

## Leptogenesis with small violation of $B - L$

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**ABSTRACT:** We analyze leptogenesis in the context of seesaw models with almost conserved lepton number, focusing on the  $L$ -conserving contribution to the flavoured CP asymmetries. We find that, contrary to previous claims, successful leptogenesis is feasible for masses of the lightest heavy neutrino as low as  $M_1 \sim 10^6$  GeV, without resorting to the resonant enhancement of the CP asymmetry for strongly degenerate heavy neutrinos. This lower limit renders thermal leptogenesis compatible with the gravitino bound in supersymmetric scenarios.

**KEYWORDS:** Neutrino Physics, Beyond Standard Model.

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## 1. Introduction

Two of the main evidences of physics beyond the Standard Model (SM), namely the observed (tiny) neutrino masses and the baryon asymmetry of the universe (BAU), are naturally explained in the context of the seesaw mechanism [1, 2, 3, 4]. In type I seesaw models, the SM is extended with at least two singlet Majorana neutrinos which can produce the observed BAU via leptogenesis [5, 6]: a lepton asymmetry is dynamically generated in the out of equilibrium decay of the heavy Majorana neutrinos, and then partially converted into a baryon asymmetry due to  $(B + L)$ -violating non-perturbative sphaleron interactions [7].

For a very hierarchical spectrum of heavy singlet neutrinos  $M_1 \ll M_2 \ll M_3$ , the  $L$ -violating CP asymmetry generated in the decay of the lightest singlet has an upper bound proportional to  $M_1$ , the so-called Davidson-Ibarra (DI) bound [8]. This implies a lower bound  $\sim 10^9$  GeV for the mass of the sterile neutrinos in order for  $N_1$ -dominated leptogenesis to be successful. Careful numerical studies show that the DI bound can be evaded for moderate hierarchies, e.g. the lower bound on  $M_1$  is relaxed by more than one order of magnitude with respect to the hierarchical limit one for  $M_3/M_2 \sim M_2/M_1 \sim 10$  [9]. However to reach these low values of  $M_1$  some unlikely cancellations are needed, which are not motivated by any underlying symmetry. Flavour effects [10, 11, 12, 13, 14, 15] do not substantially change this result. The authors of [16] make an analysis of the parameter space for successful leptogenesis relaxing also the condition of hierarchical neutrinos and including the  $L$ -conserving part of the flavoured CP asymmetries, finding a lower bound  $M_1 \gtrsim 10^8$  GeV.

Such lower bound on  $M_1$  in turn yields a lower bound for the reheating temperature,  $T_{RH}$ , of the same order, since to thermally produce the neutrinos  $M_1 \lesssim 5 T_{RH}$  [17, 18, 19]. On the other hand, in supersymmetric scenarios gravitinos are copiously produced in the high temperature plasma, and their late decay can jeopardize successful nucleosynthesis (BBN), leading to an upper bound on  $T_{RH}$ , which depends on the gravitino mass [20, 21].

If the gravitino is unstable, for gravitino masses,  $m_{3/2}$ , in the natural range from 100 GeV to 1 TeV, and within the minimal supergravity framework,  $T_{RH}$  should be smaller than  $10^5 - 10^7$  GeV, while for  $m_{3/2} \gtrsim 10$  TeV,  $T_{RH}$  can be of order  $10^9 - 10^{10}$  GeV [22]. As a consequence, in supersymmetric thermal leptogenesis there is some conflict between the gravitino bound on the reheat temperature and the thermal production of heavy neutrinos.

There are several possible ways out of this conflict: for instance, if the gravitino is stable, the nucleosynthesis bound depends on the next-to-lightest supersymmetric particle, but  $T_{RH} \gtrsim 10^9$  GeV can be obtained for  $m_{3/2} \gtrsim 10$  GeV [22]. Also, one can avoid the DI bound resorting to resonant leptogenesis, i.e., a resonant enhancement of the CP asymmetry which occurs when there are at least two strongly degenerated heavy neutrinos, such that  $M_2 - M_1 \sim \Gamma_N$ , being  $\Gamma_N$  their decay width [23, 24]. In this scenario, leptogenesis is feasible at much lower temperatures,  $T \sim \mathcal{O}(1 \text{ TeV})$  [25, 26, 27, 28].

Although such quasi-degeneracy of the heavy neutrinos might seem unnatural, there are well motivated seesaw models which yield a heavy neutrino quasi-degenerate spectrum, namely those with an approximately conserved  $B - L$  [29, 30]. In these models the tiny neutrino masses are proportional to small lepton number-breaking parameters, which are technically natural since a larger symmetry is realized when they vanish [31]. Moreover, the heavy neutrinos can be much lighter than in the generic seesaw, within the energy reach of LHC, and there can be a large active-sterile neutrino mixing. Also, lepton flavour violating rare decays as well as non-unitarity of the leptonic mixing matrix are present even in the limit of conserved  $B - L$ , and therefore they are unsuppressed by the light neutrino masses [32, 33, 30]. As a consequence, much attention has been devoted recently to this class of low scale seesaw models, since they have a rich phenomenology both at LHC [34, 35, 36] and at low energy charged lepton rare decay experiments, such as  $\mu \rightarrow e\gamma$ , and also lead to successful resonant leptogenesis [37, 38, 39].

In order to see more clearly the structure of seesaw models with small violation of  $B - L$ , it is useful to expand the low energy effective Lagrangian as

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \frac{c^{d=5}}{\Lambda_{LN}} \mathcal{O}^{d=5} + \sum_i \frac{c_i^{d=6}}{\Lambda_{FL}^2} \mathcal{O}_i^{d=6} + \dots, \quad (1.1)$$

where the dimensionless couplings  $c^{d=5}, c_i^{d=6}$  are assumed to be of  $\mathcal{O}(1)$ .

The only operator of  $d = 5$  is Weinberg's operator, responsible for neutrino masses, which is suppressed by the lepton number violating scale  $\Lambda_{LN}$ , while flavour changing but lepton number conserving  $d = 6$  operators  $\mathcal{O}_i^{d=6}$  are suppressed by a different scale  $\Lambda_{FL}$ . In models with approximately conserved lepton number, there is a separation of scales:  $\Lambda_{FL} \sim \mathcal{O}(\text{TeV})$  can be related to the heavy neutrinos mass scale, but  $\Lambda_{LN} \gg \Lambda_{FL}$ , since light neutrino masses only appear when small  $L$ -violating perturbations in the Yukawa couplings and/or the singlet neutrino mass matrix are introduced; in this scenario,  $\Lambda_{LN}$  does not correspond to any particle mass.

It has been noticed that even if the heavy neutrino that generates the BAU is not quasi-Dirac, or the mass splitting is outside the resonant regime, in seesaw models with almost conserved  $B - L$  the scale of leptogenesis can be lower than in the standard seesaw [40], provided flavour effects are at work. This is so because in these models leptogenesis

can be driven by the  $d = 6$  operators, which contribute only to the  $L$ -conserving part of the flavoured CP-asymmetries and escape the DI bound, even if the heavy neutrinos are hierarchical, because these operators are not linked to neutrino masses. The drawback is that the same  $d = 6$  operators induce large lepton flavour violating (but lepton number conserving) processes, which tend to equilibrate the flavoured asymmetries diminishing the total lepton asymmetry [41, 42].

In this paper we reanalyze the possibility of having successful leptogenesis driven by the purely flavoured  $L$ -conserving contribution to the CP asymmetries, in the context of seesaw models with small violation of  $B - L$ . In the numerical analysis of [16] such contribution was included, however models with almost conserved  $B - L$  (which can lead to larger CP asymmetries) were not considered in detail and the flavour equilibration processes were not taken into account. A study of the lower bound on  $M_1$  for approximately  $B - L$ -conserving models was carried out in [40], but their cross sections are different from ours and the analysis of the parameter space seems to be non-exhaustive. We have computed the cross sections for the crucial lepton flavour changing processes which lead to flavour equilibration, and we have found important differences with respect to previous results [26]. We have also taken into account decays and inverse decays of the next-to-lightest neutrinos, neglected in [40], and thoroughly scanned the parameter space to find the lowest  $M_1$  able to generate the BAU within this framework. Moreover, we have included a detailed discussion about leptogenesis in models with almost conserved  $B - L$ , which depends on the heavy neutrino spectrum.

The outline of this paper is as follows. In Section 2 we discuss the main features of leptogenesis within the framework of seesaw models with approximately conserved  $B - L$ . In Section 3 we write the set of Boltzmann equations (BE) relevant for leptogenesis in this scenario and compute the cross sections for the lepton flavour violating (but total lepton number conserving) processes, which play a crucial role in models with null (or negligible) total CP asymmetry. In Section 4 we perform a detailed analysis of the parameter space which leads to successful leptogenesis with only  $L$ -conserving CP asymmetries, and we conclude in Section 5.

## 2. Leptogenesis in models with small violation of $B - L$

In this section we want to determine what can be distinct, regarding leptogenesis, among different models with small violation of  $B - L$ . Let us call  $N_1$  the SM fermion singlet that is mainly responsible for the generation of the lepton asymmetry. Additionally we call  $N_2$  the SM fermion singlet that makes the most important (non resonant) virtual contribution to the CP asymmetry in the  $N_1$  decays. They have Yukawa interactions with the lepton doublets of the SM  $\ell_\alpha$  ( $\alpha = e, \mu, \tau$ ) described by the Lagrangian

$$\mathcal{L}_Y = -\lambda_{\alpha i} \tilde{h}^\dagger \overline{P_R N_i} \ell_\alpha - \lambda_{\alpha i}^* \bar{\ell}_\alpha P_R N_i \tilde{h} . \quad (2.1)$$

If  $B - L$  is only slightly violated, then for each  $N_i$  the conditions (i) or (ii) described below must be satisfied:

- (i)  $N_i$  is a Majorana neutrino with two degrees of freedom, whose Yukawa interactions violate lepton number and therefore the couplings  $\lambda_{\alpha i}$  must be small. This is analogous to the standard seesaw, so the  $N_i$  contribution to the  $d = 5$  Weinberg operator is given by

$$(c_M^{d=5})_{\alpha\beta} = \lambda_{\alpha i} \lambda_{\beta i} . \quad (2.2)$$

- (ii) The  $N_i$  is a Dirac or quasi Dirac neutrino with four degrees of freedom; this means that there are two Majorana neutrinos  $N_{ih}$  and  $N_{il}$  with masses  $M_i + \mu_i$  and  $M_i - \mu_i$  respectively. The parameter  $\mu_i \ll M_i$  measures the amount of  $B - L$  violation, so that if  $B - L$  is conserved,  $\mu_i = 0$  and  $N_i = (N_{ih} + iN_{il})/\sqrt{2}$  is a Dirac fermion. The Yukawa interactions can be expressed as

$$\mathcal{L}_{Y_{N_i}} = -\lambda_{\alpha i} \tilde{h}^\dagger P_R \frac{N_{ih} + iN_{il}}{\sqrt{2}} \ell_\alpha - \lambda'_{\alpha i} \tilde{h}^\dagger P_R \frac{N_{ih} - iN_{il}}{\sqrt{2}} \ell_\alpha + h.c., \quad (2.3)$$

where  $\lambda'_{\alpha i} \ll 1$ . The terms proportional to  $\lambda'_{\alpha i}$  induce lepton number violation even when  $\mu_i \rightarrow 0$  and hence they are similar in nature to the ones described in (i). Neglecting these terms, the Yukawa couplings of  $N_{ih}$  and  $N_{il}$  are equal to  $\lambda_{\alpha ih} = \frac{\lambda_{\alpha i}}{\sqrt{2}}$  and  $\lambda_{\alpha il} = i \frac{\lambda_{\alpha i}}{\sqrt{2}}$ , respectively. These  $\lambda_{\alpha i}$  can be large, because they do not vanish in the  $B - L$  conserved limit: in the absence of  $\mu_i$  and  $\lambda'_{\alpha i}$ , a perturbatively conserved lepton number can be defined, by assigning  $L_N = 1$  to  $N_i$ , and  $L_{\ell_\alpha} = 1$  to the SM leptons.

At leading order in the small  $B-L$  breaking parameters ( $\lambda'_{\alpha i}, \mu_i/M_i$ ), the contribution of a quasi Dirac heavy neutrino to the Weinberg operator is

$$\frac{(c_{QD}^{d=5})_{\alpha\beta}}{\Lambda_{LN}} = (\lambda'_{\alpha i} - \frac{\mu_i}{M_i} \lambda_{\alpha i}) \frac{1}{M_i} \lambda_{\beta i} + \lambda_{\alpha i} \frac{1}{M_i} (\lambda'_{\beta i} - \frac{\mu_i}{M_i} \lambda_{\beta i}) + \dots \quad (2.4)$$

As a result, one can reproduce the tiny light neutrino masses with large Yukawa couplings  $\lambda_{\alpha i}$  and heavy neutrino masses  $M_i$  as low as the TeV scale. Moreover, the admixture among singlet and doublet  $SU(2)$  neutral leptons (and the corresponding violation of unitarity in the light neutrino sector) is of order  $\lambda_{\alpha i} v/M_i$  and can be large.

In this work we will consider the cases

- (iia)  $\mu_i \ll \Gamma_{N_{ih}}, \Gamma_{N_{il}}$  (Dirac limit), and
- (iib)  $\Gamma_{N_{ih}}, \Gamma_{N_{il}} \ll \mu_i \ll M_i$  (Majorana limit),

where the decay widths  $\Gamma_s$  are given by

$$\Gamma_{N_{ih}} = \frac{M_i + \mu_i}{8\pi} \frac{(\lambda^\dagger \lambda)_{ii}}{2} \approx \frac{M_i - \mu_i}{8\pi} \frac{(\lambda^\dagger \lambda)_{ii}}{2} = \Gamma_{N_{il}} . \quad (2.5)$$

Instead we leave the study of  $\mu_i \sim \Gamma_{N_{ih,il}}$  for future research because this case cannot be described well with the simple BE used here [43, 44, 45, 46].

Then, the study of leptogenesis for a large class of models with small violation of  $B - L$  can be covered by considering the different combinations of options for both  $N_1$  and  $N_2$ . Next we comment on some key points for the different possibilities, while the complete quantitative analysis is made in Sec. 4. We assume that  $N_1$  is lighter than  $N_2$ , leaving some remarks of the opposite case  $M_2 < M_1$  for Sec. 4.

- I.  $N_1$  and  $N_2$  satisfy (i): This case is not very interesting since the CP asymmetries  $\epsilon_{\alpha 1}$  in the decays of  $N_1$  into leptons of flavour  $\alpha$ ,  $\epsilon_{\alpha 1} \equiv \frac{\Gamma(N_1 \rightarrow \ell_{\alpha h}) - \Gamma(N_1 \rightarrow \bar{\ell}_{\alpha \bar{h}})}{\sum_{\alpha} \Gamma(N_1 \rightarrow \ell_{\alpha h}) + \Gamma(N_1 \rightarrow \bar{\ell}_{\alpha \bar{h}})}$ , being proportional to the square of the Yukawa couplings of  $N_2$ , are very small, so it is not different from the standard seesaw.
- II.  $N_1$  satisfies (i) and  $N_2$  (ii): The  $L$ -violating part of the CP asymmetry,  $\epsilon_{\alpha 1}^{\cancel{L}}$ , is suppressed by  $\mu_2$  while the  $L$ -conserving part,  $\epsilon_{\alpha 1}^L$ , is not. To see this, let us write  $\epsilon_{\alpha 1} = \epsilon_{\alpha 1}^{\cancel{L}} + \epsilon_{\alpha 1}^L$  as [47, 48]

$$\epsilon_{\alpha 1} = \sum_{j=2h,2l} f(a_j) \text{Im} \left[ \lambda_{\alpha j}^* \lambda_{\alpha 1} (\lambda^\dagger \lambda)_{j1} \right] + \sum_{j=2h,2l} g(a_j) \text{Im} \left[ \lambda_{\alpha j}^* \lambda_{\alpha 1} (\lambda^\dagger \lambda)_{1j} \right], \quad (2.6)$$

where  $f(a_j)$  and  $g(a_j)$  are functions of  $a_j \equiv M_j^2/M_1^2$  and contain the factor  $1/(\lambda^\dagger \lambda)_{11}$ . The terms proportional to  $f(a_j)$  and  $g(a_j)$  come from the  $L$ -violating and  $L$ -conserving contributions, respectively. To lowest order in  $\mu_2$ ,  $f(a_{2h}) = f(a_{2l})$  and  $g(a_{2h}) = g(a_{2l})$ . Taking into account the alignment between the Yukawa couplings of  $N_{2l}$  and  $N_{2h}$ ,  $\lambda_{\alpha 2l} = a \lambda_{\alpha 2h}$  with  $a = i$ , it is clear that

$$\begin{aligned} \epsilon_{\alpha 1} &\xrightarrow{\mu_2 \rightarrow 0} f(a_{2h}) \text{Im} \left[ \lambda_{\alpha 2h}^* \lambda_{\alpha 1} (\lambda^\dagger \lambda)_{2h1} \right] (1 + a^{*2}) + \\ &g(a_{2h}) \text{Im} \left[ \lambda_{\alpha 2h}^* \lambda_{\alpha 1} (\lambda^\dagger \lambda)_{12h} \right] (1 + |a|^2) \\ &= 0 + 2g(a_{2h}) \text{Im} \left[ \lambda_{\alpha 2h}^* \lambda_{\alpha 1} (\lambda^\dagger \lambda)_{12h} \right]. \end{aligned} \quad (2.7)$$

Hence we see that the contributions from  $N_{2h}$  and  $N_{2l}$  add up in the  $L$ -conserving part and cancel in the  $L$ -violating one.

While the total CP asymmetry  $\epsilon_1 \equiv \sum_{\alpha} \epsilon_{\alpha 1}$  is proportional to  $\mu_2$  because  $\sum_{\alpha} \epsilon_{\alpha 1}^L = 0$ , the amount of matter-antimatter asymmetry produced during leptogenesis generally is not suppressed by  $\mu_2$  due to flavour effects (it is possible to have successful leptogenesis even with  $\epsilon_1 = 0$  [13]). Nevertheless, this asymmetry does vanish or becomes proportional to  $\mu_2$  when the  $L$ -violating parameter  $(\lambda^\dagger \lambda)_{11} \rightarrow 0$ , independently of whether the density of  $N_1$  at the onset of leptogenesis is null or thermal. This is because for an initial null density of heavy neutrinos and small values of  $(\lambda^\dagger \lambda)_{11}$  the final baryon asymmetry is proportional to  $(\lambda^\dagger \lambda)_{11}^2$  [19], while for an initial thermal  $N_1$  density the washouts become negligible and so do the flavour effects, in which case the total asymmetry generated is proportional to  $\epsilon_1 \propto \mu_2$ .

- III.  $N_1$  satisfies (ii) and  $N_2$  (i): Since the Yukawa couplings of  $N_2$  are small, the virtual contribution of  $N_2$  to the CP asymmetries  $\epsilon_{\alpha 1l}$  and  $\epsilon_{\alpha 1h}$  in the decays of  $N_{1l}$  and  $N_{1h}$ , respectively, is small. More interesting is the virtual contribution of  $N_{1l}$  to

$\epsilon_{\alpha 1h}$  and of  $N_{1h}$  to  $\epsilon_{\alpha 1l}$ , due to the resonant enhancement of the CP asymmetry for degenerate neutrinos. However these contributions are suppressed by  $\frac{\lambda'_{\alpha 1}}{\sqrt{(\lambda^\dagger \lambda)_{11}}}$  because of the alignment of the Yukawa couplings of  $N_{1l}$  and  $N_{1h}$  for  $\lambda'_{\alpha 1} = 0$ , therefore the CP asymmetry cannot reach the maximum value 1/2 of more generic resonant leptogenesis models. A detailed study of this case was made in [37], in the limit  $\mu_1 \gg \Gamma_{N_{1l}, N_{1h}}$ : they found that it is possible to lower the scale of leptogenesis provided that the parameter which controls lepton number violation in the Yukawa couplings  $\epsilon = \lambda'/\lambda$  is much larger than the one describing the mass splitting between the quasi-Dirac heavy neutrinos,  $\epsilon_M = \Delta M/M$ . For instance, in order to obtain successful leptogenesis with  $M = 10^6$  GeV (1 TeV), one needs  $\epsilon \sim 10^{-3}$  and  $\epsilon_M \sim 10^{-8}$  ( $10^{-11}$ ). The case  $\mu_1 \lesssim \Gamma_{N_{1l}, N_{1h}}$  has been considered in [38], however it has been shown that in this maximal resonant regime the classical Boltzmann picture breaks down [43, 44, 45, 46].

IV.  $N_1$  and  $N_2$  satisfy (ii): Besides the resonant contributions which are the same as in III, now also the virtual contribution of  $N_2$  to the CP asymmetries  $\epsilon_{\alpha 1l}$  and  $\epsilon_{\alpha 1h}$  can be large, so we focus on these terms. Neglecting the  $\lambda'_{\alpha i} \ll \lambda_{\alpha i}$  couplings, the CP asymmetries in the decays of  $N_{1l}$  and  $N_{1h}$  become equal. Here it is very important to distinguish between the cases  $\mu_1 \gg \Gamma_{N_{1l}, N_{1h}}$  and  $\mu_1 \ll \Gamma_{N_{1l}, N_{1h}}$ . In the first one,  $N_{1l}$  and  $N_{1h}$  behave as two independent Majorana neutrinos regarding the generation of the lepton asymmetry, which can hence be roughly double with respect to II. However in the second case  $N_1$  is (or effectively behaves as) a Dirac neutrino, i.e. lepton number is conserved in its decay, and therefore the only possibilities to end up with a non-zero baryon asymmetry is to have important washouts from the two Majorana components of  $N_2$  (if  $\mu_2 \gg \Gamma_{N_{2l}, N_{2h}}$ ) or let the sphalerons freeze out during leptogenesis [49].

Summarizing, in models with small violation of  $B - L$  the CP asymmetries in the decay of  $N_1$  can be enhanced in cases II, III and IV. Since the resonant contribution to the  $L$ -violating CP asymmetry has been widely studied, we develop in Sec. 4 a quantitative analysis which covers all the non-resonant interesting cases, i.e., II and IV with  $\mu_1 \gg \Gamma_{N_{1l}, N_{1h}}$ . With respect to  $N_2$ , we consider two possibilities,  $\mu_2 \gg \Gamma_{N_{2l}, N_{2h}}$  and  $\mu_2 \ll \Gamma_{N_{2l}, N_{2h}}$  (see [49] for leptogenesis with both  $N_1$  and  $N_2$  satisfying  $\mu_i \ll \Gamma_{N_{il}, N_{ih}}$ ).

### 3. Boltzmann equations

Motivated by the above discussion we consider a scenario for leptogenesis involving three fermion singlets  $N_1, N_{2l}, N_{2h}$  (each of them having two degrees of freedom), with respective masses  $M_1, M_2 - \mu_2, M_2 + \mu_2$  and Yukawa couplings given by the Lagrangian

$$\mathcal{L}_Y = -\lambda_{\alpha 1} \tilde{h}^\dagger \overline{P_R N_1} \ell_\alpha - \lambda_{\alpha 2} \tilde{h}^\dagger P_R \frac{N_{2h} + iN_{2l}}{\sqrt{2}} \ell_\alpha + h.c. . \quad (3.1)$$

Recall that the parameters  $\lambda_{\alpha 1}$  violate lepton number and hence  $\lambda_{\alpha 1} \ll \lambda_{\alpha 2}$ . In turn this implies that the CP asymmetry in  $N_1$  decays is the dominant one. As a first approximation,

in the Eq. (3.1) we have neglected the  $L$ -violating couplings  $\lambda'_{\alpha 2} \ll \lambda_{\alpha 2}$ , because we expect that their contributions to the CP asymmetries and washouts are negligible. We have checked indeed that this is the case in the parameter space region relevant for leptogenesis that can also accommodate the observed light neutrino masses (see Sec. 4). As we will explain below, it is convenient to take  $M_1 < M_2$  in order to obtain the lowest energy scale for leptogenesis within this framework.

The amount of leptons and antileptons present in the Universe can be described by density matrices in flavour space. The evolution equations of these density matrices take the simplest form in the basis that diagonalize them, which are determined by the fastest interactions in flavour space (if there is a hierarchy among the different interactions). Let us suppose that the fastest interactions during  $N_1$ -leptogenesis are the  $N_2$ -Yukawa interactions<sup>1</sup>. In this case the BE would be diagonal in a basis  $(\ell_2, \ell_{2\perp}, \ell'_{2\perp})$ , with  $\ell_{2\perp}$  and  $\ell'_{2\perp}$  two orthogonal lepton flavour states perpendicular to the  $N_2$ -decay eigenstate  $\ell_2$ . But  $\epsilon(N_1 \rightarrow \ell_{2\perp} h, \ell'_{2\perp} h) = 0$  (with  $\epsilon(N_1 \rightarrow \ell_a h)$  the CP asymmetry in the decay of  $N_1$  into  $\ell_a h$ ), and since the total CP asymmetry is null, then also  $\epsilon(N_1 \rightarrow \ell_2 h) = 0$ , with the result that no lepton asymmetry would be produced. Moreover, no asymmetry is generated when the  $N_1$ -Yukawa interactions are the fastest ones, because in this case the BE would be diagonal in the basis  $(\ell_1, \ell_{1\perp}, \ell'_{1\perp})$ , and since the CP asymmetry in the  $N_1$ -decay eigenstate  $\epsilon(N_1 \rightarrow \ell_1 h)$  coincides with the total CP asymmetry, it also vanishes. So we demand that the couplings of  $N_1$  and  $N_2$  be small enough, such that the Yukawa interactions of the  $\tau$  are the dominant ones. This implies that the BE are diagonal in the orthogonal basis  $(\ell_\tau, \ell_{\tau\perp}, \ell'_{\tau\perp})$ , with  $\ell_{\tau\perp}$  and  $\ell'_{\tau\perp}$  being determined by the fastest interaction acting in the plane perpendicular to  $\ell_\tau$ . For simplicity we take the  $N_1$  and  $N_2$  decay eigenstates to be perpendicular to  $\ell_e$ , in which case the aforementioned basis is  $\ell_\tau, \ell_\mu, \ell_e$ . Since no asymmetry is generated in the  $\ell_e$  flavour, two BE will be enough to determine the evolution of the lepton asymmetry and consequently the flavour indices  $\alpha$  and  $\beta$  run over the species  $\mu$  and  $\tau$ . Later we will comment on the more general case with three non-null flavour asymmetries.

Besides the charged lepton Yukawa interactions, the most relevant processes are the decays and inverse decays of the different heavy neutrinos, and the  $L$ -conserving but  $L_\alpha$ -violating scatterings  $\ell_\beta h \rightarrow \ell_\alpha h$ ,  $\ell_\beta \bar{h} \rightarrow \ell_\alpha \bar{h}$ , and  $h \bar{h} \rightarrow \ell_\alpha \bar{\ell}_\beta$ , hereafter called generically flavour changing interactions (FCI). We do not consider finite temperature corrections to the particle masses and couplings [18], moreover we also neglect spectator processes during leptogenesis and the asymmetry developed among the degrees of freedom of the Higgs [50, 51], as well as  $\Delta L = 1$  scatterings [52, 53]<sup>2</sup>. The relevant set of BE for the case

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<sup>1</sup>When the effects of  $N_{2l}$  and  $N_{2h}$  simply add up we will refer to these states generically as  $N_2$ . In the models considered here the only situation which requires a differentiated treatment of the degrees of freedom associated to  $N_2$  is when  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ , because a density asymmetry can develop between the states produced by  $\ell h$  and  $\bar{\ell} \bar{h}$  (see below).

<sup>2</sup>We have checked that the inclusion of spectator processes and  $\Delta L = 1$  scatterings modifies the results by at most a few tens of percent.

$\mu_2 \gg \Gamma_{N_{2l,2h}}$  is

$$\frac{dY_{N_1}}{dz} = \frac{-1}{sHz} \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \gamma_{D_1}, \quad (3.2)$$

$$\begin{aligned} \frac{dY_{\Delta_\alpha}}{dz} = \frac{-1}{sHz} \left\{ \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \epsilon_{\alpha 1} \gamma_{D_1} - \sum_i \gamma_{\ell_\alpha h}^{N_i} y_{\ell_\alpha} \right. \\ \left. - \sum_{\beta \neq \alpha} \left( \gamma_{\ell_\alpha h}^{\ell_\beta h'} + \gamma_{\ell_\alpha \bar{h}}^{\ell_\beta \bar{h}} + \gamma_{\ell_\alpha \bar{\ell}_\beta}^{h \bar{h}} \right) [y_{\ell_\alpha} - y_{\ell_\beta}] \right\}, \quad (3.3) \end{aligned}$$

where  $z \equiv M_1/T$  (with  $T$  the temperature),  $H$  is the Hubble expansion rate,  $Y_X \equiv n_X/s$  is the number density of a single degree of freedom of the particle species  $X$  normalized to the entropy density  $s$ ,  $y_X \equiv (Y_X - Y_{\bar{X}})/Y_X^{eq}$  is the asymmetry density normalized to the equilibrium density  $Y_X^{eq}$ , and  $Y_{\Delta_\alpha} \equiv Y_B/3 - Y_{L_\alpha}$ , with  $Y_B$  the baryon asymmetry and  $Y_{L_\alpha}$  the lepton asymmetry in the flavour  $\alpha$ . Since we are neglecting spectator processes during leptogenesis, the asymmetry in the lepton doublets  $\ell_\alpha$  can be expressed in terms of  $Y_{\Delta_\alpha}$  using the simple relation  $Y_{\Delta_\alpha} = -Y_{L_\alpha} = -2y_{\ell_\alpha} Y_{\ell_\alpha}^{eq}$  [51]. The sphaleron processes convert part of the final  $B - L$  asymmetry ( $Y_{B-L}^f = \sum_\alpha Y_{\Delta_\alpha}^f$ ) into a baryon asymmetry ( $Y_B^f$ ) and the relation between these quantities is taken to be  $Y_B^f = 28/79 Y_{B-L}^f$  [54]. In Eq. (3.3) we have also introduced the notation  $\gamma_{c,d,\dots}^{a,b,\dots} \equiv \gamma(a,b,\dots \rightarrow c,d,\dots)$  for the reaction density of the process  $a,b,\dots \rightarrow c,d,\dots$ , the prime in  $\gamma_{\ell_\alpha h}^{\ell_\beta h'}$  indicating that the on-shell contribution has to be subtracted, and  $\gamma_{D_i} \equiv \sum_\alpha \gamma_{\ell_\alpha h}^{N_i} + \gamma_{\bar{\ell}_\alpha \bar{h}}^{N_i}$ . Finally notice that the lepton asymmetries generated by  $N_2$  decays,  $\epsilon_{\alpha 2}$ , have not been taken into account, given that  $\epsilon_{\alpha 2} \ll \epsilon_{\alpha 1}$ .

Instead, if  $\mu_2 \ll \Gamma_{N_{2l,2h}}$  then  $N_{2l}$  and  $N_{2h}$  combine to form a Dirac neutrino  $N_2 \equiv (N_{2h} + iN_{2l})/\sqrt{2}$ , and therefore there is an asymmetry generated among the degrees of freedom of  $N_2$  which has to be taken into account [49]. An appropriate set of BE for this case is

$$\frac{dY_{N_1}}{dz} = \frac{-1}{sHz} \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \gamma_{D_1}, \quad (3.4)$$

$$\frac{dY_{N_2 - \bar{N}_2}}{dz} = \frac{-1}{sHz} \sum_\alpha \gamma_{\ell_\alpha h}^{N_2} [y_{N_2} - y_{\ell_\alpha}], \quad (3.5)$$

$$\begin{aligned} \frac{dY_{\Delta_\alpha}}{dz} = \frac{-1}{sHz} \left\{ \left( \frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \epsilon_{\alpha 1} \gamma_{D_1} - \gamma_{\ell_\alpha h}^{N_1} y_{\ell_\alpha} + \gamma_{\ell_\alpha h}^{N_2} [y_{N_2} - y_{\ell_\alpha}] \right. \\ \left. - \sum_{\beta \neq \alpha} \left( \gamma_{\ell_\alpha h}^{\ell_\beta h'} + \gamma_{\ell_\alpha \bar{h}}^{\ell_\beta \bar{h}} + \gamma_{\ell_\alpha \bar{\ell}_\beta}^{h \bar{h}} \right) [y_{\ell_\alpha} - y_{\ell_\beta}] \right\}. \quad (3.6) \end{aligned}$$

The reaction densities for the FCI are obtained integrating the corresponding cross section  $\sigma_{\text{FCI}}$ ,  $\gamma_{\text{FCI}} = \frac{T}{64\pi^4} \int_0^\infty ds s^{1/2} \hat{\sigma}_{\text{FCI}} K_1 \left( \frac{\sqrt{s}}{T} \right)$ , with the reduced cross sections  $\hat{\sigma}_{\text{FCI}} \equiv 2s\sigma_{\text{FCI}}$ . In order to give general expressions for these cross sections we take the FCI to be mediated by any number of Dirac or Majorana neutrinos  $N_i$  ( $i = 1, 2, \dots$ ) with masses  $M_i$ ,

whose Yukawa couplings with the lepton doublets  $\ell_\alpha$  are called  $h_{\alpha i}$ . We find <sup>3</sup>:

$$\begin{aligned}
\hat{\sigma}(\ell_\beta h \rightarrow \ell_\alpha h) &= \frac{1}{4\pi} \sum_{i,j} h_{\alpha i}^* h_{\beta i} h_{\alpha j} h_{\beta j}^* \frac{s^2 (s - M_i^2 - iM_i \Gamma_i) (s - M_j^2 + iM_j \Gamma_j)}{\left[ (s - M_i^2)^2 + M_i^2 \Gamma_i^2 \right] \left[ (s - M_j^2)^2 + M_j^2 \Gamma_j^2 \right]}, \\
\hat{\sigma}(\ell_\beta \bar{h} \rightarrow \ell_\alpha \bar{h}) &= \frac{1}{2\pi} \sum_{\substack{i,j \\ i \neq j}} h_{\alpha i}^* h_{\beta i} h_{\alpha j} h_{\beta j}^* \frac{1}{M_i^2 - M_j^2} \left\{ M_i^2 \ln \frac{s + M_i^2}{M_i^2} - M_j^2 \ln \frac{s + M_j^2}{M_j^2} \right\} \\
&\quad + \frac{1}{2\pi} \sum_{\substack{i,j \\ i=j}} h_{\alpha i}^* h_{\beta i} h_{\alpha j} h_{\beta j}^* \left\{ \ln \frac{s + M_i^2}{M_i^2} - \frac{s}{s + M_i^2} \right\}, \\
\hat{\sigma}(h \bar{h} \rightarrow \ell_\alpha \bar{\ell}_\beta) &= \frac{1}{2\pi} \sum_{\substack{i,j \\ i \neq j}} h_{\alpha i}^* h_{\beta i} h_{\alpha j} h_{\beta j}^* \left\{ -1 + \frac{M_i^2 (s + M_i^2)}{s(M_i^2 - M_j^2)} \ln \frac{s + M_i^2}{M_i^2} + \right. \\
&\quad \left. \frac{M_j^2 (s + M_j^2)}{s(M_j^2 - M_i^2)} \ln \frac{s + M_j^2}{M_j^2} \right\} + \\
&\quad \frac{1}{2\pi} \sum_{\substack{i,j \\ i=j}} h_{\alpha i}^* h_{\beta i} h_{\alpha j} h_{\beta j}^* \left\{ -2 + \frac{s + 2M_i^2}{s} \ln \frac{s + M_i^2}{M_i^2} \right\}. \tag{3.7}
\end{aligned}$$

Moreover, the subtracted reaction density is given by

$$\gamma_{\ell_\alpha h}^{\ell_\beta h'} = \gamma_{\ell_\alpha h}^{\ell_\beta h} - \sum_i \gamma_{N_i}^{\ell_\beta h} \text{Br}(N_i \rightarrow \ell_\alpha h), \tag{3.8}$$

with  $\text{Br}(N_i \rightarrow \ell_\alpha h)$  the branching ratio of the process  $N_i \rightarrow \ell_\alpha h$ .

For the models we are considering the dominant contribution to the FCI comes from the exchange of a neutrino with mass  $M_2$  and Yukawa couplings  $h_{\alpha 2} = \lambda_{\alpha 2}$  <sup>4</sup>.

## 4. Results

We will find the lower bound for  $M_1$  under the requirement that leptogenesis be successful in these scenarios. The relevant parameters for leptogenesis are  $M_1$ ,  $M_2/M_1$ ,  $(\lambda^\dagger \lambda)_{11}$ ,  $(\lambda^\dagger \lambda)_{22}$ , the projectors  $K_{\alpha i} \equiv \lambda_{\alpha i} \lambda_{\alpha i}^* / (\lambda^\dagger \lambda)_{ii}$  ( $\alpha = \mu, \tau$ ;  $i = 1, 2$ ), and  $\mu_2$ . Next we comment on the role of these parameters.

- $M_1$ : For a fix value of the Yukawa couplings (and in particular of the CP asymmetry when  $M_2/M_1$  is kept fixed) the lower is  $M_1$ , the stronger the washouts become (because the expansion rate becomes slower with decreasing  $T$ ). This is why  $M_1$  cannot be very low.

<sup>3</sup>Our expressions for these cross sections differ from the ones used in [40], which were taken from [26].

<sup>4</sup>Note that the contribution to the FCI of the two Majorana neutrinos  $N_{2l}, N_{2h}$  with the Yukawa couplings of Eq. (3.1) is the same as the contribution of one Dirac neutrino with mass  $M_2$  and Yukawa couplings  $\lambda_{\alpha 2}$ , in the limit  $\mu_2 \ll M_2$ .

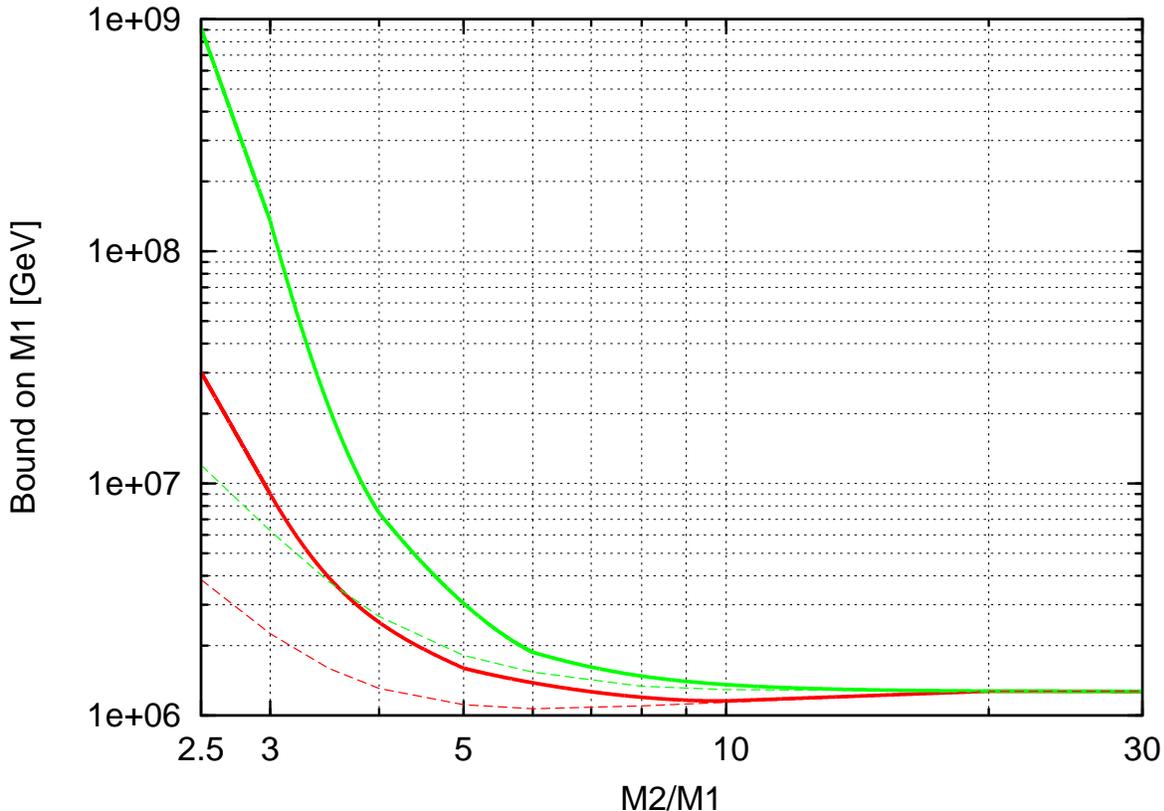
- $M_2/M_1$ : The hierarchy among the masses appears in the CP asymmetry, with  $\epsilon_{\alpha 1} \propto (M_2/M_1)^{-2}$  for  $M_2 \gg M_1$ , and in the FCI when these are mediated mainly by  $N_2$  (as always happens in the interesting cases for this work),  $\gamma_{\text{FCI}}(T) \propto (M_2/M_1)^{-4}$  for  $T \sim M_1 \ll M_2$ . Also note that if  $M_2/M_1 \lesssim 20$  the inverse decays of  $N_2$  could erase part of the asymmetry generated during  $N_1$  leptogenesis and therefore they should be included in the BE.
- $(\lambda^\dagger \lambda)_{11}$ : The intensity of the washouts due to processes involving  $N_1$  is determined by the effective mass  $\tilde{m}_1 \equiv (\lambda^\dagger \lambda)_{11} v^2 / M_1$ . In order to make full use of the -mandatory- flavour effects we take  $\tilde{m}_1 \gtrsim m_* \simeq 10^{-3}$  eV<sup>5</sup>. However, since  $(\lambda^\dagger \lambda)_{11}$  violates lepton number it cannot be very large in the scenarios we are considering.
- $(\lambda^\dagger \lambda)_{22}$ : The CP asymmetry  $\epsilon_{\alpha 1}$  is directly proportional to  $(\lambda^\dagger \lambda)_{22}$ , hence this parameter should be taken as large as possible. There are two reasons that prevent it from being very large. One is that the washouts due to processes involving  $N_2$  increase with  $(\lambda^\dagger \lambda)_{22}$ . The other is that, as explained at the beginning of Sec. 3, the Yukawa interactions of  $N_2$  must be slower than those of the  $\tau$ . Below we explain in more detail how this constraint has been dealt with.
- $K_{\alpha i}$ : The flavoured CP asymmetries depend on the square root of the projectors while the washouts in the flavour “ $\alpha$ ” depend linearly on  $K_{\alpha i}$ , hence the washouts decrease faster than the CP asymmetries with decreasing projectors. This fact must be taken into account in order to maximize the production of lepton asymmetry. Note that since  $\sum_\alpha K_{\alpha i} = 1$  and we are taking  $K_{ei} = 0$ , there are only two independent projectors, which can be chosen as  $K_{\mu 1}$  and  $K_{\mu 2}$ .
- $\mu_2$ : Since we are not considering the case  $\mu_2 \sim \Gamma_{N_{2l,2h}}$ ,  $\mu_2$  only enters as a discrete parameter, the baryon asymmetry taking one  $\mu_2$ -independent value when  $\mu_2 \gg \Gamma_{N_{2l,2h}}$  and another when  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ .

Given that the hierarchy  $M_2/M_1$  is an interesting and crucial parameter, we have determined the minimum value of  $M_1$  compatible with successful leptogenesis as a function of  $M_2/M_1$ , maximizing the final baryon asymmetry over the remaining parameters, i.e. over the relevant combinations of Yukawa couplings:  $(\lambda^\dagger \lambda)_{11}$ ,  $(\lambda^\dagger \lambda)_{22}$ ,  $K_{\mu 1}$ , and  $K_{\mu 2}$ . To obtain the baryon asymmetry we have solved numerically the appropriate set of BE, and to get successful leptogenesis we have required  $Y_B = 8.75 \times 10^{-11}$  [55]. The result is represented with the thick continuous curves in Fig. 1, the red line corresponding to the case  $\mu_2 \gg \Gamma_{N_{2l,2h}}$  and the green one to  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ .

As can be seen it is possible to have neutrino masses as low as  $M_1 \sim 10^6$  GeV, i.e. around three orders of magnitude below the lower bound for the standard case of type I seesaw with hierarchical heavy neutrinos. An interesting consequence is that this value

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<sup>5</sup>The quantity  $m_*$  is the *equilibrium mass*, which is defined by the condition  $\frac{\Gamma_{N_1}}{H(T=M_1)} = \frac{\tilde{m}_1}{m_*}$ , so that  $m_* = \frac{16}{3\sqrt{5}} \pi^{5/2} \sqrt{g_{*SM}} \frac{v^2}{m_{pl}} \simeq 1,08 \times 10^{-3}$  eV ( $g_{*SM}$  is the number of SM relativistic degrees of freedom at temperature  $T$  and  $m_{pl}$  is the Planck mass).



**Figure 1:** Lowest value of  $M_1$  yielding successful leptogenesis as a function of  $M_2/M_1$ . The red curves are for the case  $\mu_2 \gg \Gamma_{N_{2l,2h}}$  and the green ones for  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ . The thick continuous curves give the physically correct bound, while the thin dashed ones show the result that would be obtained if the Yukawa couplings of  $N_2$  were allowed to take values as large as 1 for all values of  $M_2/M_1$ .

can be compatible with the upper bound on the reheating temperature required to avoid the gravitino problem in SUGRA models. Moreover,  $M_1$  values around  $10^6$  GeV can be achieved for a wide range of  $N_2$  masses and also for different values of the Yukawa couplings. In particular, the maximum value of the baryon asymmetry is obtained for  $\tilde{m}_1 \sim 10^{-2}$  eV and  $K_{\mu 1}$  (or  $K_{\tau 1}$ )  $\sim 0.1$ , but it does not change much if  $\tilde{m}_1$  is varied within the range  $5 \times 10^{-3}$  eV  $\lesssim \tilde{m}_1 \lesssim 10^{-1}$  eV as long as the projectors are adjusted in order to have weak washouts in one flavour and strong washouts in the other, e.g.  $K_{\mu 1} \tilde{m}_1 \lesssim m_*$  and  $K_{\tau 1} \tilde{m}_1 \gtrsim (5 - 10)m_*$ . Regarding the Yukawa couplings of  $N_2$ ,  $(\lambda^\dagger \lambda)_{22}$  lies approximately between<sup>6</sup> 0.01 and 1, while for intermediate to large  $N_2$  masses,  $M_2 \gtrsim 5M_1$ , the smallest projector can take values as large as 0.1 to 0.5, without changing the bound shown in Fig. 1 by more than a factor 2. Instead, for  $M_2 \lesssim 5M_1$ , some hierarchy among the Yukawa couplings of  $N_2$  is required, namely  $\lambda_{\mu 2}/\lambda_{\tau 2}$  (or  $\lambda_{\tau 2}/\lambda_{\mu 2}$ )  $\sim (1 - 3) \times 10^{-2}$  to achieve the

<sup>6</sup>In the analysis presented here we have restricted  $|\lambda_{\alpha 2}| \leq 1$ , but we have checked that the results do not change significantly when allowing these Yukawa couplings to take somewhat larger values (but below the perturbative bound).

values shown in Fig. 1.

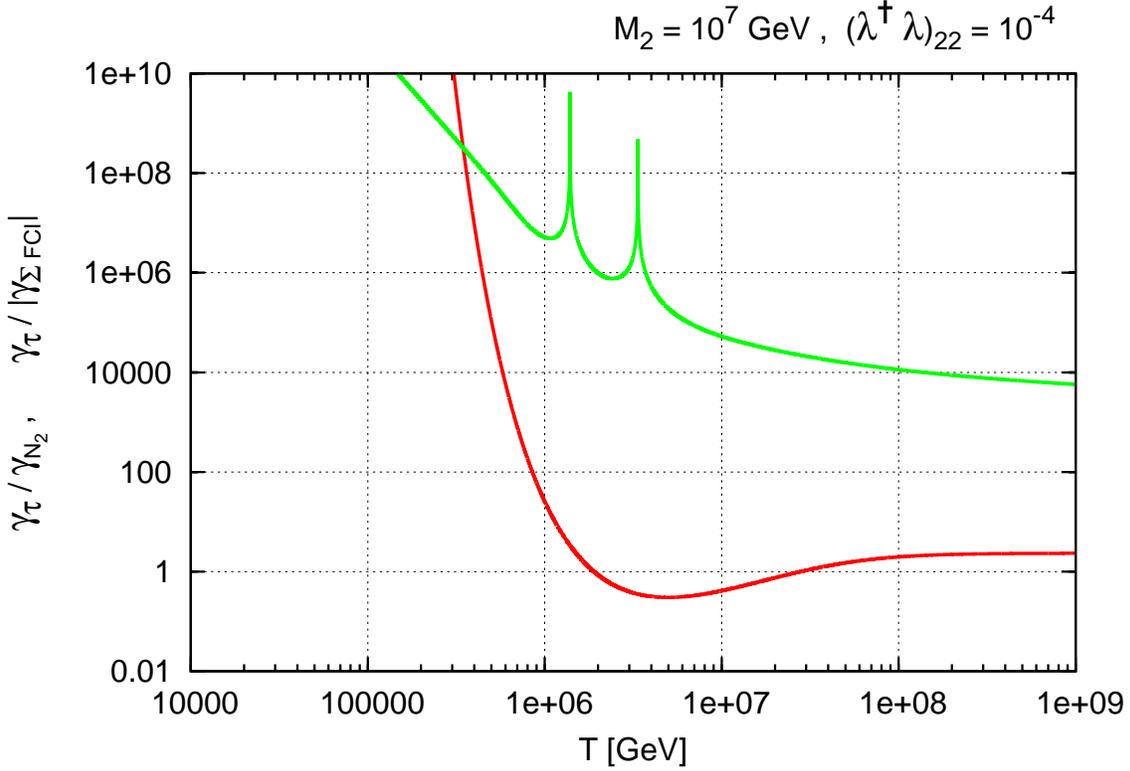
An important issue for obtaining the bound on  $M_1$  is to determine how large the Yukawa couplings of  $N_2$  can be without violating the condition that the rates of processes involving  $N_2$  be slower than the rates of the  $\tau$ -Yukawa interactions<sup>7</sup>. For reference purposes we plot in Fig. 2 the ratio between the rates of these interactions, distinguishing between processes with a real or virtual  $N_2$ , for  $M_2 = 10^7$  GeV and  $\sqrt{(\lambda^\dagger\lambda)_{22}} = 0.01 \sim h_\tau$  (the rates of the  $\tau$ -Yukawa interactions have been taken from [56, 57]). On one hand, note that the  $\tau$ -Yukawa interactions tend to become dominant over the  $N_2$  ones as the temperature decreases. On the other hand, for each set of parameters there exists a temperature  $T_{in}$  such that the final baryon asymmetry ( $Y_B^f$ ) does not depend on what happens at  $T > T_{in}$  (because there are strong washouts for all the interesting regions of the parameter space). Then the largest possible value for  $(\lambda^\dagger\lambda)_{22}$  has been set by requiring that the  $\tau$ -Yukawa interactions be faster than the  $N_2$  ones for all temperatures *below*  $T_{in}$ . In practice we have determined  $T_{in}$  finding the lowest value of  $T$  such that  $Y_B^f$  does not change by more than 10% when the initial conditions at  $T$ , namely  $Y_{N_i}(T)$  and  $Y_{\Delta_\alpha}(T)$ , are varied. The value of  $T_{in}$  as a function of  $M_2/M_1$  is shown in Fig. 3, reinterpreted as an approximate lower bound for the reheating temperature. Also note that the independence from the initial conditions is an interesting feature by itself, since the models become more predictive.

For comparison we have also plotted in Fig. 1 the -wrong- bound that would be obtained if  $(\lambda^\dagger\lambda)_{22}$  were allowed to be as large as 1, ignoring the above discussion. It is clear that as  $M_2$  approaches  $M_1$  the requirement of an upper bound for  $(\lambda^\dagger\lambda)_{22}$  becomes very relevant. This requirement is also the reason why it is convenient to take  $M_2 > M_1$ , because in the opposite case (when the lepton asymmetry is produced by the decay of the next-to-lightest singlet neutrino and that is usually called “ $N_2$  leptogenesis”) the bound on  $(\lambda^\dagger\lambda)_{22}$  would not allow for large values of the CP asymmetry.

Another feature apparent in Fig. 1 is the constant behaviour of the bound on  $M_1$  for large values of  $M_2/M_1$ . This can be understood quite easily as follows. The asymmetry  $Y_B^f$  is proportional to the CP asymmetry in  $N_1$  decays,  $\epsilon_{\mu 1} = -\epsilon_{\tau 1}$ , which in turn satisfies  $\epsilon_{\mu 1} \propto (M_1/M_2)^2(\lambda_{\mu 2}\lambda_{\tau 2})$ . For  $M_2 \gg M_1$ ,  $Y_B^f$  increases with  $(\lambda_{\mu 2}\lambda_{\tau 2})$  up to a certain value  $(\lambda_{\mu 2}\lambda_{\tau 2})_{\max}$  for which the FCI become important, equilibrating the asymmetries generated in the two flavours with the consequent decrease of the baryon asymmetry [41, 42]. At the temperatures relevant for leptogenesis,  $T \sim M_1 \ll M_2$ , the rates of the FCI scale as  $\gamma_{\text{FCI}} \propto (M_1/M_2)^4(\lambda_{\mu 2}\lambda_{\tau 2})^2$ . Hence we see that if  $(M_1/M_2)$  is decreased by a factor  $a$ ,  $(\lambda_{\mu 2}\lambda_{\tau 2})_{\max}$  can be increased by a factor  $a^2$  keeping constant the rates of the FCI. Therefore  $(\lambda_{\mu 2}\lambda_{\tau 2})_{\max} \propto (M_1/M_2)^{-2}$  and hence the maximum value of  $Y_B^f$ , being proportional to  $(M_1/M_2)^2(\lambda_{\mu 2}\lambda_{\tau 2})_{\max}$ , becomes independent of  $M_2/M_1$ .

As mentioned in Sec. 3, for simplicity we have taken  $\ell_e$  to be perpendicular to the decay eigenstates of  $N_1$  and  $N_2$ , so that only two flavour asymmetries are generated. We have checked that in the more general three flavour case it is possible to lower the bound on  $M_1$  by a factor up to almost 4 with respect to the two flavour case. This is due to the combination of two effects. On one hand,  $Y_B^f$  can be almost twice as large compared to

<sup>7</sup>There are not relevant experimental bounds on the Yukawa couplings of  $N_2$  for the masses of heavy neutrinos we are considering.



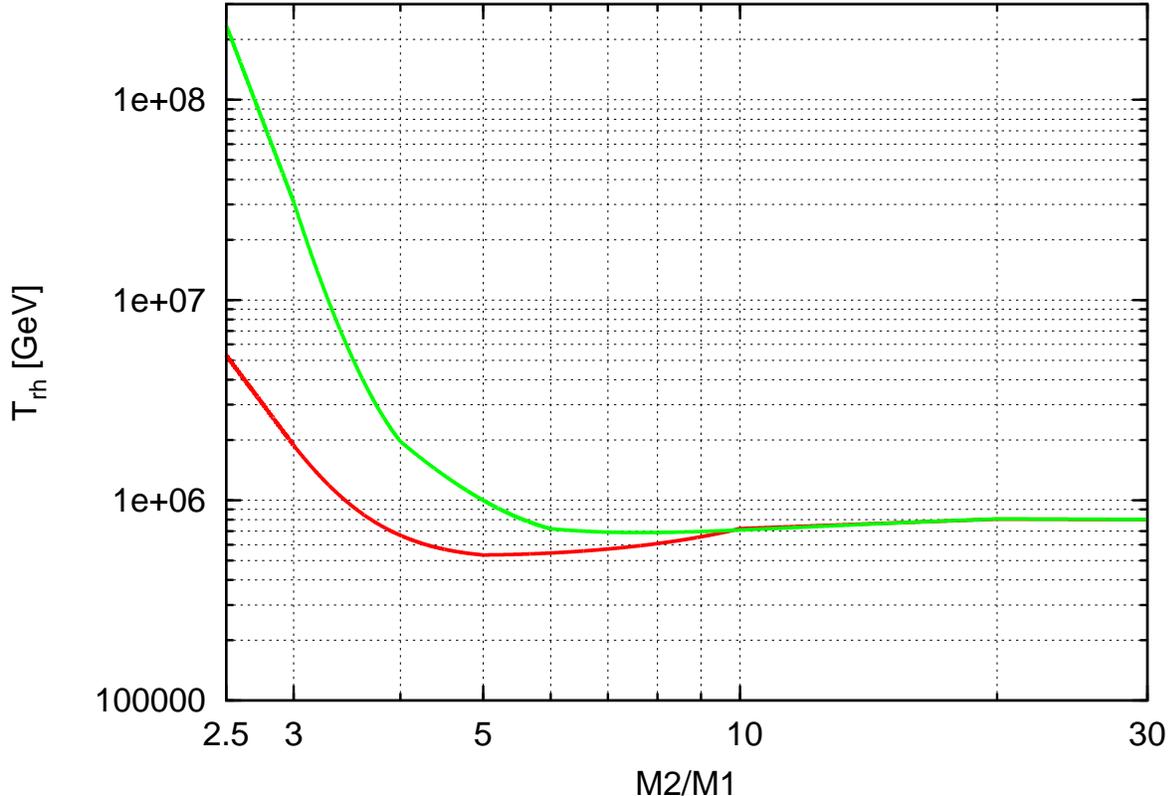
**Figure 2:** Comparison between the rates of the Yukawa interactions of  $N_2$  and  $\tau$  as a function of the temperature. The red line gives  $\gamma_\tau/\gamma_{N_2}$  and the green one  $\gamma_\tau/|\gamma_{\Sigma\text{FCI}}|$ , where  $\gamma_\tau$  is the rate of the Yukawa interactions of the  $\tau$ ,  $\gamma_{N_2}$  is the sum of the rates of processes involving a real  $N_2$  (here we have included the decay and scatterings with the top quark summing over all lepton flavours), and  $\gamma_{\Sigma\text{FCI}}$  is basically the sum of the rates of the FCI mediated by an off-shell  $N_2$ ,  $\gamma_{\Sigma\text{FCI}} \equiv (\gamma_{\ell_\mu h}^{\ell_\tau h'} + \gamma_{\ell_\mu \bar{h}}^{\ell_\tau \bar{h}} + \gamma_{h \bar{h}}^{\ell_\tau \ell_\mu}) / (K_{\mu 2} K_{\tau 2})$ . We have taken  $M_2 = 10^7$  GeV and  $(\lambda^\dagger \lambda)_{22} = 10^{-4} \simeq h_\tau^2$  (with  $h_\tau$  the Yukawa coupling of the  $\tau$ ). Note that the interactions involving a real  $N_2$  scale as  $(\lambda^\dagger \lambda)_{22}$ , while the FCI are proportional to  $(\lambda^\dagger \lambda)_{22}^2$ . The corresponding curves for other values of  $M_2$  can be obtained simply by making the appropriate translation along the  $T$ -axis. The spiky shape of the green curve for  $0.1M_2 \lesssim T \lesssim M_2$  is due to the subtraction of the on-shell contribution to the process  $\ell_\tau h \rightarrow \ell_\mu h$ .

the two flavour case (this happens when  $K_{\alpha 1} \tilde{m}_1 \lesssim m_*$  for two different flavours). On the other hand, it can be shown that for a given value of  $M_2/M_1$ , the maximum value of  $Y_B^f$  is proportional to  $\sim \sqrt{M_1}$  in the relevant region of the parameter space. Hence an increase in  $Y_B^f$  by a factor 2 leads to a decrease in the lower bound of  $M_1$  by a factor 4.

Finally let us comment on the relation between the parameters defined above and the light neutrino masses. To lowest order in the  $L$ -violating parameters, the light neutrino masses  $m_\nu$  are given by

$$(m_\nu)_{\alpha\beta} \sim \lambda_{\alpha 1} \frac{v^2}{M_1} \lambda_{\beta 1} + (\lambda'_{\alpha 2} - \frac{\mu_2}{M_2} \lambda_{\alpha 2}) \frac{v^2}{M_2} \lambda_{\beta 2} + \lambda_{\alpha 2} \frac{v^2}{M_2} (\lambda'_{\beta 2} - \frac{\mu_2}{M_2} \lambda_{\beta 2}) \quad (4.1)$$

where  $v = \langle h \rangle = 174$  GeV is the vev of the Higgs field.



**Figure 3:** The lower bound on the reheating temperature as a function of  $M_2/M_1$ . The red curve is for the case  $\mu_2 \gg \Gamma_{N_{2l,2h}}$  and the green one for  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ .

We have seen that  $M_1$  is minimized for values of  $\tilde{m}_1 \equiv (\lambda^\dagger \lambda)_{11} v^2 / M_1$  in the range  $10^{-2} \text{ eV} \lesssim \tilde{m}_1 \lesssim 10^{-1} \text{ eV}$ , therefore the contribution of  $N_1$  to light neutrino masses is expected to be of the same order, barring cancellations due to phases. Imposing that such contributions are of order  $m_{atm} \sim 0.05 \text{ eV}$ , we get  $\lambda_{\alpha 1} \sim 10^{-5} - 10^{-4}$ .

To reproduce the observed solar and atmospheric mass scales, at least one of the contributions from  $N_2$  in (4.1) should be also of order  $10^{-2} \text{ eV}$ ; for the parameters that minimize  $M_1$ , this requirement leads to  $\mu_2/M_2 \sim 10^{-8} - 10^{-6}$ , independently of the ratio  $M_2/M_1$ . Moreover,  $\Gamma_{N_{2l,2h}}/M_2 \sim 5 \times (10^{-4} - 10^{-2})$ , therefore typically  $\mu_2 \ll \Gamma_{N_{2l,2h}}$ . However, for  $M_1 \gtrsim 5 \times 10^6 \text{ GeV}$ , and still not considering large fine tunings related to phase cancellations, smaller values of  $\lambda_{\alpha 2}$  can lead to successful leptogenesis, and in this region it is possible to have  $\mu_2 \gtrsim \Gamma_{N_{2l,2h}}$ .

With respect to the  $L$ -violating parameters  $\lambda'_{\alpha 2}$ , their contribution to the masses of the light neutrinos is  $m_\nu \sim m_{atm}$  typically for  $\lambda'_{\alpha 2} \sim 10^{-8} - 10^{-7}$ . We have checked for consistency that these small values give a negligible contribution to leptogenesis (more specifically to the CP asymmetries and washouts).

## 5. Conclusions

We have studied leptogenesis in the framework of the seesaw mechanism with small violation of  $B - L$ . If  $B - L$  is only slightly broken, then either the heavy neutrinos which generate the baryon asymmetry are almost degenerate and combine to form quasi-Dirac fermions, which can have large lepton number conserving Yukawa couplings, or they are Majorana fermions with small lepton number violating Yukawa couplings.

In both cases there are interesting consequences for leptogenesis: in the first one, the strong degeneracy of the heavy neutrinos leads to a resonant enhancement of the CP asymmetry, avoiding the DI bound on  $M_1$  which applies to hierarchical SM singlets. In the second case, the  $L$ -conserving part of the flavoured CP asymmetries in  $N_1$  decays can be much larger than the  $L$ -violating one, since the former is not linked to light neutrino masses and also escapes the DI bound, even if the heavy neutrinos are hierarchical. As a consequence, in models with almost conserved  $B - L$  successful leptogenesis may be possible at lower temperatures than in the standard seesaw, alleviating the gravitino problem in supersymmetric scenarios.

In this paper we have focused on the second possibility, i.e., we have not resorted to the resonant enhancement of the CP asymmetry. Since the  $L$ -conserving contributions to the flavoured CP asymmetries cancel in the total CP asymmetry,  $\epsilon_i = \sum_{\alpha} \epsilon_{i\alpha}$ , it is mandatory that flavour effects are at work for these terms to have an impact in leptogenesis.

We have exhaustively scanned the parameter space of seesaw models with almost conserved  $B - L$ , in which the  $L$ -conserving piece of the flavoured CP asymmetries dominates (purely flavoured leptogenesis). We have found that the largest baryon asymmetry is generated by the lightest SM singlet, and it rapidly decreases if  $|M_2 - M_1| \lesssim 2M_1$  (but far from the resonance region,  $|M_2 - M_1| \sim \Gamma_2$ ), being  $N_2$  the next-to-lightest heavy neutrino; thus we have restricted the heavy neutrino masses to the region  $M_2 \gtrsim 2M_1$ . The relevant parameters are  $M_1$ ,  $M_2/M_1$ ,  $(\lambda^\dagger\lambda)_{11}$ ,  $(\lambda^\dagger\lambda)_{22}$  and the flavour projectors  $K_{\alpha i}$ . The same  $N_2$ -Yukawa couplings which enhance the  $L$ -conserving CP asymmetries induce large FCI mediated by  $N_2$ , which tend to equilibrate the asymmetries in the different lepton flavours diminishing the total lepton asymmetry, especially if both  $N_1$  and  $N_2$  have similar masses and are simultaneously present in the thermal bath. Therefore, for each value of  $M_2/M_1$  we have determined the minimum  $M_1$  compatible with successful leptogenesis, maximizing the final baryon asymmetry over the remaining parameters. We have solved numerically the relevant set of BE, including decays and inverse decays of the two singlet neutrino species, as well as the FCI. We have considered two possibilities:  $N_2$  is a pseudo-Dirac fermion (i.e., two Majorana neutrinos with masses  $M_2 \pm \mu_2$ ,  $\mu_2 \ll M_2$ ) and  $N_2$  is approximately Dirac, in which case an asymmetry between  $N_2$  and  $\bar{N}_2$  is generated and should be taken into account. In both cases we have found that leptogenesis is possible for  $M_1 \gtrsim 10^6$  GeV, as long as  $M_2/M_1 \gtrsim 5$  and  $(\lambda^\dagger\lambda)_{22} \sim 0.01 - 1$  (see Fig. 1). So purely flavoured leptogenesis in seesaw models with slightly broken  $B - L$  provides a solution to the conflict between the upper bound on  $T_{RH}$  required to solve the gravitino problem of supersymmetric scenarios and the lower bound on  $T_{RH}$  needed for successful thermal leptogenesis. However, such heavy neutrinos are far outside the reach of present and near future colliders and do not

lead to observable lepton flavour violation in non-supersymmetric frameworks.

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