moment). Finally, we have used our results to show how, in the case of non-vanishing U_e or U_μ well below current constraints from laboratory searches, SK data could be used to set strong bounds on U_τ , well below the direct limits presently available from CHARM data. We expect a similar improvement of PS191 bounds under similar assumptions, though.

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Searching for light long-lived neutralinos at Super-Kamiokande

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Light neutralinos could be copiously produced from the decays of mesons generated in cosmic-ray air showers. These neutralinos can be long-lived particles in the context of R -parity violating (RPV) supersymmetric models, implying that they could be capable of reaching the surface of the earth and decay within the instrumental volume of large neutrino detectors. In this article, we use atmospheric neutrino data from the Super-Kamiokande experiment to derive novel constraints for the RPV couplings involved in the production of long-lived light neutralinos from the decays of charged D -mesons and kaons. Our results highlight the potential of neutrino detectors to search for long-lived particles, by demonstrating that it is possible to explore regions of parameter space that are not yet constrained by any fixed-target nor collider experiments.

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I. INTRODUCTION

The discovery of the Higgs boson at the Large Hadron Collider (LHC) in 2012 [1,2] provides not only conclusive evidence of the Standard Model (SM), but also consolidates the hierarchy problem as one of the main theoretical puzzles in modern physics [3]. In this context, supersymmetry (SUSY) remains as one of the most compelling possibilities to address this problem [4,5]. At the same time, supersymmetry also provides a rich and complex phenomenology which has lead to an intensive search program at collider experiments [6].

Conventional SUSY theories assume a discrete symmetry called R -parity, which avoids conflict with experimental data on the nonobservation of baryon and lepton number violating processes, such as proton decay [7] and neutrinoless double beta decay [8]. Within the context of R -parity conserving SUSY theories, the lightest neutralino as

the lightest supersymmetric particle (LSP) provides a natural candidate for fermionic dark matter because of its stability and lack of electromagnetic interactions (see [9,10] for seminal articles and [11] for a review). Yet, it is possible to assume R -parity violation (RPV) [12] while respecting the bounds on the proton lifetime, as long as baryon numbers or lepton numbers are preserved. In such a case, the neutralino LSP is no longer stable and can decay into SM particles. The smallness of the R -parity violating couplings can make the decay macroscopic, making the neutralino, $\tilde{\chi}_1^0$, a long-lived particle (LLP) [13]. This implies that neutralinos can leave a variety of exotic signatures at colliders such as a displaced vertex (reviews on possible long-lived particle signals can be found in [14,15]). On the other hand, unlike the strong interacting sparticles whose masses have a lower bound at around 1 TeV [16–24], the neutralino mass is less constrained, and can in principle be massless [25]. For a fraction of the R -parity violating model parameter space, the ATLAS and CMS experimental collaborations at the LHC have searched for a long-lived $\tilde{\chi}_1^0$ with a mass $\mathcal{O}(100)$ GeV with a displaced vertex signature [26–28], with null results. The most up to date constraint comes from ATLAS with leptonic displaced decays, excluding long-lived neutralinos between 50 GeV–500 GeV [27].

Phenomenological prospects demonstrate that in extended regions in the R -parity violating couplings—

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leading to decays with different leptonic or hadronic final state particles—and lighter masses below 100 GeV, a long-lived $\tilde{\chi}_1^0$ has the potential to be discovered at dedicated LLP experimental facilities that could operate at the LHC, such as FASER, MATHUSLA, CODEX-b and AL3X [29–34], future colliders at the intensity frontier [34], or beam-dump experiments such as SHiP [32]. In particular, the absence of experimental constraints in the range from a few MeV to $\mathcal{O}(1)$ GeV makes it a good candidate to be studied in scenarios where they can be produced from the mesons that are abundantly created in both colliders and cosmic-ray air showers.

So far, prospects for light, long-lived neutralinos have been performed from the decays of D and B mesons in references [30,32,35–37], from Z boson decays in [29,31,34,38], and from the decay of τ leptons in Belle-II [33]. Most of these searches are based on experiments that are either in construction like FASER [39], or even in earlier stages as they are subject to funding/approval (e.g., SHiP [40], MATHUSLA [41], and others [42,43]). In contrast to this situation, large Cherenkov based neutrino detectors such as IceCube [44] or Super-Kamiokande (SK) [45] are already built and have been taking data for years, which can be used to search for LLPs, as the decay of these particles would generate a signal that is indistinguishable from the Cherenkov radiation measured in association with the neutrino interactions in the medium.

Searches for long-lived particles at large neutrino detectors were considered in the literature in references [46–51]. Of these studies, Ref. [49] used public data from IceCube and Super-Kamiokande to search for long-lived particles produced in atmospheric showers, using detailed numerical simulations. The results found demonstrated that the atmospheric neutrino data from the Super-Kamiokande experiment can be used to place stringent constraints to models predicting LLPs (see also Ref. [50]). The main reason for this is because the atmospheric neutrino data focused on the low energy regime, where the flux of the atmospheric particles peaks. In this work we build upon the general results of [49], and apply the same strategy to search for neutralinos produced from meson decays, including an appropriate treatment of the uncertainties originating from the hadronic interaction models used to simulate the production of mesons in the shower.

We consider the possibility of searching for neutralinos produced in the decays of D -mesons and kaons that are generated in the sky, when highly energetic cosmic rays collide with the upper layers of the atmosphere. The former process allows us to compare our results with the previously mentioned studies in the literature, particularly references [30,32], where the sensitivity reach at collider experiments was estimated for selected benchmarks. On the other hand, neutralino production from kaon decay was considered long before in the context of spontaneously broken supersymmetry [52]. In this work, we explore a

novel RPV channel through kaon production that is particularly well suited for our setup because the production of kaons in the atmosphere greatly exceeds the one from D -mesons [53].

The rest of the paper is organized as follows. In Sec. II we summarize the relevant phenomenological aspects of the RPV theory, and describe the neutralino production from the decay of pseudoscalar mesons generated in cosmic-ray air showers. In Sec. III we describe the expected signal from the visible decay of the neutralino in the Super-Kamiokande detector. The main results obtained for the two benchmark scenarios considered in this work are presented in section Sec. IV, along with the current best constraints from previous studies. Finally, we draw our main conclusions in section Sec. V.

II. LONG-LIVED LIGHT NEUTRALINOS IN COSMIC-RAY AIR SHOWERS

A. Neutralinos in RPV

The simplest realistic realization of a supersymmetric model is the so-called minimal supersymmetric standard model [54]. As mentioned above, in this model a Z_2 symmetry called R -parity is imposed to avoid proton decay. However, it is still possible to break R -parity and have a stable proton by imposing a different discrete symmetry like the baryon triality B_3 symmetry [55]. Without imposing R -parity, the most general Lagrangian that respects gauge and space-time symmetries contains the term [56]

$$\mathcal{L} \supset \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k^c, \quad (1)$$

where the hatted symbols denote gauge multiplets of superfields that include leptons ℓ_i and sleptons $\tilde{\ell}_i$ in $\hat{L}_i$, left-handed quarks q_{jL} , and squarks $\tilde{q}_{jL}$ in $\hat{Q}_j$, and right-handed antiquarks $\tilde{d}_{kR}$ and antisquarks $\tilde{d}_{kR}^*$ in $\hat{D}_k^c$. The family indices i, j and k run from one to three, leading to 27 independent parameters. This semileptonic operator violates lepton number by one unit, and can generate contributions to the mass and magnetic moment of neutrinos [57], neutrinoless double beta decay [58,59], and in the parameter space of our interest, neutralino production from meson decays [60–62]. Comprehensive reviews of RPV and its phenomenological implications can be found in [12,63–65].

In this work we focus on the trilinear couplings λ'_{121} and λ'_{112} , which allow us to produce neutralinos from the decays of D -mesons and kaons, respectively. These decays are always accompanied with an electron, as detailed in our two benchmarks in Table II. The introduction of RPV parameters also allows neutralino two-body decays into a meson and a lepton. In our first benchmark, we study neutralinos with masses in the range from the mass of the kaon to the mass of the D -meson, while in the second benchmark case we restrict the neutralino mass to the interval between the mass of the pion and the mass of the

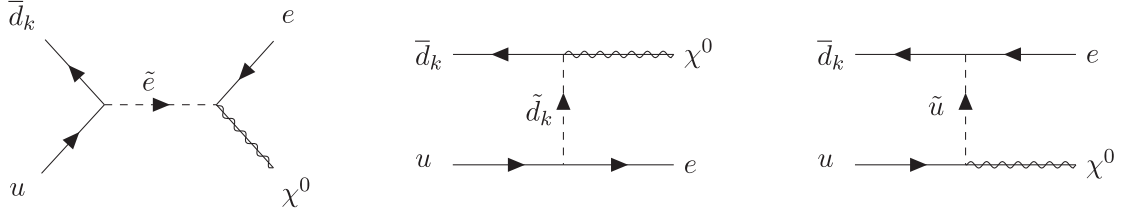


FIG. 1. Neutralino production diagrams. We display the specific case where production occurs via the decay of a charged meson M_{kj}^+ composed of quarks d_k^C and u_j . For the cases of our interest (see Table II), we set the flavor indices of the initial states in the diagram such that M_{kj}^+ corresponds to D^+ ($k = 1, j = 2$) for benchmark B1 and K^+ ($k = 2, j = 1$) for benchmark B2.

kaon. Therefore, for the second benchmark, we include the trilinear λ'_{111} , as it is the only coupling associated with pions, and is therefore necessary to make the neutralino decay when its mass is lower than the kaon mass. The relevant Feynman diagrams involved in the production of neutralinos from these mesons are shown in Fig. 1.

To calculate the decay rate of the mesons to neutralinos, we closely follow reference [32] and consider the effective interactions involving the relevant mesons, leptons, and neutralinos. Assuming that the sfermion masses involved are degenerate and large enough so that they can be integrated out, the relevant decay rates are

$$\Gamma(D^\pm \rightarrow \tilde{\chi}_1^0 + e^\pm) = \frac{\mathcal{K}^{1/2}(m_D^2, m_{\tilde{\chi}_1^0}^2, m_e^2)}{64\pi m_D^3 (m_c + m_d)^2} |G_{121}^S|^2 f_D^2 \times m_D^4 (m_D^2 - m_{\tilde{\chi}_1^0}^2 - m_e^2), \quad (2)$$

and

$$\Gamma(K^\pm \rightarrow \tilde{\chi}_1^0 + e^\pm) = \frac{\mathcal{K}^{1/2}(m_K^2, m_{\tilde{\chi}_1^0}^2, m_e^2)}{64\pi m_K^3 (m_u + m_s)^2} |G_{112}^S|^2 f_K^2 \times m_K^4 (m_K^2 - m_{\tilde{\chi}_1^0}^2 - m_e^2), \quad (3)$$

where $\mathcal{K}(x, y, z) \equiv x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ is the Källén function [66], and f_D, f_K are the D -meson and kaon decay constants, respectively. The effective couplings G_{ijk}^S are given by

$$G_{ijk}^S \equiv \frac{3g_2}{2\sqrt{2}} \frac{\lambda'_{ijk}}{m_f^2} \tan \theta_W, \quad (4)$$

where m_f represents the common value that we assume for the masses of sleptons and squarks, which are the sfermions involved in neutralino production (details can be found in reference [32]). In Eq. (4), $\theta_W \equiv \tan^{-1} g_1/g_2$ is the weak mixing angle [21], and g_1, g_2 are the gauge coupling constants associated to $U(1)_Y$ and $SU(2)_L$, respectively. The symbols m_K and m_D denote the masses of the $K^\pm$ and $D^\pm$ mesons, and m_c, m_d, m_u , and m_s are the masses of the charm, down, up, and strange quarks, respectively. In our

calculations, we use the values $f_K \simeq 156$ MeV and $f_D \simeq 213$ MeV [21].

B. Neutralinos from D-meson and kaon decays in atmospheric showers

Cosmic rays hitting our atmosphere provide us with a beam of protons (and other species) that is constantly switched on. A single cosmic ray can produce an extensive cascade of charged particles and radiation called an air shower [67–69]. Mesons in the shower, including D -mesons and kaons, decay to charged leptons and neutrinos, among other particles. The flux of leptons can be measured both at the surface of the earth and in underground experiments, while their careful reconstruction allows to estimate the expected mesonic contributions to the spectrum [53,70–77]. Similarly as with the case of neutrinos, we simulate the production of light neutralinos in the shower by solving the cascade equation involving only source terms from meson decays [74]

$$\frac{d\Phi_{\tilde{\chi}_1^0}}{dE_{\tilde{\chi}_1^0} d\Omega dX} = \sum_M \int dE_M \frac{1}{\rho \lambda_M} \frac{d\Phi_M}{dE_M d\Omega} \frac{dn}{dE_{\tilde{\chi}_1^0}}, \quad (5)$$

where the sum runs over all possible mesons that can decay to neutralinos when a given trilinear coupling λ'_{ijk} is switched on. In Eq. (5), ρ is the density of the atmosphere at a column depth X , and $\lambda_M \equiv \gamma_M \beta_M c \tau_M$ is the decay length of the meson, which includes the boost factor $\gamma_M \beta_M$ and its proper lifetime τ_M . The differential production rate of mesons in the shower per unit of solid angle is given by $\frac{d\Phi_M}{dE_M d\Omega}$. The number of neutralinos with energies between $E_{\tilde{\chi}_1^0}$ and $E_{\tilde{\chi}_1^0} + dE_{\tilde{\chi}_1^0}$ produced in the decay of the meson M is given by $\frac{dn}{dE_{\tilde{\chi}_1^0}}$. For two-body decays, this last quantity is given by

$$\frac{dn}{dE_{\tilde{\chi}_1^0}} = \frac{\text{Br}(M \rightarrow \tilde{\chi}_1^0 + e)}{p_M \sqrt{\mathcal{K}(1, \frac{m_e^2}{m_M^2}, \frac{m_{\tilde{\chi}_1^0}^2}{m_M^2})}}, \quad (6)$$

where $\text{Br}(M \rightarrow \tilde{\chi}_1^0 + e)$ is the branching fraction of meson decays to neutralinos and p_M is the meson momentum.

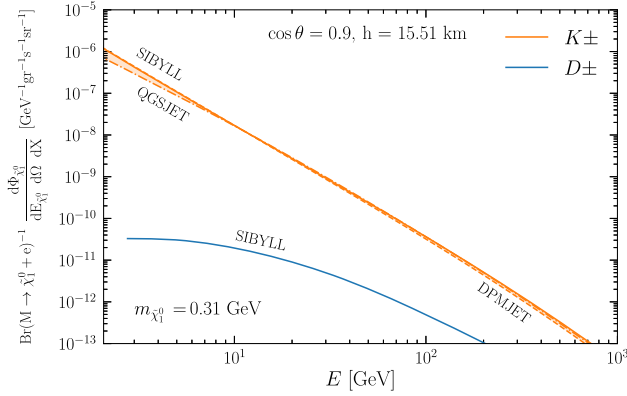


FIG. 2. *Neutralino Production Rate.* Energy spectrum of the atmospheric production rate of light neutralinos from charged D -mesons (blue curve), and kaons (orange band). The spread of the orange band reflects the difference in the production originated from the election of different hadronic event generators; SIBYLL (solid line), QGSJET (dot-dashed line), and DPMJET (dashed line). We choose a representative mass for the neutralino with a value of 0.31 GeV. The angular direction is fixed at $\cos \theta = 0.9$, while the height was chosen around 15 km, where a maximum production rate is expected.

In Eq. (5), both the density of the atmosphere and the differential production rate of mesons are extracted using the Matrix Cascade Equation (MCEq) software package [75,77]. Here, we choose the NRLMSISE-00 atmospheric model [78], while for the hadronic interaction model we consider the different event generators which are updated with LHC data (see for instance, references [79,80]). In particular, we focus on the SYBILL-2.3 [53], QGSJET-II-04 [81], EPOS-LHC [82], and DPMJET-III [83] models. These models might yield non-negligible differences in the production rate of mesons and can be a relevant source of uncertainty in our calculations.¹

As an illustration of our numerical simulations, we show in Fig. 2 the production rate for a 0.31 GeV neutralino, at a height h of 15.51 km and with a vertical direction of 25.84° degrees from the zenith, produced in the decay of both D -mesons and kaons. The spread of the orange band in the plot covers the most pessimistic and optimist cases given the uncertainties associated with the election of an event generator. Note that there is an important difference in the amount of neutralinos produced from the decay of D -mesons and kaons, as the latter are expected to be orders of magnitude more abundant than in the atmosphere [53]. Finally, it is important to remark that we were not able to assess the uncertainty that pertains to the simulations from D -mesons, as currently the only updated hadronic

interaction model for charmed hadrons is SYBILL-2.3 [53,85–87].

The uncertainties are estimated following reference [88] (see also reference [89] for a similar discussion). For a given meson, we calculate the total expected flux $\frac{d\phi}{dE}$ at the surface of the earth with different hadronic models. In order to do so, we let the mesons propagate through the atmosphere without decaying, which can be accomplished by switching off their decay in MCEq. The impact of the variation in the meson production rate for different models is then quantified with respect to a benchmark model, which we take here to be the SYBILL-2.3 model, by defining the ratio $\Delta_j(M)$ between the meson flux predicted by the event generators that are being compared

$$\Delta_j(M) = \frac{\int_{E_{\min}}^{\Lambda} dE \frac{d\phi_{\text{BM}}}{dE}}{\int_{E_{\min}}^{\Lambda} dE \frac{d\phi_j}{dE}}, \quad (7)$$

where M is the meson of interest, j is the index used to denote the model that is being compared with the benchmark model BM, $E_{\min} = 1.6$ GeV is the minimum energy available in MCEq, and $\Lambda = 10^3$ GeV is the upper energy cutoff that we use in order to obtain the neutralino production rate from a given parent meson. Table I shows the Δ coefficients obtained for the production of kaons with different event generators. The uncertainty as quantified by Eq. (7) reaches a maximum of around 66% for QGSJET. We calculate the neutralino production rate using different hadronic interaction models, and assess their impact on the region of parameter space that can be probed. We stress that dedicated efforts are needed in order to reduce the uncertainties involved in meson production in the forward direction, a problem that is also crucial for precise modeling of neutrino fluxes at the LHC [90–92].

III. NEUTRALINO SIGNALS AT SUPER-KAMIOKANDE

After light neutralinos are produced from the decay of mesons in the atmospheric shower, they decay to SM particles as they propagate. We assume that all the particles in the decay chain are highly boosted, and hence approximately collinear in their trajectories. The degree of attenuation of the flux that arrives at the detector depends

TABLE I. *Comparison of the meson flux using different hadronic interaction models.* Numerical values of Δ , the relative integrated meson flux with respect to SYBILL, for the different hadronic interaction models considered in this work.

Model	$\Delta_j(K^\pm)$
DPMJET	1.163
EPOS-LHC	1.116
QGSJET	1.660

¹Another relevant input in this calculation is the cosmic-ray model chosen for the primary spectra. We checked that the uncertainties from this election are sub-leading in the energy regime considered in this work. We use the Hillas-Gaisser cosmic-ray model H3a [84] in our calculations.

on the relation between the lifetime of the neutralino and the distance it travels before reaching the detector. In particular, it is expected that the detector signal from up-going events (i.e., the neutralinos that arrive from below the detector), will be suppressed in comparison with down-going events (i.e., the neutralinos that reach the detector from above). Assuming that the production rate of neutralinos is symmetric with respect to the azimuthal component, the expected differential flux is obtained by integrating over the column depth X , according to

$$\frac{d\Phi_{\tilde{\chi}_1^0}}{dE_{\tilde{\chi}_1^0} d\cos\theta} = 2\pi \int dX \frac{d\Phi_{\tilde{\chi}_1^0}}{dE_{\tilde{\chi}_1^0} d\cos\theta dX} e^{-\ell/\lambda_{\tilde{\chi}_1^0}}, \quad (8)$$

where ℓ corresponds to the traveled distance of the neutralino, measured from the production point, which can be obtained from the height h and the zenith angle θ using the geometrical relation

$$h = \sqrt{R_\oplus^2 + 2\ell R_\oplus \cos\theta + \ell^2} - R_\oplus, \quad (9)$$

with $R_\oplus$ the earth's radius. In Eq. (8), the lifetime of the neutralino enters via the exponential decay factor that depends on the decay length $\lambda_{\tilde{\chi}_1^0} = \gamma_{\tilde{\chi}_1^0} \beta_{\tilde{\chi}_1^0} c \tau_{\tilde{\chi}_1^0}$.

The neutralinos that reach the Super-Kamiokande detector can decay to SM particles inside its instrumental volume, which we model as a cylinder of 20 meters in radius and 40 meters in height [45]. Among the possible decay products of the neutralinos it is possible to find charged leptons and mesons, as well as neutral pseudoscalar and vector mesons. Different decay products correspond to different kinds of signals in the detector, which at Super-Kamiokande can be classified in two ways: showering (or e -like), and nonshowering (or μ -like events). The former originates from electromagnetic and hadronic showers in the detector, while the latter is primarily associated to muons, which leave a distinctive Cherenkov ring with crisp edges. In this work, we focus on the showering signals originated from the decay of neutralinos described in Table II. In particular, for the case of neutral mesons in the final states, we assume that they decay to a showering signal inside the volume of the detector. The parameter choice of the benchmark B1 also allows K_L^0 as a final state for neutralino decays, but we do not consider this as part of our signal since this particle will typically decay outside the detector due to its large decay length.

For our benchmark scenarios, the important decay width formulas of neutralinos to pseudoscalar and vector mesons accompanied with a lepton are [32]

$$\Gamma(\tilde{\chi}_1^0 \rightarrow M_{jk} + \ell_i) = \frac{\mathcal{K}^{1/2}(m_{\tilde{\chi}_1^0}^2, m_{M_{jk}}^2, m_{\ell_i}^2)}{128\pi m_{\tilde{\chi}_1^0}^3 (m_{q_j} + m_{q_k})^2} |G_{ijk}^S|^2 f_{M_{jk}}^2 \times m_{M_{ik}}^4 (m_{M_{ab}}^2 - m_{\tilde{\chi}_1^0}^2 - m_{\ell_i}^2), \quad (10)$$

$$\Gamma(\tilde{\chi}_1^0 \rightarrow M_{jk}^* + \ell_i) = \frac{\mathcal{K}^{1/2}(m_{\tilde{\chi}_1^0}^2, m_{M_{jk}^*}^2, m_{\ell_i}^2)}{2\pi m_{\tilde{\chi}_1^0}^3} |G_{ijk}^T|^2 \times (f_{M_{jk}^*}^V)^2 [2(m_{\tilde{\chi}_1^0}^2 - m_{\ell_i}^2)^2 - m_{M_{jk}^*}^2 (m_{M_{jk}^*}^2 + m_{\tilde{\chi}_1^0}^2 + m_{\ell_i}^2)], \quad (11)$$

where

$$G_{ijk}^T \equiv \frac{g_2}{4\sqrt{2}} \frac{\lambda'_{ijk}}{m_f^2} \tan\theta_W. \quad (12)$$

In the equation above, M_{ij} (M_{ij}^*) represents one of the final-state pseudoscalar (vector) mesons listed in the neutralino decays of Table II. The indices i, j designate the family of the meson's valence quarks. For the final state pseudoscalar mesons, we use the decay constants $f_\pi \simeq 130$ MeV [21], $f_{\pi^0} = f_\pi/\sqrt{2}$, and $f_{K^0} = f_K/\sqrt{2}$ [93], where f_K was defined after Eq. (4). For vector mesons, on the other hand, $f_{M^*}^V$ represents the vector meson decay constant, and for K^{*0}, K^{*+} we use the approximate value $f_{K^*}^T \simeq 230$ MeV [32,93]. As mentioned before, although for the production of neutralinos from D -mesons and kaons we only need the trilinear couplings λ'_{121} and λ'_{112} , the decay of neutralinos produced from kaons into visible showers in SK proceeds via the coupling λ'_{111} . The two benchmark scenarios

TABLE II. Parameter choices that define our benchmark scenarios B1 and B2. The CP conjugate processes of all the decays shown are also allowed and therefore contribute to the neutralino decay width. The decay modes displayed in this table are compatible with a shower (or e -like) signal in Super-Kamiokande. See text for details.

	RPV coupling	Production	Decay mode
B1	$\lambda'_{121}, \lambda'_{112}$	$D^\pm \xrightarrow{\lambda'_{121}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K_S^0 + \nu_e$
			$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{121}} K^{*0} + \nu_e$
			$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K^{(*)+} + e^-$
			$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{112}} K_S^0 + \nu_e$
B2	$\lambda'_{112}, \lambda'_{111}$	$K^\pm \xrightarrow{\lambda'_{112}} e^\pm + \tilde{\chi}_1^0$	$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^+ + e^-$
			$\tilde{\chi}_1^0 \xrightarrow{\lambda'_{111}} \pi^0 + \nu_e$

considered in this article can be seen in Table II. For benchmark B1 we consider neutralino masses in the range $m_{K^+} + m_e \leq m_{\tilde{\chi}_1^0} \leq m_{D^+} - m_e$, while for B2 the neutralino mass lies in the range $m_{\pi^0} + m_e \leq m_{\tilde{\chi}_1^0} \leq m_{K^+} - m_e$. The election of B1 allows for a direct comparison with references [30,32], where the sensitivity reach for future experiments aiming to explore the lifetime frontier was estimated. The election of B2 allows us to study a novel production channel where neutralinos can be abundantly produced in atmospheric showers. Furthermore, it can be seen from Eqs. (4) and (12) that it is possible to combine the dependence of the decay widths on the RPV parameters and the sfermion mass in the ratio λ'_{ijk}/m_f^2 , which we set as a free parameter.

To obtain the expected number of events with an energy in the range $E_{\tilde{\chi}_1^0}$ and $E_{\tilde{\chi}_1^0} + dE_{\tilde{\chi}_1^0}$, and with trajectories within $\cos \theta$ and $\cos \theta + d\cos \theta$, we include the effective surface S_{eff} for a decay to take place inside the SK detector, such that for a given time window ΔT the event rate is

$$\frac{dN}{dE_{\tilde{\chi}_1^0} d\cos \theta} = \Delta T \times \frac{d\Phi_{\tilde{\chi}_1^0}}{dE_{\tilde{\chi}_1^0} d\cos \theta} \times S_{\text{eff}}, \quad (13)$$

where the effective surface is obtained by integrating the surface of the detector that is perpendicular to the incoming direction of the neutralino flux, weighted by the probability that the particle decays inside the detector

$$S_{\text{eff}} = \int dS_{\perp} (1 - e^{-\Delta \ell_{\text{det}}/\lambda_{\tilde{\chi}_1^0}}). \quad (14)$$

Here, $\Delta \ell_{\text{det}}$ is the segment of the particle trajectory that traverses the detector, for which explicit analytical expressions can be found in the Appendix of Ref. [49]. The computation of the effective surface for decay is a purely geometrical problem. Figure 3 shows the effective surface as a function of the decay length of the neutralino, with trajectories fixed by different values of the cosine of the zenith angle given by 0.9, 0.5, and 0.1. The event distribution in Eq. (13) can be integrated in energies and trajectories within a given bin determined by the resolution of the experiment. As mentioned before, we use atmospheric neutrino data reported by the Super-Kamiokande experiment in reference [94]. This data contains the angular distribution of events involving electron and muon neutrinos, from different energy regimes; the Sub-GeV and Multi-GeV sample of events with energies below and above 1330 MeV, respectively. Taking into account the trilinear couplings considered in this article, as well as the minimum energy available in MCEq, we chose the showering (or e -like) event sample in the Multi-GeV energy window. The total events reported in reference [94] correspond to the SK-I up to the SK-IV data taking periods with a total run of 5,326 days.

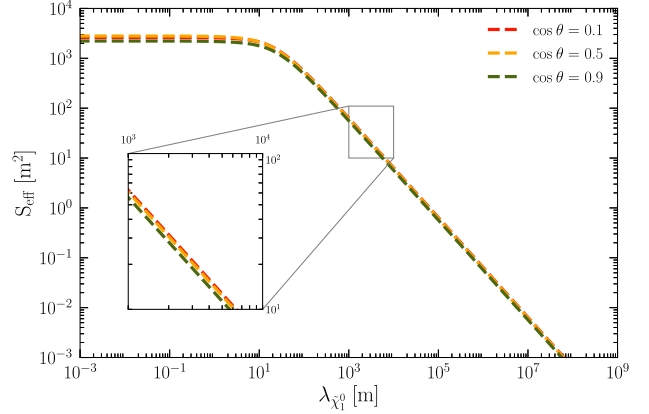


FIG. 3. *Effective Surface for decay.* The expected area of decay in Super-Kamiokande is shown as a function of the neutralino decay length for different incoming directions fixed by values of the cosine of the zenith angle at 0.1 (red line), 0.5 (yellow line), and 0.9 (green line). A similar result holds for negative values of $\cos \theta$.

The number of events contained in the Multi-GeV range for the i th bin in the cosine of the zenith angle is computed with the formula

$$N_{c_i} = \text{Br}(\tilde{\chi}_1^0 \rightarrow e\text{-like}) \int_{\cos \theta_i - 0.1}^{\cos \theta_i + 0.1} d\cos \theta \times \int_{E_{\text{min}}}^{E_{\text{max}}} dE_{\tilde{\chi}_1^0} \epsilon \frac{dN}{dE_{\tilde{\chi}_1^0} d\cos \theta}, \quad (15)$$

where ϵ is the detection efficiency, which for the Multi-GeV sample is flat in energy with a value of 0.75, while E_{min} and E_{max} are equal to 1.5 GeV and 90.5 GeV, respectively [94]. The background for our search corresponds to the expected e -like events from electron neutrinos at SK as reported in reference [94]. As an illustration, in Fig. 4 we show an example of the expected event signal generated from neutralino decays, when the couplings λ'_{112} and λ'_{111} are both fixed to a value of 0.002, and the neutralino mass is 0.31 GeV.

IV. RESULTS AND DISCUSSION

The signal computed in Eq. (15) depends on the neutralino mass and RPV couplings through its lifetime, branching fraction of production from mesons, and the fraction of neutralinos that decay to an e -like signal in the detector. We adopt the distribution for Poisson events, which implies a chi-squared statistic of the form

$$\chi^2 = 2 \sum_{c_i} \left[N_{c_i} (m_{\tilde{\chi}_1^0}, \lambda') + B_{c_i} - D_{c_i} + D_{c_i} \log \left(\frac{D_{c_i}}{N_{c_i} (m_{\tilde{\chi}_1^0}, \lambda') + B_{c_i}} \right) \right], \quad (16)$$

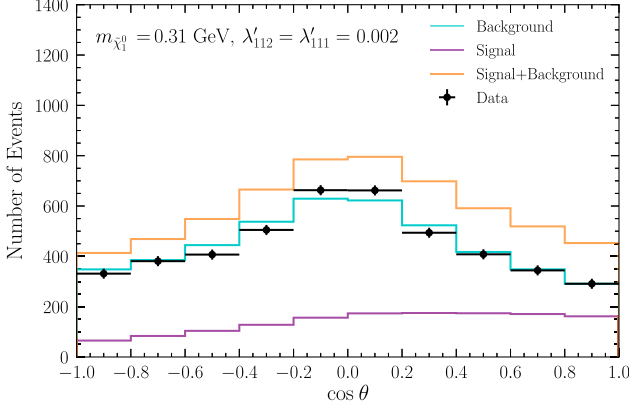


FIG. 4. *Expected events at SK.* Distribution of atmospheric events in Super-Kamiokande for incoming directions fixed by the cosine of the zenith angle. Data points are shown in black, while the background is depicted by the cyan line. The signal produced from the decay of neutralinos generated from kaons is shown in purple. The signal plus background contribution is shown in orange. The mass value and RPV couplings are fixed as indicated in the figure.

where the sum runs over the ten bins in cosine of the zenith angle, and B_{c_i} , D_{c_i} , are the background and data in the i th bin, respectively. We apply this statistical test to the SK data and derive constraints at 90% confidence level. The results for our two benchmark scenarios are displayed in Figs. 5 and 6.

There are numerous phenomenological constraints on the lepton number violating operator $\lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k^c$ (see [93] and references therein). In particular, as pointed out in [95], RPV parameters can be subject to strong constraints due to their contribution to meson oscillation observables. This is the strongest bound that applies to the benchmark B1 specified in Table II, where the tree-level contributions to kaon oscillations induced by that choice of parameters imply the sneutrino mass dependent bound [96]

$$|\lambda'_{112}\lambda'_{121}| \leq 2.2 \times 10^{-8} \left(\frac{m_{\tilde{\nu}_e}}{1 \text{ TeV}} \right)^2, \quad (17)$$

where $m_{\tilde{\nu}_e}$ is the sneutrino mass. Our results for B1 are shown Fig. 5, where we evaluated the limits for a sneutrino and squark mass of 1 TeV. As it can be seen on the left panel, the limits imposed by kaon oscillations for degenerated sneutrino and squark masses exclude all the parameter space within the projected sensitivity reach at SHiP [32], our limits from SK, and the limit from the Drell-Yan (DY) monolepton process $pp \rightarrow \ell \bar{\nu}$ at the LHC (labeled as ‘Colliders’). The latter is a single coupling bound for the parameter λ'_{112} given by [97]

$$\lambda'_{112} \leq 0.16 \frac{m_{\tilde{s}_R}}{1 \text{ TeV}} + 0.030, \quad (18)$$

where $m_{\tilde{s}_R}$ is the strange squark mass, which is also set to 1 TeV to compare with our results. However, since in

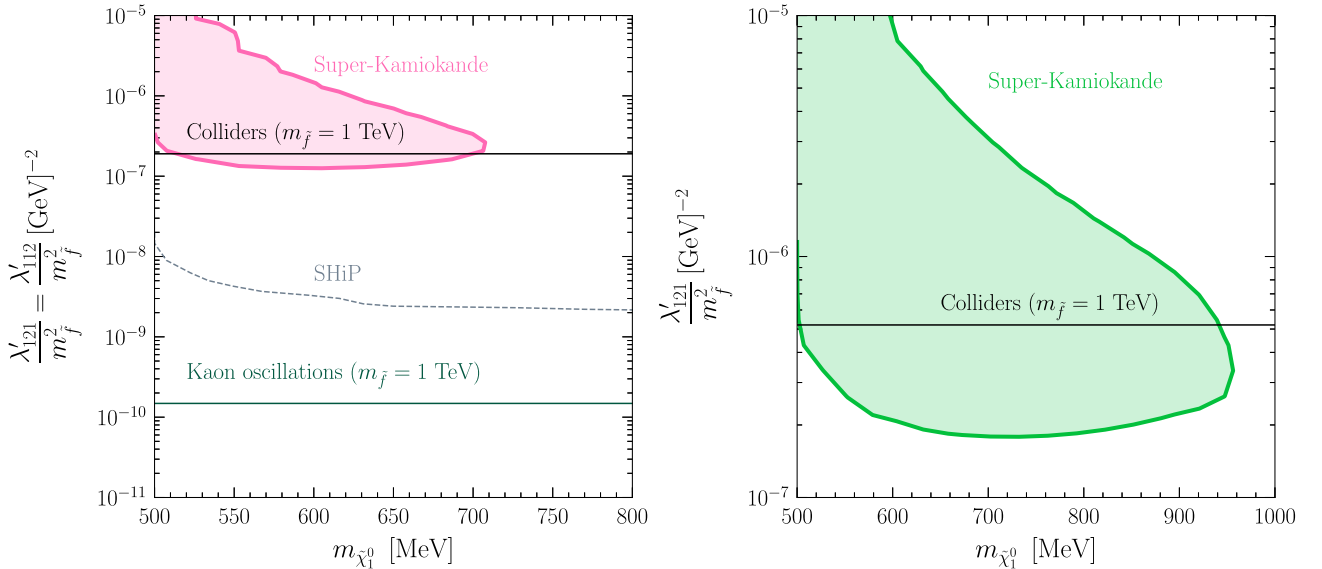


FIG. 5. *Limits for benchmark B1.* Limits derived at 90% confidence level in the parameter space defined by the neutralino mass $m_{\tilde{\chi}_1^0}$ and the trilinear RPV parameters λ'_{121} , λ'_{112} . On the left panel, the pink contour indicates the parameter space excluded by SK when the two RPV couplings are set to the same value, while the dashed line shows the projected SHiP sensitivity reach [32]. On the right panel, the green contour indicates the excluded parameter space when $\lambda'_{121} \neq 0$ and $\lambda'_{112} = 0$. Solid lines indicate constraints from neutral kaon oscillations [96] in green and DY processes at the LHC (labeled as ‘Colliders’) [97] in black, both evaluated at a sfermion mass of 1 TeV (for details, see text).

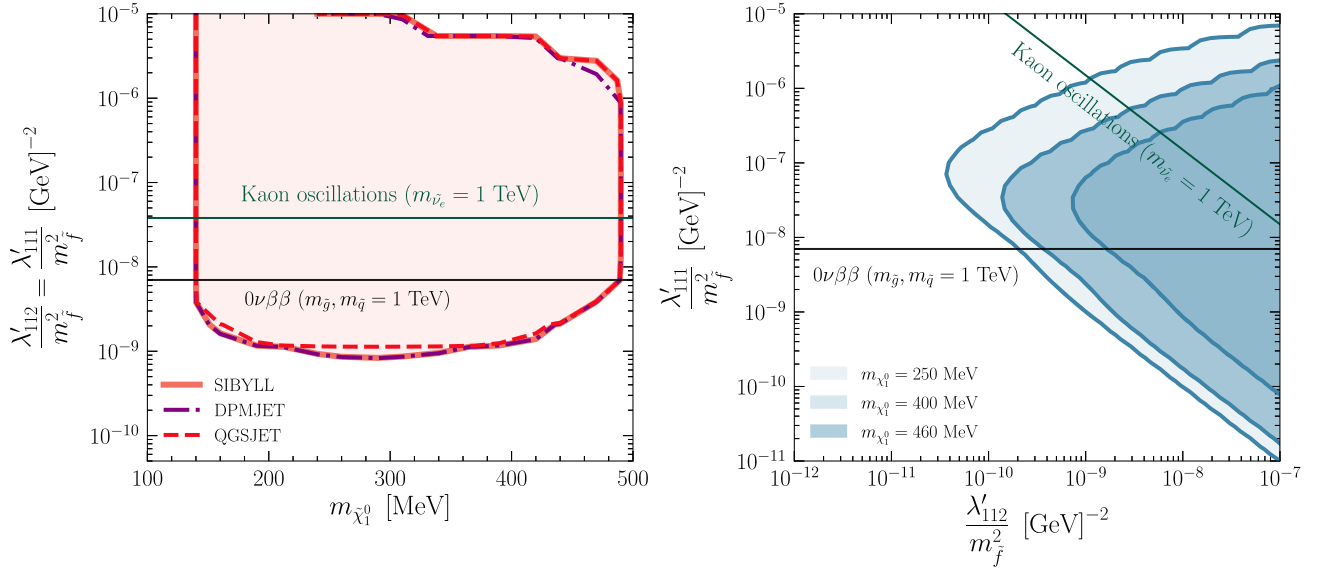


FIG. 6. *Limits for benchmark B2.* Limits derived at 90% confidence level in the parameter space defined by the neutralino mass $m_{\chi_1^0}$ and the trilinear RPV parameters λ'_{112} , λ'_{111} . On the left panel, the red contours indicate the parameter space excluded by SK when the two RPV parameters included in B2 have the same value. The different types of lines are used to distinguish the contours obtained with three different event generators, SYBILL-2.3, QGSJET-II-04, and DPMJET-III. On the right panel, the SYBILL-2.3 hadronic interaction model was used to generate three blue contours that indicate the parameter space excluded by SK when both RPV parameters vary freely. The different shades of blue represent different values of the neutralino mass. Solid lines indicate constraints from kaon oscillations [96] in green and neutrinoless double beta decay ($0\nu\beta\beta$) [99] in black. The bounds are evaluated for degenerated sfermion masses set to a value of 1 TeV.

Super-Kamiokande all the neutralino decay products listed in Table II are visible as e -like events, we can still have a signal when the parameter λ'_{112} is set to zero and we are left with neutral kaons in the final states. In such a case, the limits from kaon oscillations do not apply (see Fig. 5, right panel), and we are left with the collider constraint from the DY dilepton process $pp \rightarrow \ell^+ \ell^-$ for the parameter λ'_{121} , which is given by [97]

$$\lambda'_{121} \leq 0.34 \frac{m_{\tilde{q}}}{1 \text{ TeV}} + 0.18, \quad (19)$$

where $m_{\tilde{q}}$ is the squark mass. In this case, we find that the limit obtained from Super-Kamiokande data is better than the current constraint for the corresponding mass range.

On the other hand, in the case of benchmark B2, there is also a stringent limit from kaon oscillations on the product of the parameters, namely [96]

$$|\lambda'_{111} \lambda'_{112}| \leq 1.5 \times 10^{-3} \left(\frac{m_{\tilde{\nu}_e}}{1 \text{ TeV}} \right)^2. \quad (20)$$

Nevertheless, the most stringent constraint comes from its contribution to $0\nu\beta\beta$. The current limit for the half life of this process as determined by the KamLAND-Zen experiment [98], imposes the upper limit [99]

$$\lambda'_{111} \leq 2.2 \times 10^{-3} \left(\frac{m_{\tilde{q}}}{1 \text{ TeV}} \right)^2 \left(\frac{m_{\tilde{g}}}{1 \text{ TeV}} \right)^{1/2}, \quad (21)$$

where $m_{\tilde{q}}$ and $m_{\tilde{g}}$ are the masses of the squarks and gluinos, respectively. To contrast these bounds with our results, again we set both mass parameters to the value of 1 TeV (see Fig. 6). Remarkably enough, due to the higher kaon flux with respect to D -mesons, the limits from Super-Kamiokande in this parameter space turn out to be more stringent than the bounds that come from both kaon oscillations and $0\nu\beta\beta$. The excluded parameter space for this benchmark scenario changes marginally with different choices of hadronic interaction models, as it can be seen from the lines that indicate the exclusion region obtained with each model (see the left panel of Fig. 6). The difference between the kaon flux obtained with EPOS-LHC and with our benchmark model SYBILL-2.3 does not translate to any visible impact in the excluded parameter space, therefore we omit its contour in the figure. As a caveat, we note that the bound shown in Eq. (21) relies upon the gluino dominance assumption, where the neutralino mass is a fraction of the order of 10^{-2} times the gluino mass [59].

Finally, we emphasize that our results do not depend on the value of the sfermion masses. However, the bounds shown in Eqs. (17)–(21) become more strict as sfermion masses increase and less strict when they are lowered.

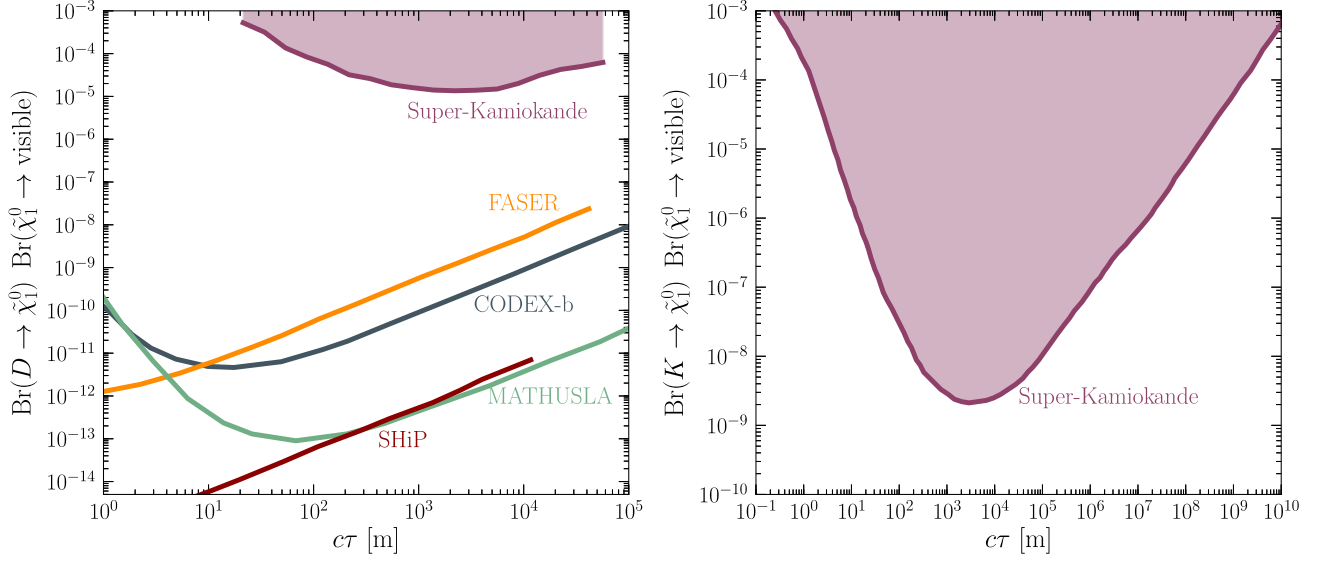


FIG. 7. *Comparison with lifetime frontier experiments.* Limits at 90% confidence level in the plane defined by the total branching ratio and the proper decay length of the long-lived neutralino. On the left panel, the purple contour shows the limits for SK derived in this article for benchmark B1, and the colored lines displays the projected sensitivity reach at future experiments for the same benchmark scenario derived in references [30,32]. On the right panel, the limit obtained using Super-Kamiokande data for the benchmark scenario B2 is shown. There are no other studies for this case in the literature.

Moreover, in the case where the squarks and sneutrino masses are kept fixed and the gluino mass is increased, the bound from $0\nu\beta\beta$ becomes less stringent.

The excluded combinations of values for the neutralino masses and RPV couplings allow us to determine the region of the parameter space that can be probed in the plane defined by the total branching fraction, which is the production branching ratio of the neutralino times the branching ratio of its decay to a visible signal, and the lifetime of the neutralino. The results can be seen in Fig. 7. Given that for benchmark B1 we are considering the same production and decay channels as in references [30,32], we can compare the current region that can be excluded by Super-Kamiokande, with the sensitivity reach expected for FASER, and other possible future experiments including CODEX-b, MATHUSLA and SHiP (left panel of Fig. 7).

In the case where the neutralinos are produced in kaon decays, there are no studies on the expected sensitivity reach at the next generation of detectors. The results obtained in this case demonstrate the great capacity of the Super-Kamiokande neutrino detector to place stringent limits for long-lived particles, owed to its ability to probe particles with lifetimes with a peak sensitivity around 10 km and total branching fractions of order 10^{-9} .

V. CONCLUSIONS

The lifetime frontier has emerged as a powerful line of exploration to search for beyond the Standard Model physics, specially in the absence of new signals at the LHC. R -parity-violating supersymmetry with light, long-lived

neutralinos constitutes a well-motivated scenario for new physics, with a rich phenomenology that has been gathering attention in recent years. In this work, we demonstrate a new way to search for long-lived neutralinos with masses of order 0.1 GeV–1 GeV that could be produced from the decay of charged D -mesons and kaons in cosmic-ray air showers, and whose visible decay can take place within large neutrino detectors such as Super-Kamiokande. We have analyzed two benchmark scenarios that include the couplings λ'_{121} , λ'_{112} , and λ'_{111} . In both cases, it is possible to improve the excluded region in parameter space when compared to existing constraints from colliders and neutrinoless double beta decay. Note, however, that this comparison requires to fix the mass of the sfermions, which we set to 1 TeV. An interesting feature about benchmark B1 is that in this case it is possible to compare the lifetime range that can be probed using SK data, against the expected sensitivity reach of next-generation, long-lived particle detectors. We find that the sensitivity for Super-Kamiokande peaks for lifetimes of the order of 1.0 km. In the case of the benchmark scenario B2, the advantage is twofold. On the one hand, there are currently no searches for light neutralinos produced from kaon decays. On the other hand, since kaons have a considerably higher production rate in air showers (as it can be inferred from Fig. 2), we find that it is possible to probe neutralinos with lifetimes of the order of hundreds of kilometers, while also being able to achieve a limit in $\lambda'_{111}/m_{\tilde{f}}^2$ and $\lambda'_{112}/m_{\tilde{f}}^2$ of order 10^{-9} GeV $^{-2}$, assuming these RPV couplings are equal.

When possible, we quantified the uncertainty that results from the choice of different hadronic interaction models. This was not possible for the benchmark scenario B1, since SYBILL-2.3 is the only event generator that provides a state of the art simulation of charmed mesons. For the case of the benchmark scenario B2, we find no strong dependence on the hadronic interaction model chosen to simulate the production of kaons in the atmospheric shower.

Overall, the results found in this study demonstrate the potential of large neutrino detectors to place limits on beyond the SM scenarios predicting long-lived particles, specially when contrasted with supersymmetric searches at colliders, where signals must be carefully chosen, and the reinterpretation of results is usually a complicated task. Finally, we stress that further scenarios can be

pursued systematically along the direction presented in this work.

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Millicharged particles from the heavens: single- and multiple-scattering signatures

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ABSTRACT: For nearly a century, studying cosmic-ray air showers has driven progress in our understanding of elementary particle physics. In this work, we revisit the production of millicharged particles in these atmospheric showers and provide new constraints for XENON1T and Super-Kamiokande and new sensitivity estimates of current and future detectors, such as JUNO. We discuss distinct search strategies, specifically studies of single-energy-deposition events, where one electron in the detector receives a relatively large energy transfer, as well as multiple-scattering events consisting of (at least) two relatively small energy depositions. We demonstrate that these atmospheric search strategies — especially the multiple-scattering signature — provide significant room for improvement beyond existing searches, in a way that is complementary to anthropogenic, beam-based searches for MeV-GeV millicharged particles. Finally, we also discuss the implementation of a Monte Carlo simulation for millicharged particle detection in large-volume neutrino detectors, such as IceCube.

KEYWORDS: Beyond Standard Model, Neutrino Physics

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1 Introduction

Since the earliest times, humans have stared at the sky and wondered. By observing the low-orbit Earth skies, they discovered the presence of extraterrestrial charged particles, now called cosmic rays [1–4]. These experimental efforts started from attempts to understand the origin of environmental radioactivity, and through their study led to the discovery of muons in 1936 [5] and pions in 1947 [6]. In this work, we return to contemplating cosmic-ray air showers’ products in order to shed light on the nature of electric charge by looking for signatures of small electrically charged particles, often known as millicharged particles [7].

In the Standard Model (SM), all electric charges are multiples of the d -quark charge, a principle known as charge quantization. However, though the SM imposes charge conservation by its gauge symmetries and anomaly cancellation, there is no firm theoretical evidence for the principle of charge quantization [7]. There are two ways in which this principle can be tested: by searching for small deviations between proton and positron charges or by looking for particles with small electric charges, $\varepsilon \equiv q/e \ll 1$. In this work, we will follow the second approach, but first we briefly summarize the leading results on these two directions.

The introduction of quarks [8] prompted the search for fractional charge particles between the 1960s and the early 1980s [9], an effort diminished by the discovery of color confinement [10]. It was suggested in ref. [8] that quarks, produced by cosmic-ray interactions in the Earth's surface [11, 12], could be detected by examining the electrical neutrality of atoms and of bulk matter [9, 13]. This led to constraints on the number of free quarks per unit mass [14] and constraints on the difference between proton and electron charges [15], $\eta = |q_p + q_e|/q_e$. In addition to suggestions from fundamental particle physics, it was pointed out in ref. [16] that the expansion of the universe could be accounted for by a small difference between the electron and proton charges at the level of $\eta \sim 10^{-18}$. However, soon after this proposal, laboratory experiments using de-ionized gas constrained it to be $\eta < 2 \times 10^{-20}$ [15]. Currently, constraints on this quantity using diverse methods — such as gas efflux, acoustic resonators, Millikan drop style experiments, and atomic and neutron beams, among others [9] — have limited η to be less than 10^{-21} [17].

Complementary to the above, direct searches for particles with small charge and sub-GeV mass have been motivated by dark matter models [18, 19] and cosmological puzzles [13]. In recent years, more attention has been given to these so-called millicharged particles (MCP) in the MeV-GeV mass regime. Searches for this scenario have been proposed for beam-based neutrino experiments [20, 21] and for dedicated experiments situated near accelerator complexes [22–27]. Existing data have been reanalyzed in this context in refs. [20, 28]. Recently, the first dedicated neutrino-experiment analysis for MCP in this mass range was carried out by the ArgoNeuT collaboration [29], setting the strongest constraints for a range of MCP masses and demonstrating the capability of neutrino experiments for these searches for decades to come.

In tandem with accelerator-based searches, atmospheric-based production of MCP has been proposed [30] (see also refs. [31, 32] which considered even more distant production mechanisms that can contribute to these searches), with large-volume neutrino detectors like Super-Kamiokande serving as the best candidates to search for these particles. Ref. [30] demonstrated that atmospheric MCP searches can be as or more powerful than beam-based ones. We build on this previous work, revisiting calculations of MCP production, discussing the uncertainty on the MCP flux, and proposing further searches that can be done with atmospheric MCP. Motivated by ref. [21], we explore multiple-hit signatures in which a given MCP traversing a detector can scatter off two or more electrons, leaving a faint track. This search is highly advantageous in detectors that can identify low-energy electrons, such as the upcoming multi-kiloton-scale unsegmented liquid-scintillator neutrino experiment JUNO [33, 34].

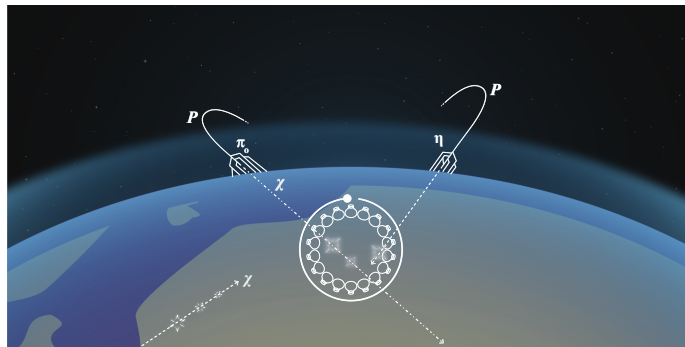


Figure 1. *Artistic rendition of this work.* Millicharged particles, χ , are produced in cosmic-ray showers from neutral meson decay. They lose energy as they travel through the Earth, with only the highest energy ones reaching the detector preferentially from directions of small overburden. These produce electromagnetic signatures in underground neutrino and dark matter detectors. In this work, we study both single- and double-hit signatures and relate their sensitivities.

Figure 1 offers an artistic rendering of this work’s main ideas. High-energy SM particles, like protons, bombard the atmosphere, producing rich particle showers. If MCP exist, then they can emerge in these showers and travel through the Earth. Large-volume detectors provide excellent targets for these MCP, which can scatter once (right track), potentially imparting enough energy on the target electron to be a strong signal in these detectors. If the MCP scatters multiple times (left track), the faint track it provides is difficult for background processes to mimic. Some flux of MCP could also travel through the Earth (bottom track), if its mean free path through the Earth is long enough, and come upwards through the detector.

The remainder of this work is organized as follows: in section 2 we revisit previous simulations of atmospheric MCP production and discuss the tools we use for our simulation. Section 3 discusses MCP propagation through the Earth, including energy loss and possible absorption, leading to an attenuation of the upward going MCP flux passing through a detector. In section 4 we provide the details of our experimental simulations for both current and upcoming experiments and construct sensitivity estimates for both the single-scattering and multiple-scattering searches. Section 5 discusses some Monte Carlo techniques for searches in even larger detectors, such as IceCube and the upcoming IceCube Upgrade. Finally, in section 6 we offer some concluding remarks.

2 Millicharged particle production

A careful consideration of the production of millicharged particles in the upper atmosphere is crucial to this proposed search strategy. We provide the details of this approach with respect to its formalism in section 2.1. Several sources of uncertainty are relevant as well, and these uncertainties, unfortunately, can plague any search for millicharged particles that relies on their production from the decays of neutral pseudoscalar/vector mesons. We discuss these uncertainties in section 2.2.

We make a standard assumption regarding the nature of this new millicharged particle — that it has some mass and small coupling to the Standard Model photon via a small electric charge. It is possible that a particle of such nature constitutes some fraction of the Dark Matter in the universe, or even an intermediate state in which Dark Matter can decay [35–37]. In the former case, a number of additional constraints apply. This has been explored as a potential solution to the EDGES anomaly [38, 39] in, for instance, refs. [40–45]. Various effects of millicharged particles as dark matter, leading to stringent constraints in the MeV-GeV mass range, are discussed in refs. [31, 32, 46–50]. If a millicharged particle is discovered via the atmospheric-production and scattering-detection approach we propose, then it is *at most* a very small fraction of the relic dark matter in the universe. After such a discovery, a great deal of scrutiny must be applied to determine a consistent picture of this new particle with cosmological and astrophysical observations.

2.1 Formalism

Cosmic rays are a population of energetic elementary particles and nuclei. Their collision with the upper layers of Earth’s atmosphere mimics the setup encountered in a proton beam dump experiment, with nuclei in the air playing the role of the target. An extensive cascade of radiation, ionized particles, and hadrons is generated with energies ranging from a few GeV up to dozens of EeV [51]. For the energy range of interest for our work, $\sim$ GeV-TeV, the cosmic ray spectrum is primarily composed by protons and helium, both of which are well-measured by the AMS experiment [52]. Among the mesons produced in the cascade, it is possible to find pseudoscalar mesons — such as π^0 , η — and vector mesons — such as ρ , ω , ϕ and J/Ψ . If a millicharged particle exists with a mass below half of any given meson mass, then it can be produced in two- or three-body decays, replacing the final-state electrons in relatively common processes such as $\pi^0 \rightarrow \gamma e^+ e^-$ and $J/\Psi \rightarrow e^+ e^-$. The η' meson could also contribute to this flux, however its contributions are suppressed relative to η and heavier meson production. At low MCP masses, η' production is suppressed relative to η due to the lower production rate of η' in the atmosphere. At higher masses, it is suppressed relative to heavier meson production, such as from ω , due to the lower branching ratio of η' into MCP. We have validated this suppression with our simulations and therefore do not include η' production in what follows. In principle, heavier mesons such as Υ and direct, Drell-Yan production of MCP could allow for searches for particles above $m_{J/\Psi}/2$. However, we expect that the production rate of MCP at these masses will be limited and we would not expect much sensitivity from these production mechanisms.

All of these mesons are unstable and have very short lifetimes and although none of them reaches the surface of the Earth, they can decay via a photon-mediated process to millicharged particles. The MCP are assumed to be stable particles and can reach the surface of the Earth and propagate through it¹ in such a way that they could be detected in underground detectors such as neutrino and dark matter experiments.

¹Despite having charge, the interaction rate of MCP with the Earth is small enough that, for most of the parameter space of interest, they can travel through the bulk of the Earth without significant energy loss. We discuss this effect and how we include it in our simulation in section 3.

In this work, we will adopt a model where the MCP (χ) is described by a stable particle coupled to the photon with strength $\varepsilon \times e$ and mass m_χ . Taking this into account, the production profile of millicharged particles generated in air-showers from a parent meson $\mathbf{m}$, can be described by the cascade equation [53]

$$\frac{d\Phi_\chi}{dE_\chi d\cos\theta dX} = \int dE_\mathbf{m} \frac{1}{\rho(X)\lambda_\mathbf{m}(E_\mathbf{m})} \frac{d\Phi_\mathbf{m}}{dE_\mathbf{m} d\cos\theta}(E_\mathbf{m}, \cos\theta, X) \frac{dn}{dE_\chi}(\mathbf{m} \rightarrow \chi; E_\mathbf{m}, E_\chi), \quad (2.1)$$

where $\rho(X)$ is the atmospheric density at column depth X , $\lambda_\mathbf{m} = \gamma_\mathbf{m}\beta_\mathbf{m}c\tau_\mathbf{m}$ is the decay length of the parent meson $\mathbf{m}$, and $\frac{d\Phi_\mathbf{m}}{dE_\mathbf{m} d\cos\theta}$ is the production rate of the meson with energy $E_\mathbf{m}$ at zenith angle θ . Here, $\frac{dn}{dE_\chi}$ is the energy distribution of the millicharged particle in the decay, which can be written in terms of the branching fraction and the decay rate distribution as

$$\frac{dn}{dE_\chi}(\mathbf{m} \rightarrow \chi) = \text{Br}(\mathbf{m} \rightarrow \chi) \times \frac{1}{\Gamma(\mathbf{m} \rightarrow \chi)} \frac{d\Gamma}{dE_\chi}(\mathbf{m} \rightarrow \chi). \quad (2.2)$$

The branching fraction $\text{Br}(\mathbf{m} \rightarrow \chi)$ and decay width $\Gamma(\mathbf{m} \rightarrow \chi)$ must be specified accordingly for the three-body decays of the pseudoscalar mesons and the two-body decay of the vector mesons. The expressions for these quantities, as well as the kinematic integration limits in eq. (2.1), are specified in appendix A. The production rate of the mesons in the atmospheric cascade can be solved numerically. We have used the Matrix Cascade Equation (MCEq) software package [54, 55], which includes several models for the cosmic ray spectrum, hadronic interactions, and atmospheric density profiles. In this work, for our benchmark results we have used the SYBILL-2.3 hadronic interaction model [56], the Hillas-Gaisser cosmic-ray model *H3a* [57], and the NRLMSISE-00 atmospheric model [58] to obtain the production rate of mesons, which have been extracted using the same procedure of references [59] and [60].

Given that millicharged particles are stable, a direct integration of eq. (2.1) in the column depth X can be performed to find the differential energy spectra for a fixed zenith angle θ . In figure 2, we show the energy and angular dependence of the MCP flux at the surface of the Earth, for each one of the parent mesons considered in this work.

2.2 Uncertainties

The uncertainties in the production of MCPs are mostly due to the cosmic-ray model (CRM) chosen to generate the primary spectrum, as well as the hadronic interaction model (HIM) used to simulate the production rate of mesons in the atmospheric shower, the latter providing the dominant source of uncertainties, which grow as the energy increases. Ref. [61] recently explored these uncertainties in a fashion similar to how we have done here.

The composition and energy spectra of the primary cosmic radiation are characterized by a CRM that describes the primary spectrum with a power law that is ultimately fitted to air-shower data. The power law spectrum may break or not depending on the origin of the cosmic ray and the energy range observed. Additionally, the spectrum is typically characterized by the steepening that occurs for proton energies between 10^6 and 10^7 GeV, the so-called “knee,” and an extra feature around 10^9 GeV called the “ankle.” The origin

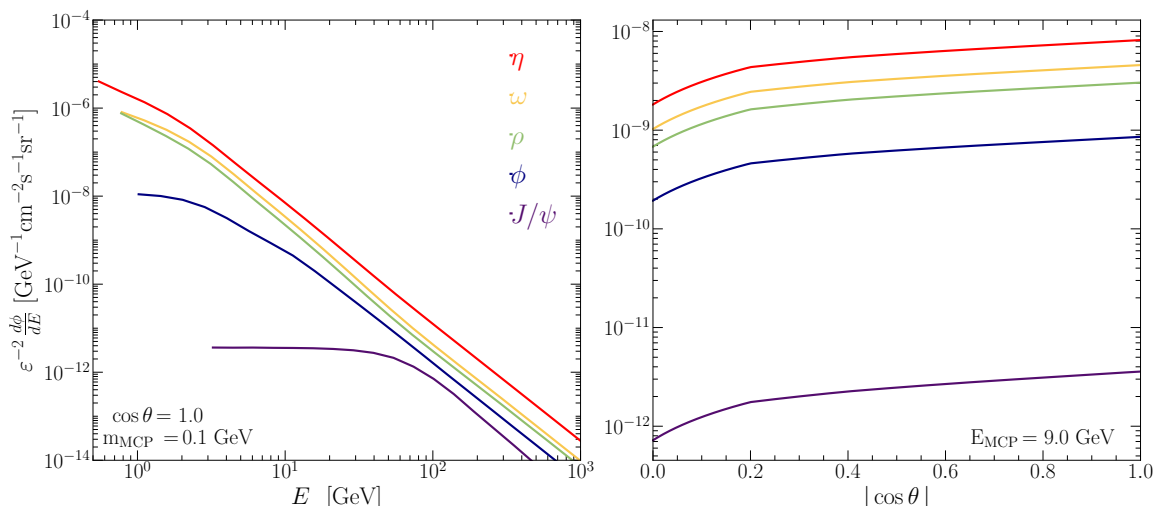


Figure 2. *Surface flux of millicharged particles from different parent mesons.* Left: spectra, which scales as ε^2 , of millicharged particles arriving at the surface (i.e., before any potential attenuation/deflection) of the Earth with zenith angle 0° , for each one of the different parent mesons. Right: angular dependence of the flux at 9 GeV. In both panels the MCP mass is set to 100 MeV.

of these features remains unclear, and it is an active research topic (see, for instance, refs. [62, 63]).

As mentioned before, the benchmark CRM we use to compute the MCP flux is the *H3a* model, which is widely used for calculations of atmospheric lepton fluxes [56, 64]. However, there are other realistic CRMs which could be considered and that would yield to a more optimistic/pessimistic estimate of the MCP flux. To illustrate this point, we have considered the production rate of the π^0 meson in other CRMs, such as the Thunman-Ingelman-Gondolo model [53], the Gaisser-Stanev-Tilav model [65], and the poly-gonato model [66].² The comparison of the differential production rate for these other CRMs as well as the ratio with respect to our benchmark model *H3a* as a function of the energy, can be seen in the right panel of figure 3.

On the other hand, to model the interactions of a primary cosmic ray with the atmosphere, we need a suitable model of hadron-hadron, hadron-nucleus, and nucleus-nucleus collisions. Yet, hadronic collisions at very high energies involve the production of particles with low transverse momenta, where present theoretical tools such as QCD are not enough to understand this feature. To address this problem, phenomenological models with Monte Carlo implementations are used. Since a hadronic interaction model provides the interaction coefficients of the coupled cascade equation used to describe the production rate of mesons, we can directly prove its impact by estimating the production rate, for a fixed CRM. In this work, we choose SYBILL-2.3c as benchmark model, and the QGSJET-II-02 [67], DPMJET-III [68], and EPOS-LHC [69] HIMs for comparison. The left panel of figure 3 displays the production rate of π^0 for all of these models as well as the ratio to our benchmark HIM.

²This last model is however not applicable at energies above the “knee”.

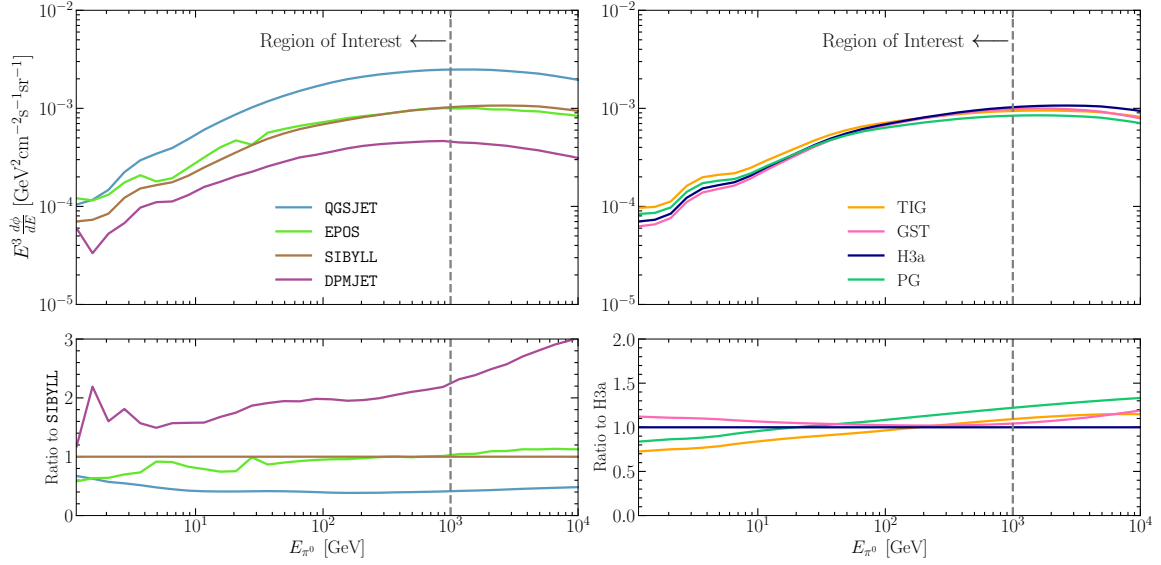


Figure 3. *Production Rate of π^0 in cosmic showers and variations in different HIMs and CRMs.* Left: comparison of different hadronic interaction models, including QGSJET (blue), EPOS (green), SIBYLL (brown), and DPMJET (purple). The bottom panel depicts the production ratio relative to SIBYLL.2.3.c, which we use for our production estimates. Right: comparison of production rates for different cosmic ray models: TIG (orange), GST (pink), H3a (blue), and PG (green). The bottom panel shows the ratio relative to H3a, which we use for our production estimates. See text for further discussion of CRMs and HIMs.

To estimate the impact of the benchmark models used in this work, we evaluate the ratio of total production with respect to a different CRM/HIM, defined by

$$\Delta_j(\mathbf{m}) = \frac{\int_{E_{\min}}^{\Lambda} dE \frac{d\phi_{\text{BM}}}{dE}}{\int_{E_{\min}}^{\Lambda} dE \frac{d\phi_j}{dE}}, \quad (2.3)$$

where $\mathbf{m}$ is the meson of interest, j the index used to denote the model that is being compared with the benchmark model BM, $E_{\min} = 1.66$ GeV is the minimum energy available in MCEq, and $\Lambda = 10^3$ GeV is the upper energy cut that we use in order to obtain the MCP flux from a parent meson. Table 1 shows the Δ coefficients for the different CRMs and HIMs for π^0 , η , and ϕ . A similar result can be found for the other mesons, except for J/Ψ , in which case there are not enough statistics on the meson production rate to evaluate the uncertainty properly. As can be seen from this result, the biggest source of uncertainty comes from the hadronic interaction models which would induce a difference in the MCP flux from about $\sim 16\%$ up to $\sim 68\%$, depending on the parent meson.

Even though all of the event generators considered here are updated with LHC data, the cross-section measurements for hadron production have rather large uncertainties [70]. On the other hand, we must also take into account the differences encountered in the phenomenological models that arise from the treatment of inelastic hadronic collisions within the framework of Reggeon Field Theory [71]. We stress that HIM uncertainties

Model	$\Delta(\pi^0)$	$\Delta(\eta)$	$\Delta(\phi)$
TIG	0.754	0.766	0.827
PG	0.868	0.88	0.944
GST	1.106	1.101	1.072
QGSJET	0.570	1.160	—
DPMJET	1.586	1.677	1.541
EPOS	0.664	0.707	0.750

Table 1. Comparison of meson production rate using different hadronic interaction models. Numerical values of Δ , the relative integrated production rate of a given meson, for the various cosmic ray models (top three rows) and hadronic interaction models (bottom three rows) considered in this work.

are present in most if not all modern MCP searches, as any beam-based³ search needs to simulate hadronic interactions at some level.

As a final remark, we note that for low CR energies, there are geomagnetic and solar modulation effects which imply an additional source of uncertainty. Quantifying this effect in a precise way lies beyond the scope of this work, but this may be included in future experimental analyses. We note that ref. [75] has shown that these effects are small ($\lesssim 10\%$) for CR energies above 10 GeV, which constitute the important energy range for MCP production in this work.

3 Propagation through Earth and energy loss

As we will show in section 4, underground neutrino detectors provide an excellent opportunity to search for MCPs produced in cosmic-ray air showers. However, the detection of MCPs requires detailed knowledge about charged particle propagation in the medium, since these experiments use the Earth’s crust to shield from atmospheric muons, which usually constitutes the main source of background. Because of this, it is expected that the MCPs that arrive at the detector lose part of the energy they had when reaching the Earth’s surface. Just as with any other charged particle, the MCP would lose energy by ionizing the medium through which they propagate and by interacting with nuclei. The average energy loss along the MCP trajectory can be parametrized by [17]

$$-\frac{dE}{dX} = \varepsilon^2 \left(a_{\text{ion.}} + b_{\text{el.-brem.}} \varepsilon^2 E + b_{\text{inel.-brem.}} E + b_{\text{pair}} E + b_{\text{photo-had.}} E \right) \approx \varepsilon^2 (a + bE), \quad (3.1)$$

where dX is the column density traversed, ε is the MCP coupling, and a_x and b_x are various categories of energy loss with (potentially) different scaling with the MCP energy E and ε . We simplify this expression, adopting the right-hand side of eq. (3.1) in our simulations, where a and b are the energy loss parameters given in units of [GeV/mwe] and

³Notable exceptions include the SLAC mQ experiment [72] and NA64 [73, 74], which consider production of MCP in an electron beam dump. Such production mechanisms would be subject to far smaller uncertainties.

[mwe] respectively (“mwe” being “meters of water equivalent”). Generally, b is inversely proportional to the mass of the charged particle. We take these parameters as measured for muons [76] in our simulations and neglect mass-dependence of these terms in our simulations, leading to a conservative estimate of the particle absorption during propagation for MCP masses above m_μ . For masses below m_μ , the above simplification yields results only marginally stronger than if those were included. This is because the ionization term (a) dominates in our energy range, and the radiative losses term (b) is subleading. The latter is expected to have a linear dependence on the MCP mass; namely, it scales like m_{MCP}/m_μ [77]. For the masses considered in this analysis, the latter contribution, even accounting for the mass correction, is negligible, and for simplicity, we ignore it. We can estimate the overburden length traversed by an MCP approaching a detector located at depth d and along a trajectory with angle θ by

$$D = \sqrt{(R_\oplus - d)^2 \cos^2 \theta + d(2R_\oplus - d)} - (R_\oplus - d) \cos \theta, \quad (3.2)$$

with $R_\oplus$ being the Earth’s radius.

The probability that a millicharged particle with energy at the surface E_i arrives at the detector with an energy E_f will depend on the coupling $\alpha_{em}\varepsilon$ and the incident direction $\cos \theta$. This probability is given by:

$$\mathcal{P} = \exp \left(-\frac{D(\cos \theta)}{R(E_i, \varepsilon, E_f)} \right), \quad (3.3)$$

where the average distance $R(E_i, \varepsilon, E_f)$ can be obtained from eq. (3.1), and is given by

$$R = \frac{1}{\varepsilon^2 b} \ln \left(\frac{1 + \frac{a}{b} E_i}{1 + \frac{a}{b} E_f} \right). \quad (3.4)$$

We have taken the energy loss parameters for a standard rock shielding with $a = 0.223 \text{ GeV/mwe}$ and $b = 4.64 \times 10^{-4} \text{ mwe}$ as reported in ref. [77]. The standard Earth matter density profile is used as well [54, 55]. The left panel of figure 4 shows the probability that an MCP arrives at a detector at a depth of 1.5 km with $E_\chi > 1 \text{ GeV}$ for two different values of the coupling ε , with a variety of different initial energies as depicted in the label. The flux at the detector can be obtained by convolving the flux at the surface with the survival probability $\mathcal{P}$. The right panel of figure 4 shows the angular distribution of the flux at detector for an MCP with a mass of 100 MeV arriving at the detector with energies of 1.0 and 2.1 GeV.

We note here that for relatively large $\varepsilon^2 = 10^{-2}$, the upward-going flux, $-1 \leq \cos \theta \leq 0$ is highly attenuated. On the other hand, for smaller $\varepsilon^2 = 10^{-5}$, because energy losses are suppressed by two orders of magnitude, the flux is mostly independent of $\cos \theta$. For experiments that can be sensitive to such small millicharges, the effective area of the sky that the detector can search will nearly double. Moreover, searches for upward-going events can assist in reduction of background events, for instance from atmospheric muons entering the detector from above.

Finally, as detailed in ref. [21], there is the possibility that the same scattering process that we used to determine MCP absorption while travelling through the Earth can lead

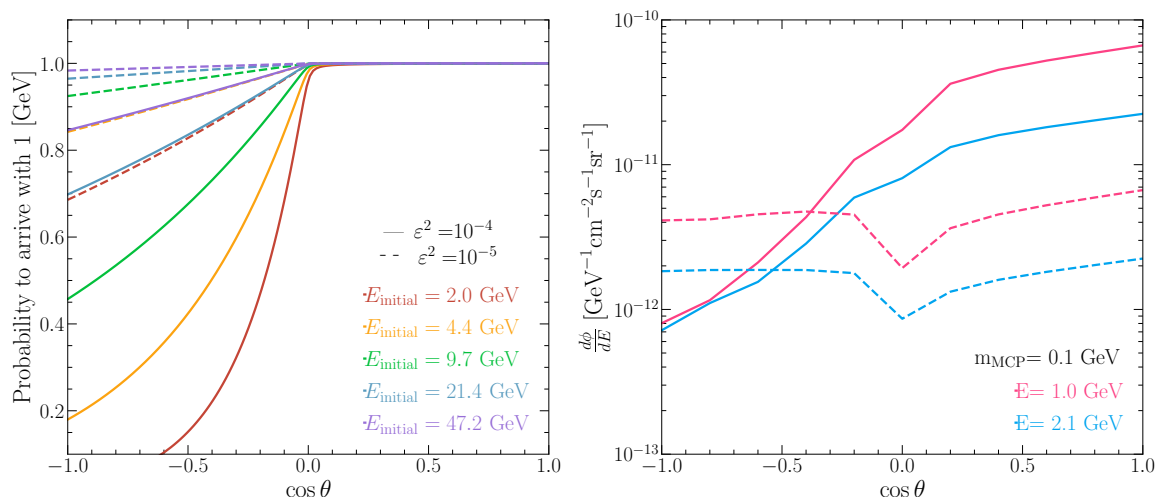


Figure 4. *Survival probability of MCPs and expected angular distribution.* Left: probability that an MCP with initial energies as indicated by the labels arrives at the detector with a total energy of 1 GeV as a function of the cosine of the zenith angle. Solid lines are shown for $\epsilon = 10^{-4}$ and dashed lines for $\epsilon = 10^{-5}$. Right: flux at detector as a function of cosine of the zenith angle for a 100 MeV mass MCP, arriving with an energy of 1.0 GeV (pink) and 2.1 GeV (blue). Solid and dashed lines corresponds to the same values of ϵ as in the left panel.

to scattering and deflection of MCP particles. Using the arguments of ref. [21], we find that for MCPs travelling through the bulk of the Earth (path lengths of $\mathcal{O}(10^4)$ km), the typical deflection is less than $\sim 10^\circ$ for $\epsilon^2 \lesssim 10^{-4}$ and 1° for $\epsilon^2 \lesssim 10^{-6}$. Our simulations, which include the energy losses discussed cf. figure 4, account for the “absorption” of MCP that scatter, but not the angular deflection, which can smooth out the flux with respect to $\cos \theta$ a small amount (without changing the integrated flux). This effect could be important for analyses that use the $\cos \theta$ distribution in the statistical analysis, however, the ones we perform all integrate over this information. This effect would be most pronounced for large ϵ^2 where many deflections happen, and we do not expect it to have a large impact on our resulting experimental sensitivity to MCP.

4 Current and future experimental searches

To date, the most stringent searches for millicharged particles in the $\sim \text{MeV-GeV}$ mass range have used antropogenic beams, where the MCP are produced in either the collision of two beams (collider searches) or when a beam impacts a target (beam-dump searches). Colliders place the strongest constraint on MCPs with masses above ~ 1 GeV, primarily through a combination of direct searches at LEP and measurements of the invisible width of the Z boson [78]. The LHC can probe even larger, $\gtrsim 30$ GeV, masses [79]. Below 100 MeV, the strongest constraint is from the SLAC mQ electron fixed-target experiment [72]. Future dedicated experiments have been proposed in the context of both collider [22, 25]⁴ and

⁴Including prototype results found in refs. [80, 81].

fixed-target [24, 26, 27] environments, reaching sensitivity to $\varepsilon^2 \gtrsim 10^{-6}$ and a wide range of masses, up to nearly 5 GeV.

Recently, searches for MCP in neutrino experiments have garnered attention. Sensitivities for accelerator production have been estimated in refs. [20, 21] and carried out by the ArgoNeuT collaboration in ref. [29]. The latter has set some of the strongest constraints over a wide MCP range. Atmospheric production has been considered in a similar fashion to this work for the Super-Kamiokande detector and its future successor Hyper-Kamiokande [82] in ref. [30], as have limits from particles produced in the interstellar medium and from Earth relics recently discussed in refs. [31] and [32], respectively.

In the remainder of this section, we will discuss and derive constraints for MCP masses between 10 MeV and 1.5 GeV. Our results for current experiments, particularly Super-Kamiokande, demonstrate an improvement on current constraints for some masses, in agreement with the results of ref. [30]. We divide the search strategy into two different types of analyses, those relying on single-hit signals (section 4.1) and those that rely on multiple energy depositions from a single MCP particle traversing the detector (section 4.2). The former is advantageous for large-volume, high-energy-threshold experiments (e.g., water Cherenkov detectors like Super-Kamiokande), whereas the latter is strongest in scintillator experiments with precision timing and low-energy thresholds (for instance, JUNO). We find that these multiple-scattering searches offer great potential for discovering millicharged particles in the 10 MeV to 1.5 GeV mass range and warrant additional focus in the coming decade.

4.1 Single-scattering searches

Via the same coupling ε that allows for the production of MCP in the upper atmosphere, the particles that reach a detector are capable of scattering (via a t -channel photon exchange) off detector materials. We will consider scattering off electrons for the remainder of this work. This massless-mediator scattering yields a differential cross section that peaks for small electron recoil energy, and so detectors with capabilities of identifying and reconstructing low-energy electrons will be advantageous in this endeavor. However, as we consider electrons of lower and lower energy, more and more backgrounds become relevant. In the following subsections, we will discuss the characteristics of the signal events, the various backgrounds, and the experimental limits we are able to derive in this single-hit analysis including the statistical techniques employed.

4.1.1 Signal

Once the flux of MCPs arrives at the detector, the millicharged particles can interact with the detector medium by scattering off electrons. We will be interested in low-energy recoils, as those are more frequent due to the shape of the differential cross section. The differential cross section between MCPs and electrons is given by [20]

$$\frac{d\sigma}{dE_r} = \varepsilon^2 \alpha_{EM}^2 \pi \frac{m_e(E_r^2 + 2E_\chi^2) - E_r(m_e(2E_\chi + m_e) + m_\chi^2)}{E_r^2 m_e^2 (E_\chi^2 - m_\chi^2)}, \quad (4.1)$$

where E_χ is the energy of the incoming MCP, E_r is the recoil energy (assuming an initial stationary electron in the laboratory frame), α_{EM} is the electromagnetic coupling constant,

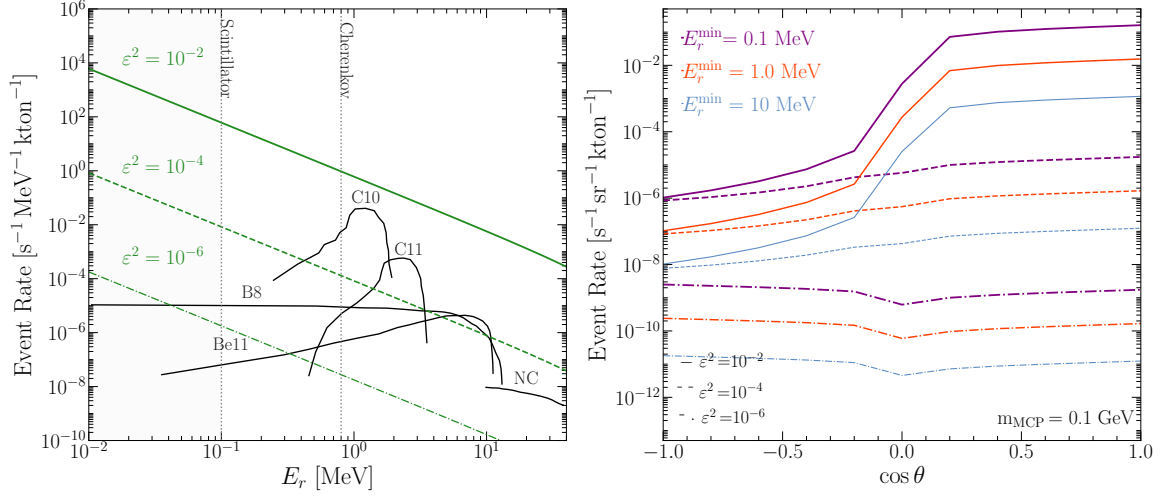


Figure 5. MCP event rate angular and energy distributions as well as relevant backgrounds for neutrino detectors. Left: expected event rate as a function of the electron recoil energy for MCPs coming in all directions. The vertical dotted lines indicate the different thresholds for scintillator and Cherenkov detectors. The solid black lines indicate the background expected in JUNO from different components (see background section for details). Right: expected event rate as a function of the incoming direction of the MCPs. Solid, dashed, and dot-dashed lines correspond to different values of ϵ^2 as depicted in the legend. The colors indicate the three different minimum electron recoil energies with values shown in the legend. Both figures are shown for an incident millicharged particle with a mass of 100 MeV.

and m_χ and m_e are masses of the MCP and the electron, respectively. From the equation above, it is easy to verify that in the ultrarelativistic limit $E_\chi \gg E_r, m_\chi, m_e$, the differential cross section scales like E_r^{-2} , and therefore as the threshold of a given experiment becomes lower, the signal becomes much stronger.

The event rate from MCPs can be obtained by a convolution of the millicharged particle flux with the differential cross section multiplied by the number of targets and the detection efficiency. This is given by

$$\frac{dN}{dE_r d\cos\theta} = n_e \epsilon(E_r) \int dE_\chi \frac{d\phi}{dE_\chi} \frac{d\sigma}{d\cos\theta dE_r}, \quad (4.2)$$

where n_e is the total number of electrons in the detector and $\epsilon(E_r)$ is the (detector-dependent) reconstruction efficiency of these single-electron events. The event rate distribution for a 1 kton effective mass detector is shown in figure 5, as a function of the incoming direction from zenith angle and the recoil energy, for different values of ϵ^2 and an MCP mass of 10 MeV. Notice that for the angular distribution an integration in the recoil energy is needed, and the values in the label indicate different detection thresholds. The MCP event rate shown in this figure is a general result, which can be applied to any underground particle detector located at a depth of ~ 1.5 km, while the mass dependence scales in a similar way as the MCP flux, with higher event rates at lower masses, where the millicharged particle receives contributions from most of the mesons.

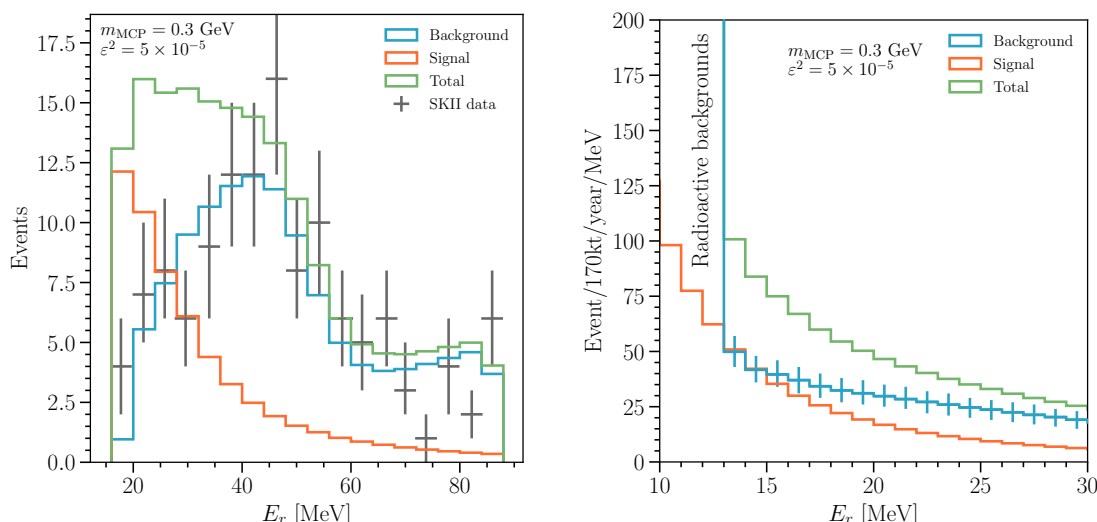


Figure 6. Event distribution illustrations for Super-Kamiokande and JUNO. Left: electron events measured in Super-Kamiokande are shown as crosses, total backgrounds are shown as a blue histogram, and a potential MCP hypothesis is shown in orange. Right: expected distribution of MCP events in JUNO overlaid with backgrounds. Note that we have chosen $m_\chi = 300$ MeV in this figure, in contrast to previous ones, to select a point in (m_χ, ϵ^2) parameter space that yields interesting event rates without being excluded by existing experiments.

4.1.2 Backgrounds

Here we discuss the different sources of background events that contribute to such single-scattering millicharged particle searches in water Cherenkov or liquid scintillator detectors. The expected background rates (summed over all contributions) in these two environments, along with a signal expectation, are shown in figure 6 for Super-Kamiokande phase II (left) and JUNO (right). More details on the Super-Kamiokande event distributions are given in appendix B.

Penetrating muons. Our atmosphere is filled with muons generated from the decays of charged mesons that are dominated by pions at the lowest energies and kaons at higher energies [83]. These muons are highly boosted, and even though they have a short lifetime ($\sim 2 \times 10^{-6}$ s), they can easily reach the surface, propagate through the Earth and leave a signal in the detector. The most efficient way to suppress this and other possible cosmogenic backgrounds is to have a sufficient amount of overburden over the location of the detector. This background may be rejected by measuring the opening angle of the Cherenkov cone of the candidate electron in the event [84].

Neutral-current neutrino events. Atmospheric neutrinos induce a source of background due to neutral-current interactions with nuclei in the medium. In the case of liquid scintillator detectors, the background is induced by interactions with ^{12}C whose de-excitation produces radiation. In Cherenkov detectors, elastic, neutral-current scattering of neutrinos off the target nuclei can lead to small energy depositions with a similar spectrum

to the millicharged particle signature. However, this is a relatively small component of the background in Cherenkov detectors like Super-Kamiokande.

Charged-current neutrino events. Specifically in water or ice Cherenkov detectors (i.e., Super-Kamiokande, Hyper-Kamiokande, or IceCube), both $\nu_e/\bar{\nu}_e$ and $\nu_\mu/\bar{\nu}_\mu$ charged-current scattering events from the atmospheric neutrino flux may mimic our desired signal of a low-energy electron. The former directly creates low-energy electrons indistinguishable from the signal as long as the hadronic particles produced in the interaction are below the Cherenkov threshold. Despite the fact that they cannot be separated on an event-by-event basis, the recoil energy distribution of the electrons has a spectral shape different than the MCP signal, which allows them to be distinguished statistically. Similarly, ν_μ events can mimic the background if the outgoing $\mu^\pm$ is below the Cherenkov threshold, but the Michel $e^\pm$ coming from the $\mu^\pm$ decay are visible. As in the previous case, the Michel electron energy distribution is distinct and can be disentangled statistically; see ref. [84] for further details.

Radioactive and anthropogenic backgrounds. Radioactive backgrounds are produced by the materials in and around the detector. These can produce electrons with energies similar to those struck by MCPs, e.g., in radioactive beta decays of nuclei, and depend on the experimental setup. For example, the dominant radioactive backgrounds in JUNO, in the energy range relevant for this analysis, are ^8B , ^{10}C , ^{11}C , and ^{11}Be , which we reproduce from ref. [33] and include as background in the right of figure 6. We note that this rate is significantly larger than currently allowed MCP signatures for $E_e \lesssim 12\text{ MeV}$ and will be the limiting factor for single-scattering searches in detectors like JUNO. Searches for signals with an associated neutron in the final state, most prominently scattering of $\bar{\nu}_e$ from the diffuse supernova neutrino background, can reject these backgrounds by requiring coincidence with the signal of neutron capture. Since MCP signatures do not produce an outgoing neutron, but only the electron recoil, this method cannot be used to reduce the background. However, as we will discuss in section 4.2, multiple-hit signatures can be used to reduce radioactive backgrounds. The dominant anthropogenic backgrounds in the search of MCPs are produced by either nearby nuclear reactors or by accelerator facilities. In JUNO, the antineutrinos produced by nuclear reactors are the source used to study neutrino oscillations and can be separated from the MCP signature by the presence of the coincident neutron capture mentioned above. In Super-Kamiokande, accelerator neutrino events produced in the Tokai accelerator facility can be removed by searching for MCPs during off-beam times.

4.1.3 Experimental setups considered

In this section, we will briefly describe the experiments considered in this work and the constraints that can be established from the lack of significant MCP signal. We will derive our constraints by performing a forward-folding binned-likelihood analysis, where the data and expectations are organized in equally sized bins of electron recoil energy. To construct our test statistic, we compute, for each one of the experiments considered, the number of signal events expected in a given bin of electron recoil energy. The number of signal events

in reconstructed electron recoil energy is given by

$$\mu_i^s = T\epsilon(E_i) \int_{E_i-E_b/2}^{E_i+E_b/2} dE_r \int_{-1}^{+1} d\cos\theta \frac{dN}{dE_r d\cos\theta}, \quad (4.3)$$

where the symbols inside the inner most integral are defined in eq. (4.2), T is the time exposure, E_b is the bin size, E_i is the bin center, and $\epsilon(E_i)$ is the average detection efficiency in a given bin.

To compute current experimental constraints, we construct a background-agnostic test statistic, comparing the observed data of recoil-electron-energies with those expected from an MCP with mass m_χ and mixing ε , ignoring any potential background contribution. We do not envision any additional triggers/cuts beyond those of the respective experiments searching for events with recoil electrons in the final state. This procedure will always result in a relatively conservative upper limit on the parameter space of MCP and cannot allow for a potential preferred region. We adopt this strategy for the current Super-Kamiokande and XENON1T experimental results for different reasons, which we will detail in their respective paragraphs to follow.

The background-agnostic test statistic is calculated using the bin-by-bin likelihood function [85]

$$\mathcal{L}_i = \begin{cases} P(d_i|\mu_i^s) & d_i < \mu_i^s \\ 1 & d_i \geq \mu_i^s \end{cases}, \quad (4.4)$$

where d_i and μ_i^s are the data and expected signal (given m_χ and ε) in bin i , respectively, and $P(d_i|\mu_i^s)$ is the Poissonian likelihood of observing d_i given expectation μ_i^s . This form of the likelihood is both background-agnostic and one-sided, guaranteeing the setting of a constraint instead of a preferred region. Our test statistic then is

$$\mathcal{TS} = -2 \log \left(\frac{\prod_i \mathcal{L}_i(m_\chi, \varepsilon)}{\prod_i \mathcal{L}_i(\varepsilon = 0)} \right), \quad (4.5)$$

where the denominator is the signal-free likelihood function and will return 1 when using the background-agnostic, one-sided likelihood of eq. (4.4).

When deriving a constraint, we will assume that Wilks' theorem holds and set limits assuming we have two degrees of freedom, m_χ and ε . In reality, the signal event distributions are all of the same shape (peaking at low electron recoil energy), where ε scales with the rate and m_χ determines which mesons can contribute to the MCP flux, also impacting the overall normalization. This relation implies that these two parameters could, in principle, be viewed as a single degree of freedom. In light of this, our use of two degrees of freedom, combined with the above background-agnostic, one-sided likelihood function approach, should be viewed as conservative.

Super-K. Our analysis uses the event selection in ref. [84], which was designed to search for diffuse supernova background neutrinos. This event selection was previously discussed in ref. [30], where they estimate event rates of MCPs produced in the atmosphere. We use the data from SK-I, SK-II, and SK-III and perform a joint likelihood analysis combining the three phases. For our analysis, we include the signal efficiencies provided in ref. [84]. The

background expectations estimated by the collaboration can be extracted from the same reference. We demonstrate the expected event rate from the collaboration’s background as well as the observed data during SK-II operation in figure 6 (left) as a function of electron recoil energy. However, it must be stated that this background expectation does not include the potential signal from diffuse supernova neutrino background (DSNB) events, the signature of interest in ref. [84]. If we perform a likelihood analysis⁵ comparing the observed data with the expected background (with no DSNB contribution) plus the MCP signal, we obtain a moderate ($\sim 90\%$ CL) preference for the existence of MCP; however, this signal is degenerate with the DSNB one. So, for this reason, we choose to adopt the conservative, background-agnostic approach discussed above to derive an upper limit on ε^2 as a function of m_χ at 90% CL.

XENON1T. Dark matter direct-detection experiments can measure electron recoil energies down to a few keV. We use the recent XENON1T experiment [86] result to search for signatures of millicharged particles produced in the Earth’s atmosphere, and refer the reader to references [87, 88] for similar studies with atmospheric Dark Matter. The background models presented in ref. [86] are notably insufficient to explain the observed rate of electron recoil events. However, one potential background to explain the excess is the presence of an unconstrained tritium beta-decay background. Because of this possibility, we choose to adopt the background-agnostic approach as we did for Super-Kamiokande. If we calculate the likelihood using the given B_0 background model of ref. [86], we observe a preferred region of parameter space at over 95% CL. This preferred region of parameter space is nearly excluded completely by other constraints at 2σ confidence and will be robustly tested by JUNO. See appendix C for more details, including the preferred region of parameter space that we obtain. When using eq. (4.2), we have assumed that all of the electrons in a detector are unbound — this is not the case in XENON1T, especially when comparing the observed recoil energy distribution with the nuclear binding energies of interest. In that context, our signal rates predicted in XENON1T will be overpredicted by a factor of several — see, for instance, refs. [31, 89] for further discussion of this effect. Given the scaling of our signal rate with ε^4 and the fact that XENON1T is not the most powerful experiment considered in this work, we disregard this effect in estimating our constraints.

JUNO. Given that JUNO offers a future search for single-scattering events, we calculate our test statistic assuming that the collected data in each bin is consistent with the background expectation and perform a comparison with the expected signal-plus-background distributions. See the right panel of figure 6 for these backgrounds from ref. [33] as well as a signal expectation. We then calculate our bin-by-bin likelihood using

$$\mathcal{L}_i = P(\mu_i^b | \mu_i^b + \mu_i^s), \quad (4.6)$$

⁵Our likelihood analysis for Super-Kamiokande includes only statistical uncertainties. We have performed a version of our analysis incorporating uncorrelated bin-by-bin normalization uncertainties of 5% (consistent with the systematic uncertainties discussed in ref. [84]) and find no realizable differences in the constraints we obtain.

where μ_i^b (μ_i^s) is the expected background (signal) in bin i . With this form of the binned likelihood, we then use the same expression for the test statistic $\mathcal{TS}$ in eq. (4.5) above and again, assume two degrees of freedom when projecting JUNO sensitivity at 90% CL. For our analysis, we assume an exposure of 170 kty, which corresponds to approximately ten years of data collection.

4.1.4 Single-scattering constraints and sensitivity projections

Figure 7 shows our new constraints at 90% CL, compared with existing upper limits from the SLAC-mQ experiment [72], ArgoNeuT [29], the milliQan Demonstrator [81], and colliders [78].⁶ Additional constraints, which we refrain from displaying because they are potentially subject to large background uncertainties and systematics associated with hand-scanning, can also be found in ref. [28]. Our Super-Kamiokande analysis, shown by the purple line and associated shaded region, yields results comparable results to those derived in ref. [30] for the bulk of parameter space. We note here again that given our mass-independent treatment of energy losses cf. eq. (3.1), our results could be mildly aggressive for $m_\chi < m_\mu$, where SLAC mQ dominates the parameter space. The small discrepancies between these two results arise from the following differences. The analysis of ref. [30] set constraints assuming an exposure of 22.5 kty at Super-Kamiokande, while here we have used the full 2853 day (or 7.81×22.5 kty) exposure of SK I-III and the statistical analysis discussed above to set this constraint. While the details of backgrounds are complicated, Hyper-Kamiokande should be able to improve on such constraints in the coming decades — we direct the reader to ref. [30] for more on millicharged particle searches at Hyper-Kamiokande and refs. [82, 90] for a detailed discussion of searches for the DSNB at Hyper-Kamiokande and associated challenges.

As shown in figure 7, the strongest current limit for $100 \text{ MeV} \lesssim m_\chi \lesssim 500 \text{ MeV}$ comes from this Super-Kamiokande analysis. However, a similar duration of a JUNO single-hit analysis (dashed orange line) will be able to improve on this, due to JUNO's ability to reach very low electron recoil energy. Although XENON1T (grey) can reach keV-scale electron recoils in this analysis, its small volume-exposure limits its capabilities relative to SK or JUNO.

4.2 Multiple scattering searches

Throughout section 4.1, we focused on scenarios where the atmospheric-produced MCP scatter inside a detector, providing enough energy to an electron to leave a distinct signature in the experiment. When discussing the signal of single-scattering events, we pointed out that the rate of signal events grows with smaller electron recoil energy E_r . However, in the case of both Super-Kamiokande and JUNO, the background do as well. Especially for JUNO, the background rate (dominated by radioactive emissions) became prohibitive for $E_r \lesssim 12 \text{ MeV}$, preventing a single-scattering analysis from reaching lower recoil energies.

One means for reducing (or even eliminating) these low-recoil-energy backgrounds is to require multiple scatterings within one small time window — the radioactive backgrounds

⁶We note here the different shape in this constraint relative to the one appearing in ref. [22], appearing to be from the same source. We believe the red region presented here represents the constraint from ref. [78].

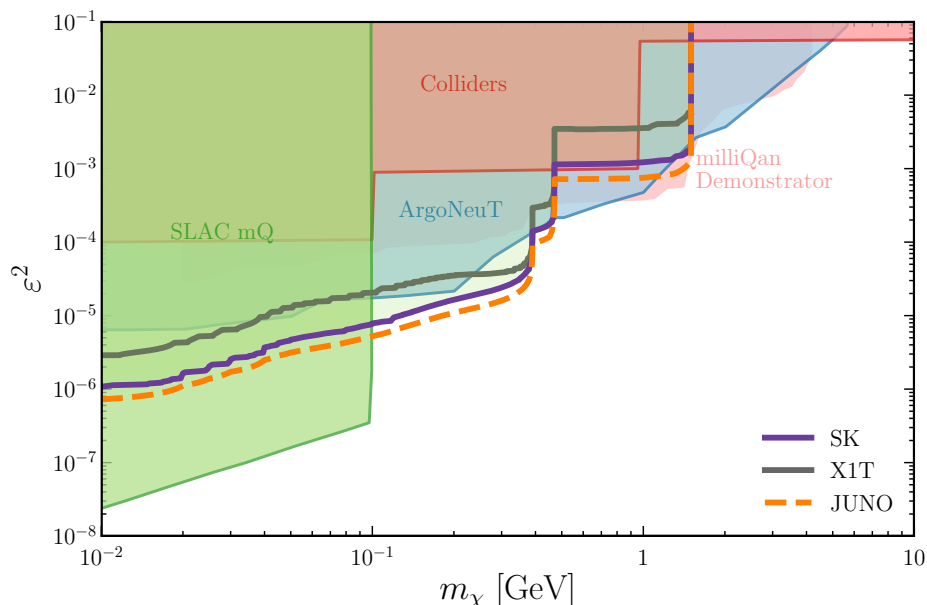


Figure 7. *Single-hit constraints on millicharged particles.* Constraints on millicharged particle parameter space from SLAC mQ [72], ArgoNeuT [29], the milliQan Demonstrator [81], and colliders [78], compared with our 90% CL single-hit constraints from Super-Kamiokande (purple), JUNO (projection, orange), and XENON1T (grey).

are each associated with some relatively large half-life and so the possibility of two such emissions in a $\sim 10^{-6}$ s window is very rare. Consequently, this type of analysis requires that a given MCP deposits energy at least twice as it traverses the detector, which will scale with even more powers of the millicharge ϵ than our single-scattering analysis. However, because this strategy allows us to search for even smaller recoil energies E_r where the scattering cross section grows, we will see that large event rates are still possible. This principle was exploited in ref. [21] for beam-produced millicharged particle searches in ArgoNeuT/DUNE.

In the remainder of this subsection, we will utilize this same approach for our atmospheric searches, focusing on the JUNO experiment. We will demonstrate that this approach can even outperform the single-scattering analysis presented in section 4.1 due to the low-energy capabilities of the JUNO detector’s liquid scintillator. A similar multiple-scatter strategy may be pursued at Hyper-Kamiokande, however, Cherenkov thresholds $\sim$ MeV limit the capabilities in such a detector.

4.2.1 Signal characteristics

We will determine the rate of multiple-scattering signal events in a given experimental setup by means of comparison with the single-scattering analyses discussed in section 4.1. For simplicity, we will consider total event rates instead of the event distribution as a function of electron recoil energy E_r . The important quantities to consider will be T_H , the minimum recoil energy for a hard-scattering event (i.e., those considered in the single-scattering analyses) and T_S , the minimum recoil energy capable of detecting events with multiple soft scatters in a small time window.

We perform the comparison between single- and multiple-scattering by computing the mean free path that a given MCP travels before either scattering in a “hard” manner (yielding an electron energetic enough to appear in single-scattering analyses) or in a “soft” manner (with electrons of too low an energy to be useful in single-hit searches, but energetic enough to be detected and used in a multiple-scattering search).

Assuming that a given MCP travels through a region with electron density n_e and length $L_{\text{det.}}$, we can approximate the probability that the MCP interacts using its hard-scattering mean free path $\lambda_H = 1/(n_e \sigma_H)$,

$$P_1 \approx 1 - e^{-\frac{L_{\text{det.}}}{\lambda_H}}, \quad (4.7)$$

where σ_H is the hard-scattering cross section, i.e., the cross section where we require the outgoing electron to be energetic enough to be distinct from radioactive backgrounds. In JUNO, this requirement is $E_r \gtrsim 12 \text{ MeV}$. Ref. [20] approximates σ_H as

$$\sigma_H \approx \frac{2\pi\alpha_{\text{EM}}^2 \varepsilon^2}{2m_e T_H} = 2.6 \times 10^{-25} \text{ cm}^2 \frac{\varepsilon^2}{(T_H/1 \text{ MeV})}. \quad (4.8)$$

The soft-scattering cross section and the corresponding mean free path between soft scatters are given by σ_S and λ_S , respectively. The cross section σ_S is identical to eq. (4.8) with T_S replacing T_H . In principle, T_S can be significantly lower than T_H , depending on the properties of the detector. For JUNO, we estimate that $T_S \approx 0.01 - 1 \text{ MeV}$ due to the 10^4 photons/MeV and $\mathcal{O}(10^3)$ photo-electrons/MeV of the detector’s liquid scintillator [34, 91]. The ratio of mean-free-paths λ_H/λ_S then is proportional to T_H/T_S .

In addition to the single-scattering probability P_1 in eq. (4.7), we can also consider the probability that an MCP scatters softly n times inside the detector. If we divide the total detector length into “segments” using $L_{\text{det.}} = N_{\text{seg.}} L_{\text{seg.}}$, this probability is

$$P_{(n)} = \frac{N_{\text{seg.}}!}{n!(N_{\text{seg.}} - n)!} \left(1 - e^{-\frac{L_{\text{seg.}}}{\lambda_S}}\right)^n \left(e^{-\frac{L_{\text{seg.}}}{\lambda_S}}\right)^{N_{\text{seg.}} - n}. \quad (4.9)$$

Assuming we can take the $N_{\text{seg.}} \rightarrow \infty$ (or $L_{\text{seg.}} \rightarrow 0$) limit for an unsegmented detector, this probability approaches

$$P_{(n)} = \frac{1}{n!} \left(\frac{L_{\text{det.}}}{\lambda_S}\right)^n e^{-\frac{L_{\text{det.}}}{\lambda_S}}. \quad (4.10)$$

The probability of two or more hits, $\sum_{n=2}^{\infty} P_{(n)}$ can be written as

$$P_{n \geq 2} = 1 - e^{-\frac{L_{\text{det.}}}{\lambda_S}} \left(1 + \frac{L_{\text{det.}}}{\lambda_S}\right). \quad (4.11)$$

The ratio of multiple soft-scatter events to the single hard-scatter events is proportional to $P_{n \geq 2}/P_1$. This quantity depends on ε^2 , the soft- and hard-scattering minimum recoil energies T_S and T_H , and the total detector length $L_{\text{det.}}$. For two choices of T_S and T_H , we present the ratio of these probabilities as a function of $L_{\text{det.}}$ and ε^2 in figure 8. In each panel, above the solid black lines, the multiple-hit rate exceeds the single-hit one, implying that searches for these multiple-scattering events can be at least as powerful as

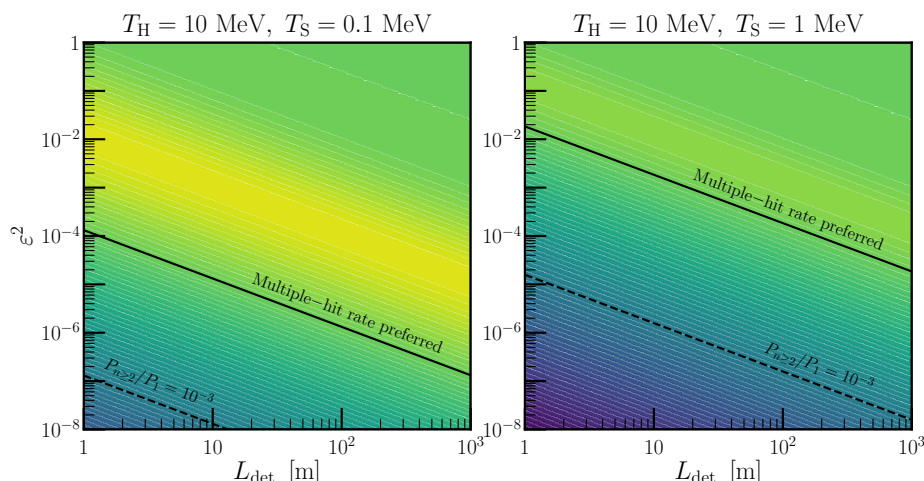


Figure 8. *Multiple-hit vs. single-hit search ratios.* Ratio of multiple soft-scatter events to single hard-scatter events as a function of total detector length $L_{\text{det.}}$ and ε^2 for two choices of T_H and T_S , as labelled. Above the solid black lines, the double-hit rate exceeds the single-hit one. Above the dashed lines, the double-hit rate is at most three orders of magnitude smaller than the single-hit one.

the single-scattering searches. If JUNO is capable of searching for 0.1 MeV soft-scatters, this implies that these multiple scattering events can be favorable for $\varepsilon^2 \approx 10^{-5}$ (assuming path lengths on the order of 10 m).

One final feature of the multiple-scattering signature that is not very useful in single-scattering searches is directionality. Because the single-scattering searches are seeking soft electrons, obtaining the direction of the incident millicharged particle is difficult. With a multiple-scattering analysis, we can obtain the angular distribution of these events, which should match the flux prediction, including attenuation through Earth, as discussed in section 3. This can be further used to statistically separate our signal from background events. As discussed cf. figure 4 and in ref. [21], deflection of MCP on the order of a few degrees for $\varepsilon^2 \approx 10^{-4}$ would smooth out the expected signal distribution with respect to $\cos\theta$ somewhat and this effect should be considered carefully in the case of a potential MCP discovery via multiple scattering.

4.2.2 Multiple-scattering backgrounds

All of the backgrounds we discussed for the single-scattering analysis can contribute as backgrounds for the multiple-scattering searches. However, they are only relevant if one or more of these stochastic backgrounds occurs within the time frame of an event for an experiment. The large background rates present at low recoil energy for JUNO (^8B , ^{10}C , ^{11}C , and ^{11}Be radioactive decays, specifically) will be suppressed by requiring multiple scatterings in this small window. For JUNO, the emission timescale of the liquid scintillator is at most ~ 200 ns [33] — if we can restrict the time window to be $\mathcal{O}(10^{-6}\text{s})$, then these backgrounds are reduced to being $\mathcal{O}(1)$ for the ten years of data collection we assume for JUNO.

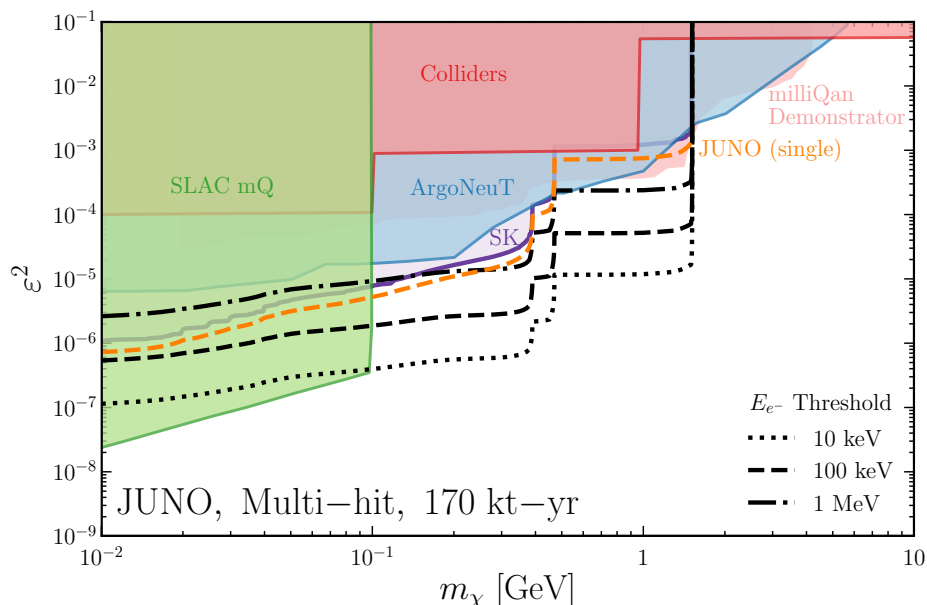


Figure 9. *Multiple-hit constraints on millicharged particles from JUNO.* Black: expected 90% CL sensitivity to millicharged parameter space using the multiple hit strategy of JUNO and a 170 kt-yr exposure. Three different assumptions about the minimum observable electron recoil energy are assumed: 10 keV (dotted), 100 keV (dashed), and 1 MeV (dot-dashed). Existing constraints (including our Super-Kamiokande analysis in purple) are shown as filled in regions, and the single-hit analysis of JUNO from figure 7 is shown as a dashed orange line.

4.2.3 Projected sensitivity with multiple scattering

Combining the above information, we project the sensitivity using this multiple-scattering search in JUNO here in figure 9. In order to determine the “detector length” of JUNO that we discussed above, we simply average the path length of incoming MCP over the 35-m diameter spherical vessel. The average path length through a sphere is $2D/3 = 23.3$ m for JUNO.

We assume that the backgrounds in JUNO are small enough that 10 signal events are statistically significant in the 170 kt-yr exposure — we do not use any spectral information about the energy of the electron events, as we expect that resolution at such sub-MeV energies will be difficult. We take three different assumptions about the minimum threshold energy of electrons that JUNO can detect — 1 MeV (dot-dashed), 100 keV (dashed), and 10 keV (dotted). If this 10-keV threshold is attainable (which could be possible given the 10^4 photons/MeV of energy deposited in the liquid scintillator), we see that this multiple-hit strategy will far exceed single-scattering searches for atmospheric MCP. Even with MeV thresholds, this is a complementary approach, especially at large m_χ . The realistic 100-keV threshold seems particularly promising as a target for JUNO. Finally, we note that additional care in the mass-dependent energy loss is required to properly derive constraints for $m_\chi < m_\mu$.

5 Millicharged particles at neutrino telescopes

The search for millicharged particles from natural sources is a statistically-limited problem, whose optimal detector would be a low-threshold, small-background, large-mass underground

device. In this work, we have so far studied kiloton-mass-scale low-threshold and small-background detectors, such as XENON1T, and tens of kilotons mass low-threshold detectors, such as JUNO, in this section we will study the sensitivity of megaton and gigaton detectors, known as neutrino telescopes. These detectors use natural transparent media to detect Cherenkov light produced by neutrino interactions and are designed to detect faint high-energy astrophysical neutrino sources. However, their large detector volume allows them to also have unique capabilities to observe low-energy neutrinos produced in supernovae.

Since MCPs are long-lived and have small energy losses their signature in neutrino telescopes will be a long, faint track. The main backgrounds for this search are coincident or dim atmospheric muons that can accidentally mimic the signature. Searches for long-lived faint-tracks have been performed by the IceCube collaboration by looking for an isotropic flux of fractional charged particles with charges motivated by the quark charges [92, 93]. This analysis yielded a sensitivity approximately ten times stronger than previous constraints from Kamiokande and MACRO, however the assumption of an isotropic flux of fractional charge particles is not consistent with the propagation of these particles through the Earth, as we discussed in section 3. Despite this, refs. [92, 93] motivate further study of the sensitivity to these signals.

The yield of Cherenkov photons produced per energy per unit path length of the distance travelled by a particle of charge ϵe is given by the Frank-Tamm equation [17]

$$\frac{d^2 N}{dE dX} = \frac{\alpha \epsilon^2}{\hbar c} \sin^2 \theta_c \approx 160 \epsilon^2 \left(\frac{\sin^2 \theta_c}{0.43} \right) \text{eV}^{-1} \text{cm}^{-1}, \quad (5.1)$$

where θ_c is the Cherenkov angle in the medium, which we have normalized to the Cherenkov angle in ice, $\theta_c \approx 41^\circ$ [94]. The relevant Cherenkov photon wavelength for IceCube is around 400 nm, which results from the convolution of the IceCube PMT quantum efficiency and the glass-housing transmission probability [95, 96]. At this wavelength the IceCube digital optical module (DOM) acceptance is approximately 0.15, and decreases by approximately half by reducing the wavelength to 350 nm or increasing it to 550 nm. Considering this, the yield of relevant Cherenkov photons is approximately $50 \epsilon^2 \text{cm}^{-1}$. This implies that an IceCube-through-going MCP will produce, on average, $5 \times 10^6 \times \epsilon^2$ relevant Cherenkov photons. Next we need to estimate how many of these Cherenkov photons could reach an IceCube module, this implies taking into account the absorption of the Antarctic ice [94] and the single-photo-electron efficiency [97]. The number of photons from a single-point light source, which would be a given MCP energy loss, is reduced by a factor of 10^{-3} at a 10 m distance and by 10^{-6} at a 100 m [98]. Thus, if an MCP passes at a 10 m distance from an IceCube DOM the Cherenkov photon yield will be approximately $10 \epsilon^2$ photons. This estimation implies that the relevant MCP parameter space that can produce enough Cherenkov photons in IceCube is above approximately 5 GeV, cf. figure 7.

In the above discussion, we have only taken into account the Cherenkov light, but there are other sources of light production from charged particles, such as ionization. To study the sensitivity of neutrino telescopes to millicharged particles, while taking into account these other losses, we have developed a dedicated Monte Carlo to estimate the trigger rate in

detectors such as the IceCube Neutrino Observatory in the South Pole. This Monte Carlo can be found in our GitHub repository at <https://github.com/Harvard-Neutrino/HeavenlyMCP>.

In order to estimate the trigger rate we need to know the precise locations of the energy depositions of the MCPs and their distance to the detection units. Our Monte Carlo consists of the following stages:

1. We produce MCPs according to a power-law energy distribution and spread them uniformly across the surface of the Earth.
2. We propagate the MCPs from the Earth surface to a cylinder that contains the detector instrumented volume. To perform this we use **PROPOSAL**, which is a Monte Carlo package that simulates the energy losses of charged leptons in different media. The energy losses implemented in **PROPOSAL** are equivalent to the ones discussed in section 3, however, unlike our previous discussion which focused on the mean energy loss, **PROPOSAL** provides a detailed simulation of continuous energy losses and stochastic ones. Thus, in this step, for each MCP particle in our Monte Carlo we obtain the location, type, and amount of energy loss along the particle trajectory.
3. We convert the energy losses produced in the detector vicinity to the number of Cherenkov photons produced in ice. In order to do this, we use the publicly available PPC photon propagation code, which uses the parameterizations given in [77] to convert each type of loss (ionization, bremsstrahlung, photo-hadronic, and pair-production) into an ensemble of Cherenkov photons.
4. The in-ice photons are then propagated by PPC through the detector instrumented volume. Within this volume, we specify the detection units in PPC by providing the coordinate of each sensor. With this information PPC returns the number of detection units hit by photons.

Using this Monte Carlo approach one can estimate the trigger rate in a given detector. However, this is beyond the scope of this article and instead we provide this Monte Carlo and MCP fluxes computed in section 2 in order for experiments to estimate the MCP yield in their specific setup.

6 Discussion and conclusions

Whether particle charges are quantized and whether any charged particles exist beyond the Standard Model are two questions that have been asked for generations. Searches for new particles with fractional charges work to address both of these questions, and significant progress has been made in these searches in recent years, particularly in the MeV to GeV mass range.

Focus in this mass range has been divided between two general categories — searches for millicharged particles produced by collider and fixed-target experiments, and searches for millicharged particles naturally produced in the atmosphere. We have focused on the latter approach. In this work, we have revisited current constraints from the Super-Kamiokande

neutrino experiment, qualitatively confirming the existing literature on searches of this type. We have also analyzed existing data from the XENON1T dark matter direct-detection experiment and projected future capabilities of the JUNO reactor neutrino oscillation experiment in this parameter space, demonstrating paths for improvement in the next decade.

Going beyond this, we have combined a number of millicharged particle search strategies by proposing the search for multiple-scattering events in a liquid scintillator detector (specifically JUNO), allowing for sensitivity to significantly smaller millicharges than conventional, single-scattering searches. This is because the multiple-scattering searches allow for analyses to probe even smaller energies where backgrounds dominate the conventional search.

We have focused predominantly on searches for these particles in tens-of-kiloton-scale detectors. However, even larger detectors, such as the IceCube Neutrino Observatory (and forthcoming IceCube Upgrade) can offer an interesting, complementary means of searching for millicharged particles. We demonstrated that detectors with longer path lengths (of traversing millicharged particles) are well-suited for multiple-scattering searches. IceCube is one such detector, and it can potentially perform a search for these “faint track” signatures in the coming years. We have developed a Monte Carlo package to simulate the propagation and energy deposition of millicharged particles through a detector such as IceCube.

As we look forward to the next decade of searches for new, fractionally charged particles, it is imperative that we combine as many search strategies as possible. This maximizes the chances of discovery, and, in the hopeful event of one discovery, a combined approach is our best way to interpret and understand such a momentous result. Atmospheric searches, particularly those looking for multiple-scattering events, offer a powerful means to search for these particles, complementary to current and upcoming collider and fixed-target based searches.

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A Details of millicharged particle production

We consider production of millicharged particles via rare decays of neutral mesons π^0 , η , ρ , ω , ϕ , and J/Ψ . The first two (pseudoscalar mesons) can result in three-body decays $\mathbf{m} \rightarrow \gamma\chi\bar{\chi}$, analogous to Dalitz decays. The latter four (vector mesons) can decay via an off-shell photon $\mathbf{m} \rightarrow \chi\bar{\chi}$. Below we give the branching ratios of these processes as well as the energy spectrum of the daughter χ particles.

The ω meson has an additional decay channel that has a relatively large width — $\omega \rightarrow \pi^0 e^+ e^-$, with a branching ratio of 7.7×10^{-4} compared to $\text{Br}(\omega \rightarrow e^+ e^-) = 7.36 \times 10^{-5}$ [17]. The corresponding width of $\omega \rightarrow \pi^0 \chi \bar{\chi}$ is more complicated than the two- and three-body decay widths we present below. In principle, there should be an additional χ flux nearly an order of magnitude larger than our estimates in the mass range $m_\eta/2 < m_\chi < 1/2(m_\omega - m_{\pi^0})$. For simplicity, and due to the narrow mass range that this impacts, we choose to neglect this decay channel in our calculations.

Two-body decays. The branching ratio of the two-body decays for vector mesons into millicharged particle pairs can be expressed as

$$\text{Br}(\mathbf{m} \rightarrow \chi \bar{\chi}) = 2\varepsilon^2 \text{Br}(\mathbf{m} \rightarrow e^+ e^-) I^{(2)}\left(\frac{m_\chi^2}{m_M^2}, \frac{m_e^2}{m_M^2}\right), \quad (\text{A.1})$$

where $\text{Br}(\mathbf{m} \rightarrow e^+ e^-)$ is the experimentally-measured branching ratio of the meson into electron/positron pairs and $I^{(2)}(x, y)$ is a dimensionless quantity relating these two processes,

$$I^{(2)}(x, y) = \frac{(1 + 2x)\sqrt{1 - 4x}}{(1 + 2y)\sqrt{1 - 4y}}. \quad (\text{A.2})$$

The energy distribution of the millicharged particles in the parent meson rest frame is flat. The lab-frame distribution can be obtained by transforming between frames using the boost of the parent meson. The allowed energy of the millicharged particle in the lab frame E_χ can be determined by requiring

$$\frac{2E_\chi}{1 + \sqrt{\lambda\left(1, \frac{m_\chi^2}{m_m^2}, \frac{m_\chi^2}{m_m^2}\right)}} \leq E_m \leq \frac{2E_\chi}{1 - \sqrt{\lambda\left(1, \frac{m_\chi^2}{m_m^2}, \frac{m_\chi^2}{m_m^2}\right)}}, \quad (\text{A.3})$$

where $\lambda(a, b, c)$ is the Källén function [99]

$$\lambda(a, b, c) = a^2 + b^2 + c^2 + 2ab + 2ac + 2bc. \quad (\text{A.4})$$

Three-body decays. For the pseudoscalar meson (π^0 and η) decays, the process of interest is $\mathbf{m} \rightarrow \gamma \chi \bar{\chi}$. The branching ratio for this decay can be related to the branching ratio of $\mathbf{m} \rightarrow \gamma \gamma$ using

$$\text{Br}(\mathbf{m} \rightarrow \gamma \chi \bar{\chi}) = 2\varepsilon^2 \alpha_{\text{EM}} \text{Br}(\mathbf{m} \rightarrow \gamma \gamma) I^{(3)}\left(\frac{m_\chi^2}{m_m^2}\right), \quad (\text{A.5})$$

where $I^{(3)}(x)$ is, similar to $I^{(2)}(x, y)$, a dimensionless function [24],

$$I^{(3)}(x) = -\frac{1}{3\pi}(1 - 4x)^{3/2}(7 + 2x) + 2\left(1 + 2x^2(4x - 9)\right) \tanh^{-1}\left(\sqrt{1 - 4x}\right). \quad (\text{A.6})$$

To obtain the lab-frame distribution of the millicharged particle energy, we use the quantity $z \equiv E_\chi/\gamma_m$, where γ_m is the meson boost. This distribution can be expressed as [59]

$$\frac{1}{\Gamma} \frac{d\Gamma}{dz} = \frac{(z - m_M)}{3z^3 Y_m(m_m, m_\chi)} F(z, m_m, m_\chi), \quad (\text{A.7})$$

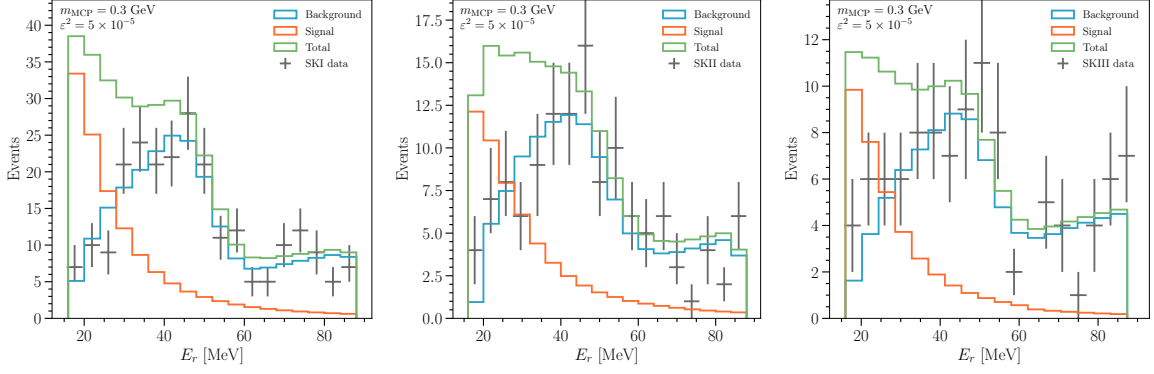


Figure 10. Distributions for our millicharged particle search for the three stages of Super-Kamiokande. Data is shown as crosses and expected background is shown as a blue histogram. An example millicharge particle distribution is shown, for the same parameters, as an orange histogram.

with

$$Y_m(m_m, m_\chi) = m_m^7 - 32m_m^5 m_\chi^2 - 348m_m^3 m_\chi^4 \log(2m_\chi/m_m) + 512m_m m_\chi^6 - 256 \frac{m_\chi^8}{m_m}, \quad (\text{A.8})$$

and

$$\begin{aligned} F(z, m_m, m_\chi) = & z(320m_m m_\chi^6 + 144m_m^3 m_\chi^4) - z^2(m_\chi^6 + 432m_m^2 m_\chi^4) \\ & + z^3(108m_m^3 m_\chi^2 - 5m_m^5) - z^4(m_m^2 m_\chi^2 + m_m^4) \\ & + 4m_m^3 z^5 - 256m_\chi^6 m_m^2. \end{aligned} \quad (\text{A.9})$$

For a three-body decay we can determine the allowed range of E_χ using

$$\frac{m_m E_\chi}{E_{\max} + \sqrt{E_{\max}^2 - m_\chi^2}} \leq E_m \leq \frac{m_m E_\chi}{E_{\max} - \sqrt{E_{\max}^2 - m_\chi^2}}, \quad (\text{A.10})$$

with $E_{\max} \equiv \frac{4m_\chi^2 + m_m^2}{2m_m}$.

B Additional details of the Super-Kamiokande analysis

When discussing expected signal and background event distributions in Super-Kamiokande (cf. figure 6 left), we showed the expected distributions with Super-Kamiokande II for simplicity. For completeness in figure 10, we provide the analogous distributions for all three stages of Super-Kamiokande data collection that enter our analysis. Each panel displays the expected background events (blue), signal assuming $m_\chi = 300$ MeV and $\epsilon^2 = 5 \times 10^{-5}$ (orange), and their sum (green), compared against the observed data (black crosses).

In practice, we do not use the expected background distributions (blue) in our analysis. This is because they do not include the possibility of any diffuse supernova neutrino background events, which would contribute at low recoil energy. For this reason, we adopt the background-agnostic, one-sided likelihood approach discussed in section 4.1.3.

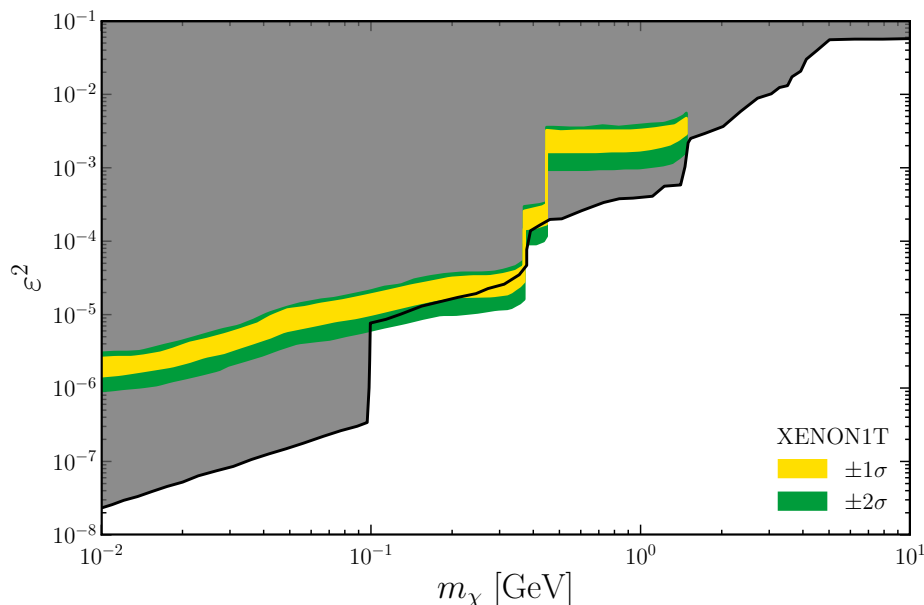


Figure 11. Preferred parameter space from XENON1T. Yellow/green bands depict the $\pm 1\sigma/\pm 2\sigma$ preferred parameter space of our XENON1T analysis when we assume that the B_0 background of ref. [86] accurately represents all backgrounds. All existing constraints are shown in grey.

C XENON1T preferred parameter space

In section 4.1 we discussed constraints on millicharged particle parameter space as a function of m_χ and ε^2 coming from single-hit searches, including Super-Kamiokande, XENON1T, and a future search in JUNO. When considering current data from Super-Kamiokande and XENON1T, we took a background-agnostic approach with a one-sided likelihood function to derive a conservative upper limit on ε^2 as a function of m_χ . For XENON1T, this was done in part due to the much-discussed potential tritium background, absent in the nominal B_0 background model of ref. [86].

In this appendix, we perform an alternate test of the XENON1T data where we assume that the B_0 model is robust and calculate the likelihood comparing μ_i^s (the expected number of signal events in bin i) plus $B_{0,i}$ (the expected background in bin i) with the observed data d_i . This process, due to the excess electron-like events in the lowest energy bins, will yield a preferred region of millicharged parameter space instead of a constraint.

The potential signals proposed in ref. [86] include neutrino magnetic moments and solar axions, which improve the fit to data over the background-only explanation by $3.2 - 3.4\sigma$. We find that the millicharged particle signature improves our test statistic over the background-only hypothesis by 10.43 units. When we (conservatively) consider two degrees of freedom, this implies a preference of 2.8σ over the null hypothesis — if we had considered the highly-correlated (m_χ, ε^2) to account for a single degree of freedom, this would imply a preference of 3.2σ , as preferable as the hypotheses presented by the XENON1T Collaboration.

Figure 11 presents our 1σ (yellow) and 2σ (green) preferred region of parameter space from XENON1T, compared against all other existing limits (including our Super-

Kamiokande analysis) in grey. We see that the bulk of the $\pm 2\sigma$ preferred parameter space is already excluded by one or more other constraint, however, some regions survive near $m_\chi \approx 350$ MeV. These regions will be tested by JUNO's single-hit analysis and thoroughly explored by JUNO's double-hit analysis.

We note here, as in the main text, that our XENON1T analysis and signal-event prediction rate from eq. (4.2) do not include the fact that the electrons in the XENON1T detector are bound [31, 89]. This is relevant because their binding energies are non-negligible compared to the recoil energy observed in the detector, and this implies that our signal rate predictions are optimistic. For this reason, in addition to the tritium background discussion above, we caution the reader from inferring that the MCP solution to the XENON1T excess is a likely explanation.

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Hadrophilic light dark matter from the atmosphere

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ABSTRACT

Light sub-GeV dark matter (DM) constitutes an underexplored target, beyond the optimized sensitivity of typical direct DM detection experiments. We comprehensively investigate hadrophilic light DM produced from cosmic-ray collisions with the atmosphere. The resulting relativistic DM, originating from meson decays, can be efficiently observed in variety of experiments, such as XENON1T. We include for the first time decays of η , η' and K^+ mesons, leading to improved limits for DM masses above few hundred MeV. We incorporate an exact treatment of the DM attenuation in Earth and demonstrate that nuclear form factor effects can significantly impact the resulting testable DM parameter space. Further, we establish projections for upcoming experiments, such as DARWIN, over a wide range of DM masses below the GeV scale.

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1. Introduction

Extensive astrophysical observations point towards the existence of dark matter (DM), currently constituting $\sim 85\%$ of matter in the Universe (see [1,2] for a review). Despite decades of searching for its possible non-gravitational interactions, the nature of the DM remains elusive. One extensively explored paradigm is that of Weakly Interacting Massive Particles (WIMPs), where the DM has a typical mass in the range of 1 GeV–100 TeV, and often appears in theories aiming to address the hierarchy problem. A less studied scenario, which could also appear in variety of well-motivated models (see e.g. [3–7]), is sub-GeV (light) DM. This later case is the focus of this article.

A cornerstone of DM searches are direct DM detection experiments, which looks for energy depositions on a target material associated with interactions within the experiment of traversing DM from the Milky Way halo. Conventional ton-scale direct DM detection experiments with keV-level energy thresholds (see e.g. [9,10]) face significant challenges to search for light DM, since it is not expected to produce sufficient nuclear recoil when DM interacts with the target material.

Observational signatures of DM could be drastically distinct from nominal expectations if the DM speed distribution contains a significant relativistic component. Such component can be naturally realized when DM is produced in cosmic-rays collisions with the atmosphere. The “atmospheric beamdump” has been historically exploited as an essential instrument for exploring the properties of neutrinos, contributing to the discovery of neutrino oscillations [11]. Recent studies have highlighted that the atmospheric beamdump is a unique tool to explore new physics beyond the Standard Model (SM), including long-lived particles [12], heavy neutral leptons [13], millicharge particles [14–16], particle DM [8,17], supersymmetric particles [18], and magnetic monopoles [19]. This natural beam dump produces a steady flux that can be exploited by terrestrial experiments. Other possible sources of a relativistic DM component include boosted DM (see e.g. [20]), DM accelerated in astrophysical sources [21–23], the result of cosmic-ray DM scattering [24–29] or evaporation of primordial black holes [30]. These DM components are distinct from conventional halo DM analyses (e.g. [31,32]).

In this article, we comprehensively analyze light hadrophilic DM produced in the atmosphere by cosmic-ray collisions. In particular, we focus on hadrophilic dark sectors, which have been discussed within variety of scenarios (see e.g. [33–36]). We improve upon earlier work [8] of hadrophilic DM producing in an atmospheric beamdump in many significant aspects which include: considering a more realistic model of the atmospheric flux, including of additional mesons whose decay produces DM, and perform

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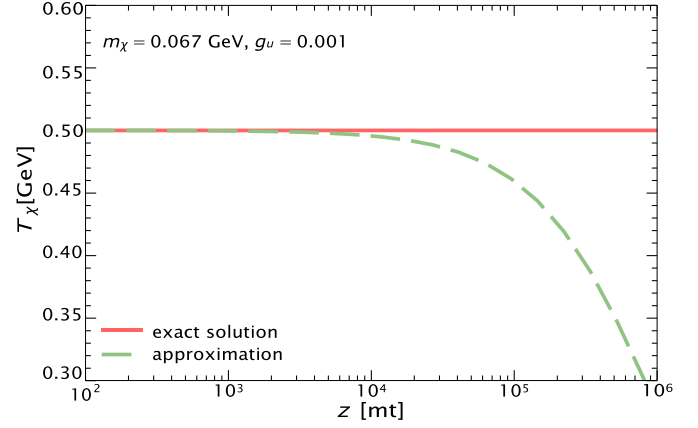
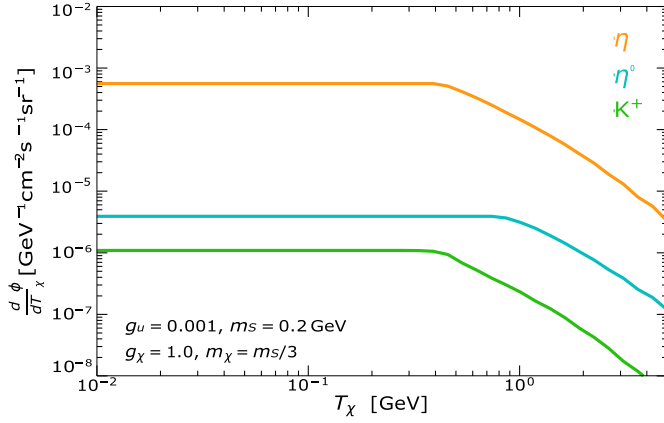


Fig. 1. Left: DM flux spectrum at surface of the Earth in terms of DM energy T_χ , produced from atmospheric decays of mesons. **Right:** DM energy loss due to attenuation as it propagates through the Earth, as a function of the depth. The exact solution of Eq. (9) (solid red curve) and the expected loss in approximate treatment of Ref. [8] (dashed green curve), which does not incorporate the nuclear form factor, are displayed.

a proper treatment of the attenuation of DM through the Earth. We also present projections of this scenario for upcoming and proposed experiments.

2. Hadrophilic light dark matter

Collisions of the almost isotropic cosmic-ray flux, primarily composed of protons, with the atmosphere results in copious production of Standard Model particles from GeV to well over TeV energies [37]. This is directly analogous to a fixed target beam dump experiment. The atmospheric beam dump can produce particles associated with physics beyond the SM.

To explore sub-GeV DM, we consider a minimal realization within a framework of flavor-specific, hadrophilic, scalar mediator [33,34]. The effective Lagrangian includes

$$\mathcal{L} \supset i\bar{\chi}(\not{D} - m_\chi)\chi + \frac{1}{2}\partial_\mu S\partial^\mu S - \frac{1}{2}m_S^2 S^2 - (g_\chi S\bar{\chi}_L\chi_R + g_u S\bar{u}_L u_R + h.c.), \quad (1)$$

where S is a singlet scalar mediator with mass m_S , χ is a singlet Dirac fermion DM with mass m_χ stabilized by a Z_2 symmetry, and g_u and g_χ are effective couplings. Here, the scalar mediates the interactions between the dark sector and the visible sector, in our work, the up quarks.

In the medium where the DM propagated and in the experiments instrumented volume, DM will scatter against a nucleus with differential cross-section given by

$$\frac{d\sigma_{\chi N}}{dT_r} = \frac{g_\chi^2 (Zy_{spp} + (A - Z)y_{snn})^2}{8\pi} \frac{(T_r + 2m_N)}{(T_\chi^2 + 2m_\chi T_\chi)} \times \frac{(2m_\chi^2 + m_N T_r)}{(2m_N T_r + m_S^2)^2} F_H^2(2m_N T_r), \quad (2)$$

where A , T_r , and m_N are the nucleon number, recoil energy, and nucleus mass, respectively. The term F_H is the nuclear Helm form factor [38] (see also e.g. [39])

$$F_H(q^2) = \frac{3j_1(qR_1)}{qR_1} e^{-\frac{1}{2}q^2 s^2}, \quad (3)$$

where $q = \sqrt{2m_N T_r}$ is the momentum transfer, $j_1(x)$ is a spherical Bessel function of the first kind, $R_1 = \sqrt{R_A^2 - 5s^2}$ with $R_A \simeq 1.2A^{1/3}$ fm and $s \simeq 1$ fm. The g_χ is the scalar-DM coupling, which we fix at 1.0, while the effective scalar-nucleon couplings, y_{spp} , y_{snn} are given by [34]

$$y_{spp} = 0.014 g_u \frac{m_p}{m_u},$$

$$y_{snn} = 0.012 g_u \frac{m_n}{m_u}, \quad (4)$$

with m_p , m_n , and m_u the masses of the proton, neutron, and up-quark.

The nuclear interaction cross-section can be related to the conventional spin-independent cross-section, σ^{SI} , by

$$\sigma_{\chi N} = \sigma^{\text{SI}} A^2 \left(\frac{m_N (m_\chi + m_p)}{m_p (m_\chi + m_N)} \right)^2. \quad (5)$$

3. Atmospheric dark matter production

The scalar mediator S can be abundantly produced in meson decays in cosmic-ray air showers. The production is modeled by the cascade equation (e.g. [40]), which, omitting the interaction terms, contains the following source term for DM production

$$\frac{d\Phi_S}{dE_S d\Omega dX} = \sum_M \int dE_M \left(\frac{1}{\rho(X)\lambda_M(E_M)} \times \frac{d\Phi_M}{dE_M d\Omega}(E_M, \cos\theta, X) \times \frac{dn}{dE_S}(E_M, E_S) \right), \quad (6)$$

where the sum M runs over all possible mesons that can decay to S , ρ is the density of the atmosphere, X is the column depth (i.e. the depth in g/cm² measured from the top of the atmosphere to production point), and $\lambda_M \equiv \gamma_M \beta_M c \tau_M$ is the laboratory decay length of the meson that includes the boost factor $\gamma_M \beta_M$ and its proper lifetime τ_M . Here, the differential production rate of mesons in the shower per unit of solid angle is given by $(d\Phi_M/dE_M d\Omega)$.

In Eq. (6), the number of scalar mediators with energies between E_S and $E_S + dE_S$ produced in the decay of the meson M is given by (dn/dE_S) . The explicit expression for the last quantity depends on whether the scalar is produced in two-body or three-body decays. We focus on the simplest case involving only two-body decays of mesons. In this case, the energy distribution in the laboratory frame is given by

$$\frac{dn}{dE_S} = \frac{\text{Br}(M \rightarrow S + P)}{p_M \sqrt{\mathcal{K} \left(1, \frac{m_S^2}{m_M^2}, \frac{m_P^2}{m_M^2} \right)}}, \quad (7)$$

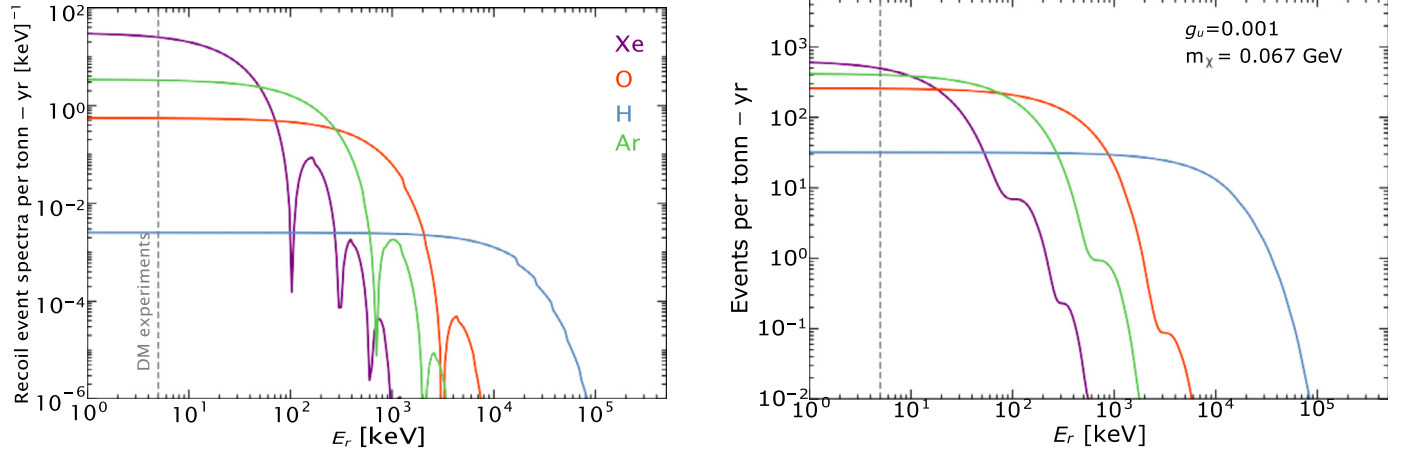


Fig. 2. **Left:** Differential event spectra for different target materials as a function of nuclear recoil energy E_r , including xenon (Xe, purple), oxygen (O, orange) and hydrogen (H, blue). The total exposure is one ton-year. **Right:** Total integrated number of recoil events above given experimental energy threshold E_r , for xenon (purple), oxygen (orange) and hydrogen (blue). The considered model values are $g_u = 10^{-3}$ and $m_\chi = 0.067$ GeV.

where $Br(M \rightarrow S + P)$ is the meson branching fraction to the scalar and an additional particle P , K is the K  llen function [41], and p_M is the meson momentum.

For scalar mediators sufficiently strongly coupled to the DM, they will promptly decay to a pair of $\chi\bar{\chi}$ in the particle shower. The resulting DM production rate can be obtained using mediator production of Eq. (6) via

$$\frac{d\Phi_\chi}{dE_\chi d\Omega dX} = \int dE_S \frac{d\Phi_S}{dE_S d\Omega dX} \frac{dn}{dE_\chi}(E_S, E_\chi). \quad (8)$$

Eq. (8) implies that the DM is produced at the same point (height) as the scalar mediator, while accounting for the energy distribution of the DM in the decay. If the coupling between the scalar and the dark matter is not strong enough, the mediator could be made a long-lived particle and the production of DM would be suppressed. Integrating Eq. (8) over the column depth yields the expected DM flux at the surface of the Earth per unit of solid angle and per energy bin.

In this work, we include for the first time the production of scalar mediator from the decays of η , η' as well as charged kaons in the atmosphere, via $\eta \rightarrow \pi^0 S$, $\eta' \rightarrow \pi^0 S$ as well as $K^+ \rightarrow \pi^+ S$ channels.¹ We review the branching ratios for these processes in Appendix A. We numerically solve for the production rate of the mesons in the atmospheric cascade using the Matrix Cascade Equation (MCEq) software package [44,45], which includes several models for the cosmic-ray spectrum, hadronic interactions, and atmospheric density profiles. For our base results we employed the SYBILL-2.3 hadronic interaction model [46], the Hillas-Gaisser cosmic-ray model *H3a* [47], and the NRLMSISE-00 atmospheric model [48] to obtain the production rate of mesons, which have been extracted following the procedure of Ref. [12,13].

The resulting DM flux at the surface of the Earth is shown in Fig. 1. We show each of the progenitor mesons η , η' and K^+ , considering $g_u = 10^{-3}$, $g_\chi = 1$, $m_S = 0.2$ GeV, and $m_\chi = m_S/3$. With a mass of $m_{\eta'} = 957.8$ MeV, η' allows for additional kinematic regimes compared to η and K^+ , since these have smaller masses; namely $m_\eta = 547.9$ MeV and $m_{K^+} = 493.7$ MeV, respectively.

¹ This is in contrast with previous work [8], which only considered η decays. Different meson decay channels can themselves be used to investigate new physics (e.g. [42,43]).

4. Dark matter flux attenuation

In analogy to SM particles such as muons, the flux of DM with sizable interaction strength is attenuated as it traverses the medium from the top of the atmosphere towards the experiment. At large cross-sections the medium becomes ‘‘optically thick,’’ causing direct DM detection experiments to dramatically lose their sensitivity in this part of the DM parameter space (e.g. [49,50]).

The DM energy loss is described by

$$\frac{dT_\chi}{dz} = - \sum_N n_N \int_0^{T_r^{\max}} T_r \frac{d\sigma_{\chi N}}{dT_r} dT_r, \quad (9)$$

where T_r is the energy lost by a DM particle in a collision with nucleus N , and

$$T_r^{\max} = \frac{T_\chi^2 + 2m_\chi T_\chi}{T_\chi + (m_\chi + m_N)^2/2m_N}. \quad (10)$$

To account for DM flux attenuation, we numerically solve Eq. (9).

Our analysis takes into account the relevant effects of the nuclear form factor of Eq. (3), which in contrast have been neglected in previous work employing analytic approximation of Eq. (9) to treat attenuation [8]. The effects of the nuclear form factor in attenuation can significantly impact which DM parameter space regimes are direct DM detection experiments are sensitive to, as has been recently highlighted in Ref. [51].

In Fig. 1 we display the DM flux attenuation using the exact numerical treatment of Eq. (9) as well as analytic approximation given in Ref. [8]. We highlight that the approximate treatment did not include the effects arising from the nuclear form factor, which suppress large momentum transfer regimes resulting in an over-estimation of the DM energy losses for underground experiments lying at a depth below 1 km of rock overburden. Having found $T_\chi(z)$ with the procedure outlined above, we can determine the resulting atmospheric DM flux as a function of depth $d\Phi_\chi/dT(z)$.

5. Experimental constraints

The differential event rate for a detector with N_t nucleons is given by

$$\frac{dN}{dT_r} = N_t \int_{E_\chi^{\min}} dE_\chi \epsilon(T_r) \frac{d\sigma_{\chi N}}{dT_r} \frac{d\Phi_\chi}{dT_\chi}, \quad (11)$$

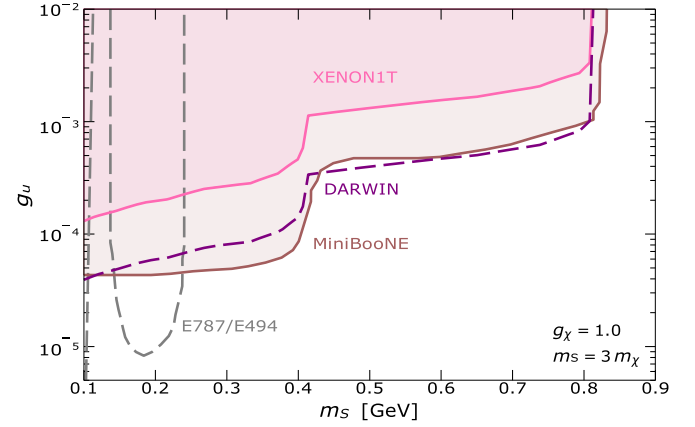
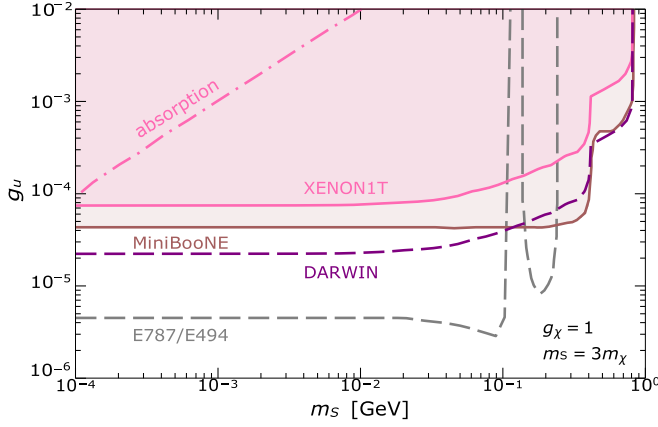


Fig. 3. Left: Limits from atmospheric DM flux from meson decays (red), for a hadrophilic scalar mediator of mass m_S on up-quark coupling g_u , considering $g_\chi = 1$ and $m_S = 3m_\chi$. **Right:** Same, with focus on parameter space for m_S above 0.1 GeV.

where $\epsilon(T_r)$ is the detection efficiency taken from reference [52]. In Fig. 2 we show the event spectra as a function of the nuclear recoil energy T_r for different target materials, including xenon (Xe), oxygen (O) and hydrogen (H). The bumps in the spectrum follow the shape of the nuclear Helm factor in the interaction cross-section. We also display in Fig. 2 the total integrated number of nuclear recoil events above experimental threshold T_r , for different materials and considering characteristic model parameter values of $g_u = 10^{-3}$ and $m_\chi = 0.067$ GeV. These results highlight the importance of low O(1) keV experimental thresholds available for large DM experiments, in accord with the literature results for other physics analyses and complementarity studies (e.g. [53–56]). We note that hadrophilic DM could also induce interactions in large neutrino experiments. However, such interactions require significantly higher thresholds and the resulting event rates are expected to be suppressed even taking into account the larger available fiducial mass.

To set limits on the model up-quark coupling g_u from XENON1T, we integrate Eq. (11) in the recoil energy window with a threshold of 5 keV and a maximum recoil value of 40.9 keV. Following Ref. [52], we set a limit by requiring to observe more than 3.56 events for an exposure of 278.8 live-days. The limit in the plane $g_u m_S$ is shown in Fig. 3. For comparison, we also display constraints on the model from E787/E494 and MiniBooNE experiments following Ref. [34]. We also display the sensitivity reach projections for the future upgrade of XENON1T, the xenon-based DARWIN experiment [57], which is expected to have a size that is about 50 times larger with a fiducial mass of ~ 50 tons.

The effects of implementing an exact numerical treatment of attenuation that also includes the nuclear form factor are shown in Fig. 3 with the line “absorption,” above which DM flux would be attenuated if the approximate treatment has been used instead. This further highlights the importance of attenuation for identifying and exploring the optimal DM parameter space with direct DM experiments. Inclusion of additional mesons allows us to set new limits above few hundred MeV, as shown on Fig. 3. This is in contrast to earlier studies of Ref. [8] that focused solely on η and hence were not sensitive to this regime. Depending on the details of the full model as well as cosmological history, additional constraints could become relevant in the parameter space of interest, such as those from Big Bang nucleosynthesis and cosmic microwave background (see e.g. [34,58]). We stress, however, that our analysis results are independent of cosmology and the mechanism of DM production.

The limits found for the parameters on the model and up-quark coupling g_u can be translated into a stringent constraint for the spin-independent DM cross-section using Eq. (5). The results are

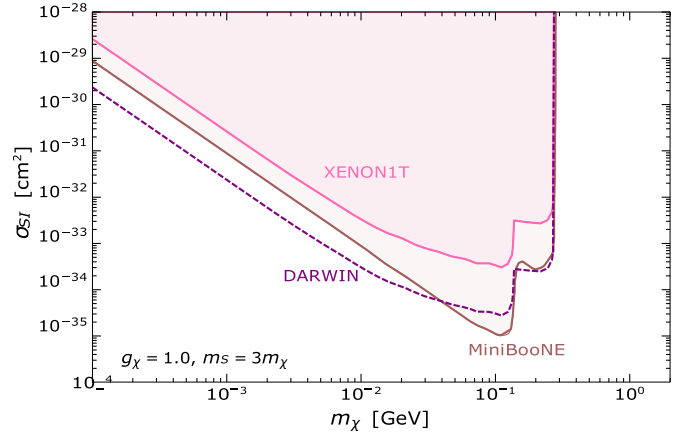


Fig. 4. Limits on the spin-independent DM nucleon cross-section σ^{SI} as a function of the DM mass m_χ .

shown in Fig. 4 for DM masses below the GeV scale. This highlights the significance of atmospheric DM production for exploration of light sub-GeV DM.

6. Conclusions

Light sub-GeV DM is a promising target for exploration beyond the traditional WIMP paradigm. However, the associated DM parameter space is challenging to probe with conventional direct DM detection experiments due to their low nuclear recoils. We have comprehensively explored production of hadrophilic light DM from cosmic ray collisions with the atmosphere, whose relativistic flux leads to observable experimental signatures. The decays of mesons resulting from collisions copiously source an irreducible flux of light DM for all terrestrial experiments. In this sense, our work applies to any hadrophilic scalars which couple to a new light neutral state, irrespective of whether or not that state accounts for the DM density.

Additionally, for the first time we analyze DM production associated with η , η' as well as K^+ mesons. We implement an exact treatment of DM flux attenuation as it traverses the Earth towards an underground detector, finding that nuclear form factor effects can dramatically impact the viable DM parameter space. We set new limits on hadrophilic light DM from XENON1T experiment and establish sensitivity projections for future experiments, such as DARWIN which will cover new low mass DM and low mass mediator parameter space. Future extensions of this work may include extending the DM model to examine inelastic DM tran-

sitions, which will impact both the recoil spectra and the DM attenuation. Another future extension would be to examine DM flux contributions beyond meson decay, such as scalar bremsstrahlung [34] which may provide access to higher mass DM.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

Data availability

Data will be made available on request.

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Appendix A. Meson decay scalar production

Sub-GeV scalar mediator S can be in meson decays from atmospheric collisions through a variety of channels. The resulting leading branching ratios for S can be found via chiral Lagrangian approach [34]:

$$\begin{aligned} \text{Br}(\eta \rightarrow \pi^0 S) &= \frac{C_\eta^2 g_u^2 B^2}{16\pi m_\eta \Gamma_\eta} \lambda_{1/2} \left(1, \frac{m_S^2}{m_\eta^2}, \frac{m_{\pi^0}^2}{m_\eta^2}\right), \\ \text{Br}(\eta' \rightarrow \pi^0 S) &= \frac{C_{\eta'}^2 g_u^2 B^2}{16\pi m_{\eta'} \Gamma_{\eta'}} \lambda_{1/2} \left(1, \frac{m_S^2}{m_{\eta'}^2}, \frac{m_{\pi^0}^2}{m_{\eta'}^2}\right), \\ \text{Br}(K^+ \rightarrow \pi^+ S) &= \frac{G_F^2 g_u^2 f_\pi^2 f_K^2 B^2}{8\pi m_{K^+} \Gamma_{K^+}} |V_{ud} V_{us}|^2 \\ &\quad \times \lambda_{1/2} \left(1, \frac{m_S^2}{m_{K^+}^2}, \frac{m_{\pi^+}^2}{m_{K^+}^2}\right), \end{aligned} \quad (\text{A.1})$$

where $B \simeq m_\pi^2/(m_u + m_d) \simeq 2.6$ GeV, $C_\eta(\eta') = \sqrt{1/3} \cos \theta' \mp \sqrt{2/3} \sin \theta'$ parametrize $\eta - \eta'$ mixing with $\theta' \simeq -20^\circ$, $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2bc - 2ac$, $G_F = 1.167 \times 10^{-5}$ GeV $^{-2}$ is the Fermi coupling constant, $|V_{ud}|$ and $|V_{us}|$ are the CKM matrix elements, f_π and f_K are decay constants and Γ_η , $\Gamma_{\eta'}$ and Γ_{K^+} are the decay widths [37]. Throughout we consider that $\text{Br}(S \rightarrow \chi \chi) = 1$.

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