Mass bounds for baryogenesis from particle decays and the inert doublet model

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Abstract. In models for thermal baryogenesis from particle decays, the mass of the decaying particle is typically many orders of magnitude above the TeV scale. We will discuss different ways to lower the energy scale of baryogenesis and present the corresponding lower bounds on the particle's mass. This is done specifically for the inert doublet model with heavy Majorana neutrinos and then we indicate how to extrapolate the results to other scenarios. We also revisit the question of whether or not dark matter, neutrino masses, and the cosmic baryon asymmetry can be explained simultaneously at low energies in the inert doublet model.

Contents

1 Introduction

There is plenty of evidence that the Universe is asymmetric, in the sense that the amount of antimatter is negligible compared to the amount of matter. This is known at different scales from different data. Notably, in [1] it has been shown that if there were matter and antimatter domains, their size should have to be larger than the observable Universe.

In addition, the size of the asymmetry is known with remarkable precision from two different ways: (a) big bang nucleosynthesis, given that the abundance of the light elements depends mainly on the baryon to photon density ratio, and (b) the anisotropies of the cosmic microwave background radiation. The most precise value comes from Planck [2], $Y_B - Y_{\bar{B}} = Y_B \simeq 8.6 \times 10^{-11}$, where $Y_B \equiv n_B/s$ is the baryon density normalized to the entropy density. This value is too large to be explained in the Standard Model (SM), therefore the Baryon Asymmetry of the Universe (BAU) is an evidence of physics beyond the SM.

The basic requirements for generating dynamically the BAU have been established long ago [3] and they have been implemented in many different scenarios (for some reviews see e.g. [4–6]). One of the simplest and most studied explanations is thermal Baryogenesis from heavy Particle Decays (BarPaDe) [7]. In particular, leptogenesis [8] models are very attractive because they can also provide an explanation to neutrinos masses, relating closely both mysteries.

However, the natural scale for these scenarios — given by the mass of the heavy particle — is above $\sim 10^8$ GeV. To illustrate this point let us take as an example type I leptogenesis (for a detailed review see e.g. [9]), with the lepton asymmetry being produced in the decay of the Majorana neutrinos N_1, and the virtual contribution to the CP asymmetry coming from another neutrino specie N_2. The final BAU can be written roughly as $Y_B^f \sim \epsilon\eta/725$, where ϵ is the CP asymmetry in the neutrino decay, η is the efficiency, with $|\eta| \leq 1$, and the numerical factor is the product of the equilibrium density of a relativistic specie of neutrino per unit entropy density and the fraction of lepton (L) asymmetry transformed into baryon (B) asymmetry by the sphalerons. On one hand, the CP asymmetry is proportional to the
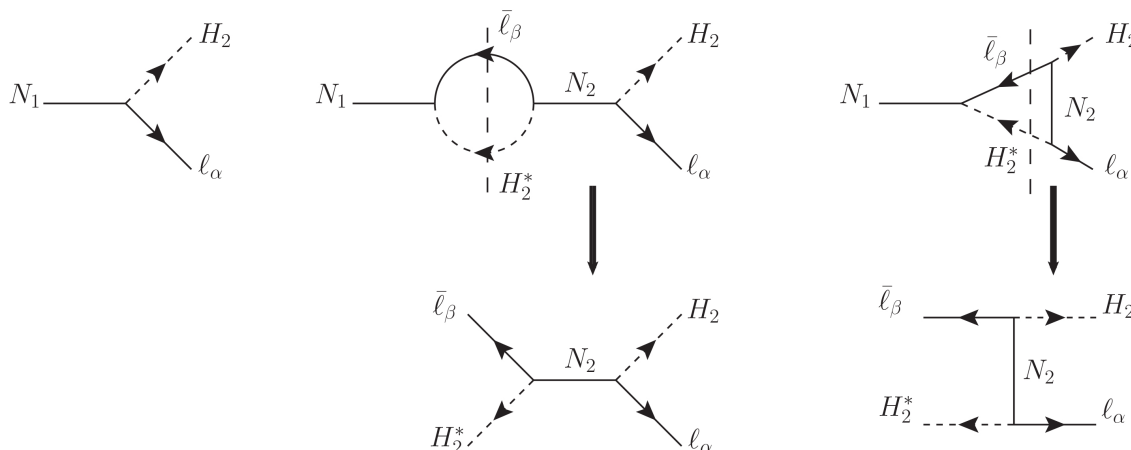
Figure 1. On the top we show a combination of tree-level and one-loop diagrams that lead to a CP asymmetry in N_1 decays. The vertical cuts through the loops indicate that those particles can go on shell, which is necessary to get a kinematical phase. In turn, this implies the existence of the tree level diagrams on the bottom, which contribute to the washout.

sum of the modulus squared of the Yukawa couplings of N_2, $(\lambda^\dagger \lambda)_{22}$, and is suppressed by the ratio of the masses, M_1/M_2 (see [10] or eq. (A.3)). Taking into account also the loop suppression factors, we have that $Y_B^f \lesssim 8 \times 10^{-5} (\lambda^\dagger \lambda)_{22}$ (for M_2 not very different from M_1, if not the bound would be smaller). Therefore $(\lambda^\dagger \lambda)_{22} \gtrsim 10^{-6}$ to obtain the BAU. On the other hand, the efficiency depends on $\tilde{m}_1 \equiv (\lambda^\dagger \lambda)_{11} v^2/M_1$, with $v = 174\,\text{GeV}$ the vev of the Higgs. This baryogenesis mechanism works better for $\tilde{m}_1 \sim 10^{-3}\,\text{eV}$, which is the condition for N_1 to decay barely out of equilibrium, yielding $\eta \sim 1$. If we assume that the Yukawa couplings of N_1 and N_2 are not very different, i.e. $(\lambda^\dagger \lambda)_{11} \sim O\left(10^{-6}\right)$, then the scale for baryogenesis is $M_1 \sim 10^{10}\,\text{GeV}$.

Moreover, it is not only that the natural scale for BarPaDe is very high, but also there are severe problems to have successful baryogenesis at lower scales. One of these problems appears in some of the most studied leptogenesis models, which yield a relation between the BAU and light neutrino masses (LNM). It turns out that the CP asymmetry is suppressed by the tiny value of the LNM. E.g. in type I leptogenesis with hierarchical heavy neutrinos, the mass of the neutrino producing the asymmetry must be $M_1 \gtrsim 10^8 - 10^9\,\text{GeV}$ [11, 12]. But the main problem is not related at all with neutrino masses and is common to all models for BarPaDe. It stems from the fact that CP violation requires not only a complex phase in the couplings, but also a kinematical phase. This one in turn implies the existence of on-shell processes that violate L or B (see figure 1) and tend to washout the $B - L$ asymmetry [13]. Since the CP asymmetry is proportional to the couplings of this L/B-violating interactions, they cannot be very small, which leads to washout processes that are typically too fast compared to the expansion rate of the Universe if baryogenesis occurs at low temperatures (when the expansion rate is small).

Without forgetting that the natural scale for BarPaDe seems to be very high, $M \gtrsim 10^8\,\text{GeV}$, it is really worth to study ways to have BarPaDe at lower scales, so that some of these scenarios become more easily accessible to experimental exploration. Moreover, some well motivated supergravity models require reheating temperatures below $10^5 - 10^7\,\text{GeV}$ to avoid the gravitino problem [14, 15], and hence baryogenesis must occur at quite "low" scales. There are also interesting models of physics beyond the SM incorporating heavy particles at

the TeV scale, therefore it is worth knowing generic mechanisms which can yield successful baryogenesis from the decay of these particles. Finally, it could also happen that baryogenesis above a few TeV becomes severely disfavored, demanding a low-scale explanation for the BAU, e.g. if some lepton number violating processes are observed at the LHC [16, 17].

The paper is organized as follows. In section 2 we analyze different mechanisms for BarPaDe and find mass bounds for them to work. We do so in the framework of the inert doublet model, while in section 3 we discuss the model dependence of the results. In section 4 we revisit the interesting question of whether LNM, Dark Matter (DM) and the BAU can all be explained at low energies in the inert doublet model without resorting to degenerate heavy neutrinos. Finally we summarize in section 5 and give the Boltzmann equations (BEs) and other relevant expressions in the appendix.

2 Mechanisms for baryogenesis from particle decays

We will explore different mechanisms of thermal BarPaDe, indicating the main requirements for them to work, in particular the lower bound on the particle's mass. We will assume that baryogenesis occurs in the radiation-dominated era and that no significant amount of entropy is created during or after the BAU is generated. To be more specific a particular model must be chosen. An attractive and simple one, that can encompass the different variants for baryogenesis we wish to study, is the inert doublet model complemented with — at least two — heavy Majorana neutrinos (hereafter referred as IDM) [18]. The singlet neutrinos N_i participating in leptogenesis are taken heavier than the inert doublet H_2, $M_i > M_{H_2}$. Both N_i and H_2 are — the only — odd particles under a discrete Z_2, therefore H_2, being the lightest, is a candidate for DM. The lagrangian of the model together with the BEs and other relevant expressions can be found in the appendix. More details on the IDM and references are given in section 4.

Quite generally, the major contribution to the lepton asymmetry comes from the decay of one of the neutrinos, that will be called N_1. For simplicity all the calculations in this section are made assuming that there is only one more neutrino specie, N_2, that contributes virtually to the CP asymmetries in N_1 decays (see figure 1). The parameters relevant for leptogenesis are the masses M_1, M_2, M_{H_2}, the Yukawa couplings $\{\lambda_{\alpha i}, \alpha = e, \mu, \tau; i = 1, 2\}$, and the temperature at which sphalerons decouple (see below). In sections 2.1, 2.3, and 2.4 we will give the lower bound on M_1 for successful leptogenesis, $M_{1\text{min}}$, obtained after scanning over all the parameter space of Yukawa couplings (whose phases are chosen to maximize the CP asymmetries, see the appendix for more details). In addition we will indicate the dependence of $M_{1\text{min}}$ on M_2/M_1 and M_{H_2}/M_1. In turn, for almost degenerate neutrinos leptogenesis is possible for very low masses, as long as $(M_2 - M_1)/M_1$ is sufficiently tiny. Then, in this *resonant* scenario it is more interesting to ask how small $(M_2 - M_1)/M_1$ needs to be to generate the BAU for a given M_1. We study this in section 2.2.

It is good to have a rough idea of the role of the different Yukawa couplings. A convenient way of expressing them is in terms of the combinations $(\lambda^\dagger \lambda)_{11}, (\lambda^\dagger \lambda)_{22}$, the projectors $K_{\alpha i} \equiv \frac{\lambda_{\alpha i} \lambda_{\alpha i}^*}{(\lambda^\dagger \lambda)_{ii}}$ and some phases. The strength of the Yukawa interactions of N_1, measured by the effective mass $\tilde{m}_1 \equiv (\lambda^\dagger \lambda)_{11} v^2 / M_1$, determines (i) the washouts induced by the inverse decays of N_1, which turn out to be relevant only for the bound determined in section 2.1, (ii) the amount of neutrinos that can be produced if no interaction other than the Yukawa creates the N_1, and (iii) the lifetime of N_1, which is an important quantity in sections 2.3 and 2.4, since a late decay of the N_1 allows to avoid potentially huge washouts. The reference value

for $\tilde{m}$ is the equilibrium mass $m_* \simeq 10^{-12}\,\mathrm{GeV}$, defined by the condition $\frac{\Gamma_1}{H(T=M_1)} = \frac{\tilde{m}_1}{m_*}$, where Γ_1 is the decay width of N_1, and H is the Hubble rate. Note that for $\tilde{m}_1$ to be equal to m_* at the TeV scale, $(\lambda^\dagger\lambda)_{11}$ has to be really small, e.g. $(\lambda^\dagger\lambda)_{11} \approx 3 \times 10^{-14}$ for $M_1 = 1\,\mathrm{TeV}$.

In turn, the Yukawa couplings of N_2 — together with the ratio M_2/M_1 — determine the size of the CP asymmetry and the intensity of the $\Delta L = 2$ washouts mediated by N_2, both quantities intimately related, as explained before and represented in figure 1. For $M_2/M_1 \lesssim 6$ it is also very important to take into account the washouts due to the inverse decays of N_2. Except in the resonant case, $(\lambda^\dagger\lambda)_{22} \gtrsim 10^{-6}$ in order to have enough CP violation, with the optimal value depending on M_2/M_1. Since in most cases described below the optimum value of $\tilde{m}_1 \lesssim m_*$ (and often much smaller), a large hierarchy among the Yukawa couplings of N_1 and N_2 is usually required to achieve baryogenesis at the TeV scale (in the resonant scenario this can be avoided by taking very tiny values of $(M_2 - M_1)/M_1$). We will comment more on this issue in the following sections.

Before going to the particular analysis of each mechanism, we point out some common technical issues and details.

- For the observed BAU we take the value $Y_B \simeq 8.6 \times 10^{-11}$ [2], without considering observational errors, which are negligible for the purpose of this work.

- We have imposed that $(\lambda^\dagger\lambda)_{11}, (\lambda^\dagger\lambda)_{22} \le 10$ in order to safely satisfy the perturbative bound on the Yukawa couplings. In fact that bound is somewhat larger, $|\lambda_{\alpha i}| < 4\pi$, but most of the results do not depend on the exact upper value imposed (we will indicate when this is not the case).

- Except in the almost degenerate case, we have taken $M_2/M_1 \ge 3$ to avoid enhancements of the CP asymmetry due to the proximity of M_2 and M_1. In this way the distinction between the resonant mechanism and the other ones described in the text will be clear.

- In all the situations of interest, there are very strong washouts at $T > M_1$, so any asymmetry that could be produced at $T \gtrsim M_1$ is quickly erased. This renders the results more independent of the initial conditions and quite insensitive to scatterings and finite temperature effects [19–22]. As we will see, the only initial condition that is relevant is the amount of neutrinos at the onset of leptogenesis.

- The scenarios we consider occur at low temperatures, when all the SM Yukawa interactions are in equilibrium. Then the BEs are diagonal in the lepton flavor basis $\{e, \mu, \tau\}$ and leptogenesis occurs in the three flavor regime [20, 23–28]. Actually, if the Yukawa interactions of N_2 are very large, the basis that diagonalizes the leptonic density matrix may be different from $\{e, \mu, \tau\}$, and moreover it can change with the temperature. However, the maximum baryon asymmetry is obtained in cases where the interactions of N_2 are not super-fast at the stage the BAU is created, so that a set of BEs diagonal in $\{e, \mu, \tau\}$ can be safely used.

- In all the mechanisms described below, we have found that the maximum baryon asymmetry is obtained when the projectors of N_1, $K_{\alpha 1}$, are equal to $1/3$ (to a good degree of accuracy). And in all but the resonant case, the maximum is obtained also for $K_{\alpha 2} = 1/3$. The reason is due to: (i) The maximum of the CP asymmetries with respect to the projectors, taking into account the constrain $\sum_\alpha K_{\alpha i} = 1$, is obtained for the symmetric case $K_{\alpha 1} = K_{\alpha 2} = 1/3$. (ii) The CP asymmetries and $\Delta L = 2$ washouts

mediated by N_2 are proportional to some power of a product of projectors of N_2 with $(\lambda^\dagger \lambda)_{22}$. Therefore it is enough to maximize the B-asymmetry yield over $(\lambda^\dagger \lambda)_{22}$ if the $\Delta L = 2$ washouts are dominant, which is usually the case. The exception is for almost degenerate neutrinos, when the $\Delta L = 1$ washouts due to the inverse decays of N_2 may be very important. In that case it is convenient to take some projectors very small, since the projector of a given flavor suppresses more the $\Delta L = 1$ washouts than the CP asymmetries (contrary to $(\lambda^\dagger \lambda)_{22}$ that enters linearly in both quantities). In any case we stress that the bounds given in the following sections have been obtained scanning over all the allowed range for the projectors, $0 \leq K_{\alpha 1}, K_{\alpha 2} \leq 1$ (except for figure 3 which explicitly deals with some restrictions on the couplings).

- The freeze out temperature of the sphalerons, $T_{\rm sfo}$, is an important quantity for leptogenesis models. Its value depends on the critical temperature of the electroweak phase transition, T_c, and on whether this is a first or second order transition. In the SM, $T_c \sim 140\,{\rm GeV}$ and $T_{\rm sfo} \sim T_c/1.7$ [29–32]. To show how results depend on $T_{\rm sfo}$ (that may change from one model to another), we have worked with two different values of $T_{\rm sfo}$, $80\,{\rm GeV}$ and $140\,{\rm GeV}$. We will also comment on bounds for models with perturbative violation of B, so that baryogenesis may occur at $T \ll T_{\rm sfo}$.

2.1 Massless decay products with zero initial density

We start considering the case with non-degenerate neutrinos, that can be produced only by the CP-violating Yukawa interactions, and decay into particles with negligible mass. In figure 2 we show the lower bound on M_1 as a function of the mass hierarchy among the neutrinos, M_2/M_1. The main conclusion drawn from this plot is that successful baryogenesis, under the conditions stated above, requires a mass for the decaying particle above $\sim 80\,{\rm TeV}$. This bound is not related at all to the small values of LNM, but to the washout processes represented in figure 1, and is therefore inherent to this baryogenesis mechanism. On the contrary, the bound $M_1 \gtrsim 10^8 - 10^9\,{\rm GeV}$ in standard type I leptogenesis [11, 12] is due to the connection with LNM and their tiny value.

Furthermore, baryogenesis at this — quite low — scale requires a large hierarchy among the Yukawa couplings of N_1 and N_2. E.g., when $M_2/M_1 = 10$ the lower bound is obtained for $\tilde{m}_1 \approx 10^{-12}\,{\rm GeV}$ ($(\lambda^\dagger \lambda)_{11} \approx 3 \times 10^{-12}$) and $(\lambda^\dagger \lambda)_{22} \approx 5 \times 10^{-5}$, i.e. the Yukawa couplings of N_2 must be 3 to 4 orders of magnitude larger than those of N_1. This holds true for all the values of M_2/M_1 represented in the plot.

Next we give some more details on the optimum values of the Yukawa couplings to get the maximum baryon asymmetry, distinguishing between cases with large and small values of M_2/M_1.

- $M_2/M_1 \gtrsim 20$: when $M_2/M_1 \gg 1$ the washouts involving an on-shell N_2 are irrelevant. In turn, the CP asymmetry and the washouts mediated by N_2 are proportional to powers of the combination $(\lambda^\dagger \lambda)_{22}/(M_2/M_1)$. Hence an increase in M_2/M_1 by a factor a is completely compensated by the same increase in $(\lambda^\dagger \lambda)_{22}$. As a consequence, the lower bound on M_1 stays the same for all $M_2/M_1 \gg 1$ after maximizing the BAU over the Yukawa couplings.[1] Furthermore, since the optimum value of $\tilde{m}_1$ keeps constant at $\sim 1 - 2 \times 10^{-12}\,{\rm GeV}$, the required hierarchy among couplings to get the maximum BAU grows mildly with M_2/M_1, $\sqrt{(\lambda^\dagger \lambda)_{22}/(\lambda^\dagger \lambda)_{11}} \propto \sqrt{M_2/M_1}$.

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[1]For $M_2/M_1 \gtrsim 10^6$ the optimum values of the Yukawa couplings become non-perturbative and this analysis is no longer valid.
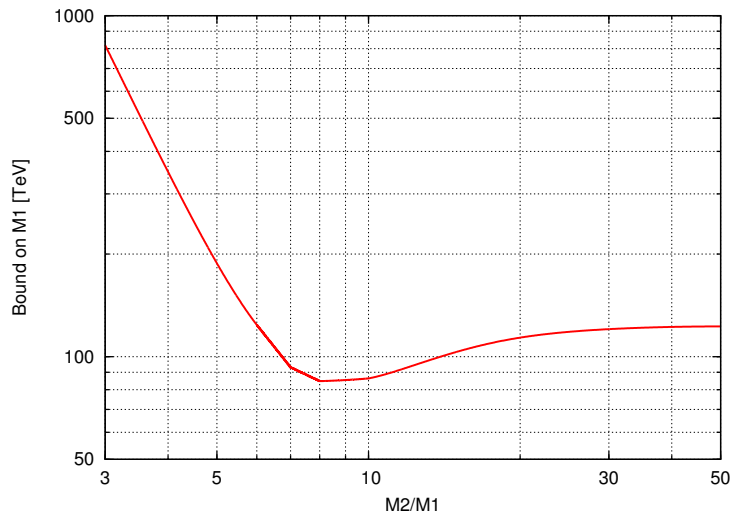
Figure 2. Lowest value of M_1 yielding successful leptogenesis as a function of M_2/M_1. This is for massless decay products and a null density of N_1 at the start of leptogenesis. The bound for each value of M_2/M_1 has been obtained scanning over all the parameter space of — perturbative — Yukawa couplings.

- $M_2/M_1 \lesssim 6$: the optimum value of $\tilde{m}_1$ decreases as M_2 approaches M_1. The reason is that the smaller the $\tilde{m}_1$, the later the neutrinos decay and start to generate the BAU, avoiding — at least partially — the dangerous washouts due to on-shell N_2 (the drawback is that the maximum amount of N_1 that can be produced decreases with $\tilde{m}_1$). However the lower bound on M_1 increases as M_2/M_1 decreases, so that the optimum value of $(\lambda^\dagger \lambda)_{11}$ does not change much, as neither does $(\lambda^\dagger \lambda)_{22}$.

2.2 Almost degenerate neutrinos

When N_1 and N_2 are almost degenerate, the CP asymmetry can be enhanced up to $O(1)$ values [33–36]. More precisely, for $\Gamma_{1,2} \ll \Delta M \ll M_1$ the CP asymmetries $\epsilon_{\alpha 1} \propto (\lambda^\dagger \lambda)_{22}/\delta$, with $\delta \equiv \Delta M/M_1$, and $\Delta M \equiv M_2 - M_1$.[2] Therefore it is possible to reduce the washouts taking $(\lambda^\dagger \lambda)_{22}$ small enough, while keeping $\epsilon_{\alpha 1}$ sizeable by choosing a tiny value for δ. This — so called resonant leptogenesis — mechanism has been widely studied (see e.g. [25, 37–39], and [40] for a model with conservation of $B-L$) and it is known that it can lead to successful leptogenesis at the TeV scale.

We want to determine how small the degeneracy parameter δ has to be to get the BAU. For this, we maximize the baryon asymmetry over all the Yukawa couplings for different values of M_1. Since the neutrinos have very similar masses, it is to expect that their Yukawa couplings are not very different. To include this aspect in the analysis we impose different restrictions on the ratio of the Yukawa couplings. For this it is convenient to define the parameter r as the minimum non-null quantity in the following list, $\left\{ \sqrt{(\lambda^\dagger \lambda)_{11}/(\lambda^\dagger \lambda)_{22}}, \sqrt{K_{\alpha i}} \ (i = 1, 2; \ \alpha = e, \mu, \tau) \right\}$. The results are shown in figure 3, where the upper bound on δ is represented as a function of M_1 for three different set of restrictions on the couplings. The upper (red) curve has been obtained imposing that $r \geq 10^{-5}$, i.e. the Yukawa couplings were allowed to differ in 5 orders of magnitude or less. In turn, the intermediate (green) and lower (blue) curves correspond to $r \geq 10^{-3}$ and $r \geq 10^{-1}$, respectively.

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[2]We will not include in the analysis the much more involved case with maximum enhancement of the CP asymmetry, that arises when $\Delta M \sim \Gamma_{1,2}$.
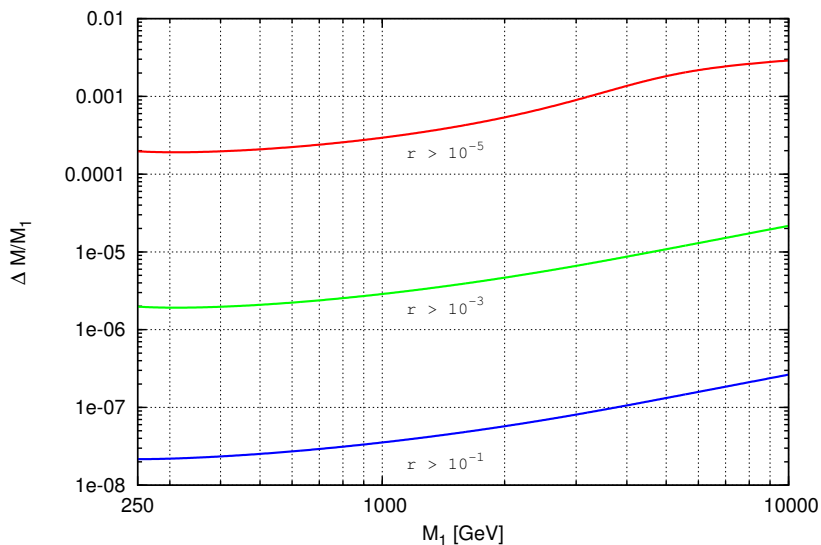
Figure 3. Upper bound on δ as a function of M_1 for three different constraints on the hierarchy among the Yukawa couplings, $r \geq 10^{-5}, 10^{-3}$, and 10^{-1}.

From the figure one can see that, e.g., to have successful leptogenesis for some $M_1 \lesssim 4\,\text{TeV}$, δ has to be lower than $\sim 10^{-7}$ if the Yukawa couplings differ from one another in less than one order of magnitude. More generally, the upper bound on $\delta \propto 1/r$. The reason is, roughly, that the CP asymmetry is $\propto 1/\delta$ and the washouts are independent of δ (if $\delta \ll 1$), while the ratio between the CP asymmetry and the washouts is $\propto 1/r$. Altogether this implies that the final baryon asymmetry $Y_B^f \propto 1/(r\,\delta)$ [41]. As a rule of thumb, to get the BAU via the resonant leptogenesis mechanism requires that

$$
\begin{aligned}
\delta \times r &\lesssim 10^{-8}\,, && \text{for} && M_1 \sim 4\,\text{TeV} && \text{and} \\
\delta \times r &\lesssim 2 \times 10^{-9}\,, && \text{for} && 250\,\text{GeV} \lesssim M_1 \lesssim 1\,\text{TeV}\,.
\end{aligned}
\tag{2.1}
$$

Some comments are in order:

- For figure 3 we have taken $T_{\text{sfo}} = 140\,\text{GeV}$, but the conclusions stated above are quite independent of the exact value of T_{sfo}. We have verified that even in models for baryogenesis with perturbative violation of B (so that Y_B is not frozen at T_{sfo}), and with similar particle content than the IDM, the rules of thumb stated above are correct within a factor $\lesssim 5$.

- For $M_1 \lesssim T_{\text{sfo}}$ leptogenesis occurs at high temperatures, $T \gtrsim M_1$, in which case scatterings and finite temperature effects [19–22] become very important and the BEs we have used are no longer appropriate. That is why we have taken $M_1 \gtrsim 250\,\text{GeV}$ in figure 3 and rule (2.1).

- Moreover the BEs used in this work are classical and therefore they are not adequate to treat the resonant limit $\Delta M \sim \Gamma_2/2$. Nevertheless, for all the points plotted in figure 3 it is possible to obtain the BAU with $\Delta M \gg \Gamma_{1,2}$, in which case a quantum treatment [42–46] is not mandatory to get a quite accurate result. To be more concrete, for the three curves in figure 3, $r \geq 10^{-5}, 10^{-3}$, and 10^{-1}, it is verified that $\Delta M/\Gamma_{1,2} \gtrsim 5 \times 10^0, 10^2$, and 10^4, respectively.
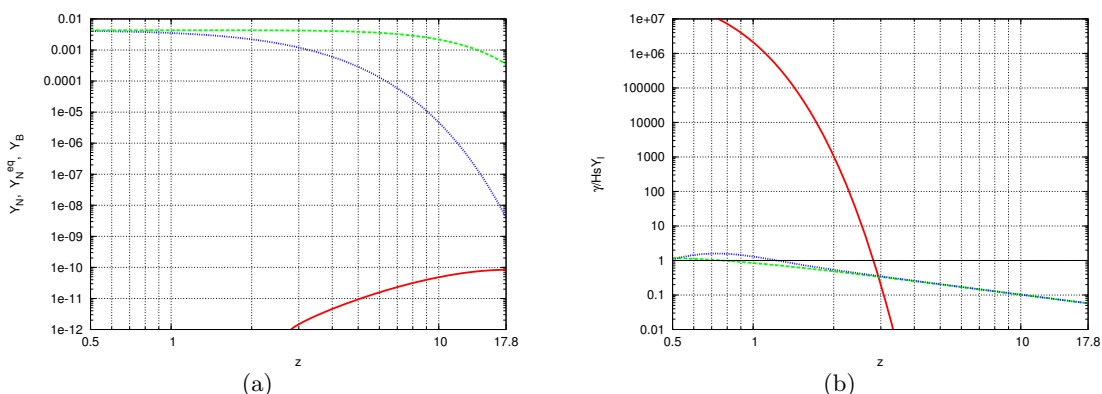
Figure 4. Figure 4(a): the evolution of the equilibrium density of N_1, Y_N^{eq} (dotted blue line), the real density Y_N (dashed green line), and the baryon asymmetry $|Y_B|$ (solid red line), as a function of z. Here Y_X denotes the number density of X normalized to the entropy density. Figure 4(b): the reaction densities $\gamma_{N_2}^{\ell H_2}$ (solid red line), $\gamma_{\bar\ell \bar H_2}^{\ell H_2}{}'$ (dotted blue line), and $\gamma_{\bar H_2 \bar H_2}^{\ell\ell}$ (dashed green line), normalized to $H n_\ell^{\text{eq}}$, with $\gamma_{c,d,\dots}^{a,b,\dots} \equiv \gamma(a,b,\dots \to c,d,\dots)$, and n_ℓ^{eq} the equilibrium number density of a relativistic particle with one degree of freedom. The prime in $\gamma_{\bar\ell \bar H_2}^{\ell H_2}{}'$ indicates that the on-shell contribution has been subtracted. The parameters are: $T_{\text{sfo}} = 140\,\text{GeV}$, $M_1 = 2500\,\text{GeV}$, $M_2 = 10 M_1$, $\tilde m_1 = 2 \times 10^{-5}\,\text{eV}$, $(\lambda^\dagger \lambda)_{22} = 2 \times 10^{-5}$, and all projectors are taken equal to 1/3. The Yukawa couplings of N_1 are chosen to let most of the neutrinos decay just before the sphalerons freeze out at $T \sim 140\,\text{GeV}$ ($z \sim 17.8$), while those of N_2 are the largest ones that do not induce a significant washout of the asymmetry. Note that the baryon asymmetry Y_B remains constant for $z \gtrsim 17.8$, since the sphaleron processes are then out of equilibrium.

2.3 Massless decay products with initial thermal density

Another mechanism to have baryogenesis at the TeV scale arises when the N_1 are produced at higher temperatures by a process different from the CP-violating Yukawa interactions. Since there is another interaction that produces the neutrinos, the Yukawa couplings $\lambda_{\alpha 1}$ can be chosen small enough so that the N_1 decay at $T \ll M_2$. In this way it is possible to have large couplings of N_2 and consequently a big CP asymmetry, but at the same time small washouts at the moment the N_1 start to decay and produce the BAU. It is also crucial that the interaction responsible for the creation of the N_1 decouples before they decay.

This mechanism is illustrated in figure 4, where the relevant densities and rates are plotted as a function of $z \equiv M_1/T$, for $M_1 = 2500\,\text{GeV}$.

The lower bound on M_1 is $M_{1\text{min}} \sim 2500\,(2000)\,\text{GeV}$ for $T_{\text{sfo}} = 140\,(80)\,\text{GeV}$. Bounds of the same order are to be expected in all models for leptogenesis, which require that the lepton asymmetry be generated at $T > T_{\text{sfo}}$. Instead, if B is violated perturbatively, we have verified that baryogenesis can occur, *at least in principle*, at any temperature above the matter-antimatter annihilation epoch, i.e. at $T \gtrsim 40\,\text{MeV}$ (see [47] and also [48]), which implies a lower bound on the mass of the decaying particle of $O(1)\,\text{GeV}$.[3] Moreover, the lower bound can be achieved for a wide range of values of M_2/M_1, with the optimal value of $(\lambda^\dagger \lambda)_{22}$ being proportional to M_2/M_1.

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[3]Models for baryogenesis below the TeV scale are very constrained by experiments (see e.g. [49]). It is not the goal of this work to analyze these issues, the claim is just that baryogenesis could in principle work at very low energies via the mechanism described in this and other sections of the paper.
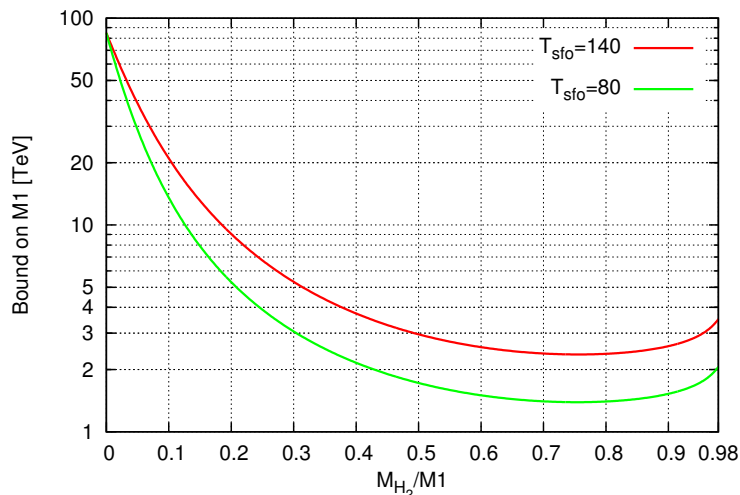
Figure 5. Lowest value of M_1 yielding successful leptogenesis as a function of M_{H_2}/M_1 for $T_{\rm sfo} = 80$ and $140\,{\rm GeV}$. The bound for each value of M_{H_2}/M_1 has been obtained scanning over all the parameter space of — perturbative — Yukawa couplings and over $M_2/M_1 \geq 3$.

We note that one possibility to produce enough neutrinos with very small Yukawa couplings is via some exotic gauge interaction, as in [50, 51]. It is important that the Z' gauge boson be heavy enough, so that this interaction is out of equilibrium at $T \lesssim M_1$.

2.4 Massive decay products

The authors of [52] showed that the BAU can be generated at low energies from DM annihilations if the DM annihilates into a SM baryon or lepton plus a massive exotic particle Ψ (see also [53]). When the mass of Ψ, m_Ψ, is of the order of the DM mass, $m_{\rm DM}$, the dangerous ΔB or $\Delta L = 2$ washouts become Boltzmann suppressed by the factor $e^{-m_\Psi/T}$. In this section we want to study the same suppression mechanism, but for baryogenesis from heavy particle decays instead of annihilations. We stress that there is a major difference between the decay and annihilating scenarios. In the later, m_Ψ can be well above $m_{\rm DM}$ — but below $2\,m_{\rm DM}$ — without reducing significantly the phase space for the necessary B — or L — violating interactions. Instead, in the former, m_Ψ is condemned to be below the mass M_1 of the decaying particle. In fact, as m_Ψ gets closer to M_1 two opposite effects arise: on one hand the washouts become more Boltzmann suppressed, but on the other hand there is a reduction in the CP asymmetry due to the shrinkage of the available phase space for the decays.

An interesting example to study quantitatively the effect of massive decay products is the IDM. Here the neutrinos decay into a lepton and the inert doublet H_2, whose mass, M_{H_2}, is a free parameter. We have calculated the lower bound on M_1 for successful leptogenesis as a function of M_{H_2}/M_1 and we show the results in figure 5. The two curves correspond to $T_{\rm sfo} = 80$ and $140\,{\rm GeV}$.

The main result that can be drawn from the figure is that the BAU can be generated at very low scales with this mechanism. More precisely, the lower bounds on M_1 are $\sim 1300\,{\rm GeV}$ and $\sim 2300\,{\rm GeV}$, for $T_{\rm sfo} = 80\,{\rm GeV}$ and $140\,{\rm GeV}$, respectively. In fact, the lower bound is basically determined by the value of $T_{\rm sfo}$, the approximate relation being $M_{1\rm min} \sim 17 \times T_{\rm sfo}$ if $T_{\rm sfo} \sim O\left(10^2\right)\,{\rm GeV}$. The number 17 is roughly the smallest value of z at which the neutrinos

can decay and produce a lepton asymmetry that is safe from $\Delta L = 2$ washouts. On the contrary, if B is violated perturbatively, baryogenesis with massive decay products can occur — in principle — at any temperature above the matter-antimatter annihilation epoch.

It is important to remark that a large hierarchy among the Yukawa couplings of N_1 and N_2 is required to generate the BAU at such low scales, with $|\lambda_{\alpha 2}/\lambda_{\alpha 1}| \sim 10^6$ for the points in figure 5. E.g., for $T_{\mathrm{sfo}} = 140\,\mathrm{GeV}$ we have verified that if the Yukawa couplings are not allowed to differ by more that 5 (4) orders of magnitude, the lower bound on M_1 rises from $\sim 2.3\,\mathrm{TeV}$ to $\sim 9\,(25)\,\mathrm{TeV}$.

Some more comments are in order:

- From figure 5 it is clear that the massive-decay-product mechanism works efficiently for a wide range of values of M_{H_2}/M_1, e.g. $M_{1\mathrm{min}} \lesssim 3\,(5)\,\mathrm{TeV}$, for $0.3 \lesssim M_{H_2}/M_1 \lesssim 1$ and $T_{\mathrm{sfo}} = 80\,(140)\,\mathrm{GeV}$. Even for H_2 masses as low as $M_{H_2} \sim 0.1 M_1$ the lower bound is significantly reduced compared to the case with massless decay products. When M_{H_2} approaches M_1 the optimum value of $(\lambda^\dagger \lambda)_{22}$ increases (to compensate for the phase space suppression of the CP asymmetry), and hence $M_{1\mathrm{min}}$ depends on the exact upper bound imposed to $(\lambda^\dagger \lambda)_{22}$. However, we have verified that this has an effect only for $M_2 \gg M_1$ and M_{H_2} very close to M_1, $M_{H_2}/M_1 > 0.95$. As an example, for $M_{H_2}/M_1 = 0.95$ and $M_2/M_1 \lesssim 20$, $M_{1\mathrm{min}} \sim 3\,\mathrm{TeV}$ (if $T_{\mathrm{sfo}} = 140\,\mathrm{GeV}$), with the maximum B-asymmetry obtained for $(\lambda^\dagger \lambda)_{22} < 1$.

- For $0.1 \lesssim M_{H_2}/M_1 \lesssim 0.95$ the lower bound on M_1 is essentially independent of the ratio M_2/M_1 in the range $3 \lesssim M_2/M_1 \lesssim 50$. If $M_2/M_1 \lesssim 3$ the CP asymmetry starts to increase due to the proximity of M_2 and M_1, and therefore $M_{1\mathrm{min}}$ decreases. However we have not included these lower values of M_2/M_1 in the analysis in order to distinguish the mechanism of this section from that of section 2.2. Instead for $M_2/M_1 \gtrsim 50 - 100$, $M_{1\mathrm{min}}$ starts to increase due to the perturbative upper bound on $(\lambda^\dagger \lambda)_{22}$.

- The results mentioned above are valid when the neutrinos are solely produced by their Yukawa interactions, i.e. if the density of neutrinos is null at the beginning of the leptogenesis epoch. Instead, if there is a population of neutrinos with a thermal density at $T \gg M_1$, the lower bound on M_1 for successful leptogenesis is somewhat below $750\,(1200)\,\mathrm{GeV}$ for $T_{\mathrm{sfo}} = 80\,(140)\,\mathrm{GeV}$, the exact value depending on M_2/M_1. The required hierarchy among the couplings is $|\lambda_{\alpha 2}/\lambda_{\alpha 1}| \sim 10^5 - 10^6$.

- There is a crucial requirement for this mechanism to work. The issue is related to the washouts and can be understood through a simple analysis of the terms in the BEs. The most dangerous washout is (see eq. (A.2)):

$$\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} = \frac{1}{sHz} \sum_\beta (1 + \delta_{\alpha\beta}) \left(\gamma_{\bar{\ell}_\beta \bar{H}_2}^{\ell_\alpha H_2}{}' + \gamma_{\bar{H}_2 \bar{H}_2}^{\ell_\beta \ell_\alpha} \right) 2 \frac{Y_{\Delta H_2}}{Y_{H_2}^{\mathrm{eq}}} + \dots . \tag{2.2}$$

The effect of this term depends on how the asymmetry $Y_{\Delta H_2}$ is related to Y_{B-L}. If the interactions $H_2 H_2 \to H_1 H_1$ and $H_2 \bar{H}_1 \to \bar{H}_2 H_1$ are out of equilibrium (with H_1 the SM Higgs), the quantity $Y_{B-L} + Y_{\Delta H_2}$ is effectively conserved (when the coupling of the quartic term $\lambda_5 (H_1^\dagger H_2)^2 / 2$ of the scalar potential is null, $\lambda_5 = 0$, it is possible to define a conserved lepton number). In that case the washout term can be rewritten as

$$\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} = -\frac{1}{sHz} \sum_\beta 2(1 + \delta_{\alpha\beta}) \left[\frac{1}{Y_{H_2}^{\mathrm{eq}}} \left(\gamma_{\bar{\ell}_\beta \bar{H}_2}^{\ell_\alpha H_2}{}' + \gamma_{\bar{H}_2 \bar{H}_2}^{\ell_\beta \ell_\alpha} \right) \right] Y_{B-L} + \dots . \tag{2.3}$$

Since $\left[\frac{1}{Y_{H_2}^{\text{eq}}} \left(\gamma_{\bar{\ell}_\beta \bar{H}_2}^{\ell_\alpha H_2}{}' + \gamma_{\bar{H}_2 \bar{H}_2}^{\ell_\beta \ell_\alpha} \right) \right]$ is not Boltzmann suppressed, this washout badly erases the asymmetry, independently of how heavy H_2 is. Instead, if the interactions associated to λ_5 are fast, as has been assumed in this work, $Y_{\Delta H_2}$ is related to Y_{B-L} as in eq. (A.7), and the washouts are Boltzmann suppressed. This important point has also been explained in [54].

3   Model (in)dependence

In the previous section we derived bounds and conditions to have baryogenesis at low scales for different mechanisms. It was done specifically within the framework of the IDM, but the essential element behind those results, namely the washout processes inherent to the CP asymmetry (see figure 1), is common to all models for thermal BarPaDe. Then it is to expect that the basic conclusions also hold for other models. In this section we go beyond this general claim and show how to relate the bounds in different models, but still assuming that the BAU is generated in an era of radiation domination and that no significant amount of entropy is created since the baryogenesis epoch.

To start it is convenient to cast the BEs of a given model in the following form, assuming that the asymmetry is produced mainly in the decay of one neutrino specie, called N_1:

$$\frac{\mathrm{d}Y_{N_1}}{\mathrm{d}z} = a\, C(z) \left(Y_{N_1} - b\, Y_{N_1}^{\text{eq}} \right) , \tag{3.1}$$

$$\frac{\mathrm{d}Y_{B-L}}{\mathrm{d}z} = c\, \frac{\epsilon_1}{3} C(z) \left(Y_{N_1} - b\, Y_{N_1}^{\text{eq}} \right) - Y_{B-L} \left[a\, d\, W_{N_1}^{\Delta L=1}(z) + e\, W_{N_2}^{\Delta L=2}(z) \right] . \tag{3.2}$$

The exact definitions of the functions $C(z)$, $W_{N_1}^{\Delta L=1}(z)$, and $W_{N_2}^{\Delta L=2}(z)$ can be obtained from eqs. (A.2), where $W_{N_1}^{\Delta L=1}$ includes the washouts induced by processes with an on-shell N_1, and $W_{N_2}^{\Delta L=2}$ those mediated by a virtual N_2. What is important for this study is that

$$C(z) \propto (\lambda^\dagger \lambda)_{11}, \ W_{N_1}^{\Delta L=1} \propto (\lambda^\dagger \lambda)_{11}, \ W_{N_2}^{\Delta L=2} \propto (\lambda^\dagger \lambda)_{22}^2, \ \text{and} \ \epsilon_1 \propto (\lambda^\dagger \lambda)_{22} . \tag{3.3}$$

To keep it simple, we have neglected the washouts associated to interactions with an on-shell N_2, which is a good approximation for $M_2/M_1 \gtrsim 6$. The analysis can be extended to include almost degenerate neutrinos, but we think the conclusions obtained in section 2.2 for the IDM are already a good indicative of the required hierarchy in couplings and degeneracy in masses. Moreover, in eq. (3.2) we have summed over all lepton flavors, defining $\epsilon_1 \equiv \sum_\alpha \epsilon_{\alpha 1}$. The reason is that for all the mechanisms studied in this work -except the resonant one- the maximum baryon asymmetry is obtained when all the projectors $K_{\alpha i}$ are equal, which allows to write a single BE for the evolution of the matter-antimatter asymmetry. However we note that again, it would be possible to extend — at least some — of the results to a more complex flavor scenario.

The quantities a, b, c, d, and e parametrize possible differences between the IDM and other models, with respect to the generation of baryon asymmetry. For the IDM, $a = b = c = d = e = 1$, while for other models it is to expect that they differ from 1 at most by a factor ~ 10. Let us call $Y_B^{\max}(M_1)$ the maximum baryon asymmetry that can be produced for a given value of M_1 (and also a fixed value of T_{sfo}). Then we have that:

- $Y_B^{\max}(M_1) \propto c$: the baryon asymmetry is proportional to the numerical constants multiplying this — so called source — term of the BEs.

- $Y_B^{\max}(M_1) \propto b$: this is because the baryon asymmetry is proportional to $Y_{N_1}^{\mathrm{eq}}(T \gg M_1)$ and b can be effectively absorbed into $Y_{N_1}^{\mathrm{eq}}$.

- $Y_B^{\max}(M_1) \propto e^{-1/2}$: since $Y_B^{\max}(M_1)$ is obtained after scanning over all possible values of $(\lambda^\dagger \lambda)_{22}$, it is clear from the relations (3.3) that a factor e multiplying $W_{N_2}^{\Delta L=2}$ is equivalent to a factor $e^{-1/2}$ multiplying the CP asymmetry $\big($this can also be seen with the change of variables $(\lambda^\dagger\lambda)_{22} = \widetilde{(\lambda^\dagger\lambda)}_{22}/\sqrt{e}\big)$.

- $Y_B^{\max}(M_1) \propto a^{-1}$: the explanation is similar to the one above, but now it must be taken into account that $Y_B^{\max}(M_1)$ is also the maximum over all values of $(\lambda^\dagger\lambda)_{11}$. Then a scan over $(\lambda^\dagger\lambda)_{11}$ is equivalent to a scan over $\widetilde{(\lambda^\dagger\lambda)}_{11} = a(\lambda^\dagger\lambda)_{11}$, and in terms of $\widetilde{(\lambda^\dagger\lambda)}_{11}$ the factor a only appears dividing the source term.

- $Y_B^{\max}(M_1) \propto d^0$ (i.e. $Y_B^{\max}$ is independent of d): actually this is only true for the mechanisms described in sections 2.3 and 2.4, because the washouts due to the inverse decays of N_1 are negligible. Instead, for the case described in section 2.1 (which cannot give baryogenesis at the TeV scale), we have not found a simple rule to describe the dependence on d.

Putting all together,

$$Y_B^{\max}(M_1) = \frac{c\,b}{a\,\sqrt{e}}Y_B^{\max\,\mathrm{IDM}}(M_1)\,,\tag{3.4}$$

where $Y_B^{\max\,\mathrm{IDM}}(M_1)$ is the maximum baryon asymmetry for a given M_1 in the IDM.

To complete the analysis we give the — approximate — dependence of $Y_B^{\max\,\mathrm{IDM}}$ on M_1, valid for values of $Y_B^{\max\,\mathrm{IDM}}$ close to Y_B^{obs}. This will allow to see how the lower bound on M_1 depends on the parameters of a model.

- Massless decay products and zero initial N_1-density (see section 2.1): $Y_B^{\max\,\mathrm{IDM}} \propto \sqrt{M_1}$. This is because the $\Delta L = 2$ washouts are proportional to $(\lambda^\dagger\lambda)_{22}^2\,m_\mathrm{P}/M_1$, with m_P the Planck mass. Therefore as M_1 increases, $(\lambda^\dagger\lambda)_{22}$ can rise proportionally to $\sqrt{M_1}$ without changing the washouts, but increasing the CP asymmetry by a factor $\sqrt{M_1}$.

  Altogether the lower bound on M_1, obtained by setting $Y_B^{\max} = Y_B^{\mathrm{obs}}$, is $M_{1\min} \propto a^2 e/(c\,b)^2$. Hence, if e.g. the CP asymmetry can be enhanced by a factor 2 ($c = 2$) without changing the washouts, then $M_{1\min}$ decreases by a factor 4. However, an increase in ϵ_1 is usually accompanied by an increase in the washouts (see figure 1), therefore it is to expect that the bound derived in section 2.1 stays well above $10\,\mathrm{TeV}$, and typically close to $\sim 100\,\mathrm{TeV}$.

- Massless decay products and initial thermal N_1-density (see section 2.3): $Y_B^{\max\,\mathrm{IDM}}$ depends roughly linearly on M_1 (more precisely we have found that, approximately, $Y_B^{\max\,\mathrm{IDM}} \propto M_1^{0.9}$). This can be understood as follows: the most important washouts are the ones mediated by an off-shell N_2, which are proportional to $\gamma_{\Delta L=2}/Hn_\ell^{\mathrm{eq}} \propto (\lambda^\dagger\lambda)_{22}^2 \frac{m_\mathrm{P}T_{\mathrm{sfo}}}{M_2^2}$ at the time the asymmetry is generated ($T \sim T_{\mathrm{sfo}}$). For a fixed value of M_2/M_1 and T_{sfo}, $\gamma_{\Delta L=2}/Hn_\ell^{\mathrm{eq}} \propto (\lambda^\dagger\lambda)_{22}^2/M_1^2$, hence a rise in M_1 by a factor f allows an increase in $(\lambda^\dagger\lambda)_{22}$ by the same amount without changing the washouts, but with a CP asymmetry — and consequently a $Y_B^{\max\,\mathrm{IDM}}$ — f times larger.

- Massive decay products (see section 2.4): $Y_B^{\max\,\mathrm{IDM}} \propto M_1^p$, where p depends on M_{H_2}/M_1. Our scan of the parameter space is not fine enough to determine p with good precision, but typically we have found that $p \gtrsim 4$. The important point for this discussion is that p is significantly above 1, even for M_{H_2}/M_1 as low as 0.1. This implies that $M_{1\mathrm{min}}$ is very insensitive to the details of a model (encoded in the parameters $a, \ldots, e$). The reason for the high values of p is related to the exponential suppression of the washouts and the existence of a limiting temperature, T_{sfo}, for generating the asymmetry. Roughly speaking, a small increase in M_1 allows the lepton asymmetry to be generated at slightly larger values of $z = M_1/T$, when the $\Delta L = 2$ washouts are much smaller due to their exponential suppression. This allows to take much larger values of $(\lambda^\dagger\lambda)_{22}$ and consequently of the CP asymmetry.

As an example of how to apply these results, consider adding a third neutrino, N_3, with $M_3 \gtrsim 6M_2 \gtrsim 36M_1$, so that the main contribution to the BAU comes from the decay of N_1 and the washouts due to processes with on-shell N_2 and N_3 are not important. The phases of the Yukawa couplings of N_2 and N_3 have to be equal if they are chosen to maximize the contribution of N_2 and N_3 to $\epsilon_{\alpha 1}$. Then, if due to N_3 the source term increases by a factor n (i.e. $c = n$), the $\Delta L = 2$ washouts increase by $\sim n^2$ (i.e. $e = n^2$). Therefore, according to eq. (3.4), $Y_B^{\max}$ stays — approximately — the same when adding the N_3.

4 Dark matter, leptogenesis, and neutrino masses?

In section 2 we have analyzed in detail how to achieve leptogenesis at the TeV scale in the IDM. The main goal has been to obtain the conditions for low energy baryogenesis that are independent of a possible connection to neutrino masses. In this regard the IDM has been considered just as a "toy" model to illustrate different mechanisms for baryogenesis at low energy scales, while in section 3 we explained how to extrapolate the conclusions to other models. But the IDM is very attractive by itself, because it is a testable model that provides a DM candidate, can explain LNM, and we have also seen that leptogenesis is feasible via different ways.

The question of whether or not LNM, DM, and the BAU can be explained all together at low energies has been addressed before [55, 56]. In [55], after some quantitative estimates, the authors conclude that it is indeed possible, even for non-degenerate heavy neutrinos, and they give a numerical example. On the contrary, in [56] the authors claim that for non-degenerate neutrinos the BAU yield is between 1 and 2 orders of magnitude below the observed value. However, their study of the parameter space does not seem exhaustive and their expression for the CP asymmetry differs from ours.[4] Therefore we revisit this interesting question here.

For an almost degenerate spectrum of heavy neutrinos we confirm the results found in [56], namely that it is possible to obtain simultaneously the LNM, DM density, and the BAU. For the less trivial non-degenerate case the key to answer this issue is, in leptogenesis, (i) to distinguish between different initial conditions for the density of the neutrinos, and (ii) to include a proper treatment of the M_{H_2} effects, as described in section 2.4.

Let us first recall the expression for the light neutrinos mass matrix in the IDM [18],

$$(m_\nu)_{\alpha\beta} = \frac{\lambda_5 v^2}{8\pi^2} \sum_k \lambda_{\alpha k}\lambda_{\beta k} \frac{M_k}{M_k^2 - M_{H_2}^2}\left(\frac{M_k^2}{M_k^2 - M_{H_2}^2}\ln\frac{M_k^2}{M_{H_2}^2} - 1\right)\,,\qquad(4.1)$$

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[4]In particular, the CP asymmetry must be null when $M_{H_2} = M_1$, but this suppression is not present in the expression given in [56].

where λ_5 is the coupling of the quartic term $\lambda_5\left(H_1^\dagger H_2\right)^2/2$ of the scalar potential. This coupling is responsible for the mass splitting of the neutral components H_0 and A_0 of the inert doublet after the electroweak symmetry breaking. There are several lower bounds on λ_5 which are described next.

To avoid inelastic scatterings with nuclei mediated by a Z boson, that are excluded by many orders of magnitude in direct DM searches, it is necessary to impose the kinematic condition (see e.g. [57])

$$\frac{\lambda_5 v^2}{M_{H_2}} \simeq |M_{A_0} - M_{H_0}| \gtrsim \frac{1}{2} M_{H_2} v_{\mathrm{DM}}^2 \ . \tag{4.2}$$

Here $v_{\mathrm{DM}} \sim 0.7 \times 10^{-3}$ is the DM velocity (in units of c), which we take approximately equal to the velocity of the sun in the galaxy. An independent constraint on λ_5 arises if one requires a Boltzmann suppression of the washouts in leptogenesis (see section 2.4), but this condition turns out to be weaker than (4.2). Finally, the DM relic density can be obtained in two mass windows, $50\ \mathrm{GeV} \lesssim M_{\mathrm{LSP}} \lesssim 75\,\mathrm{GeV}$ and $M_{\mathrm{LSP}} \gtrsim 500\,\mathrm{GeV}$,[5] where M_{LSP} is the smallest among M_{A_0} and M_{H_0} [59–61]. However in the low mass window $\lambda_5 \gtrsim 10^{-2}$ [57] for the inert doublet to be the only DM component. For these large values of λ_5, successful leptogenesis at low energies is not compatible with neutrino masses [55], unless a very large fine tuning among the phases of the Yukawa couplings is invoked.[6] Hence we will work in the high mass window.

In order to obtain the BAU the couplings of N_1 must be very tiny, resulting in a contribution of N_1 to the LNM much smaller than the square root of the solar mass splitting. Therefore in the minimal scenario there have to be three heavy neutrino species, $N_{1,2,3}$, one of the light neutrinos is — almost — massless, and another, the heaviest, has a mass $\simeq \sqrt{\Delta m_{\mathrm{atm}}^2} \simeq 0.05\,\mathrm{eV}$.

It is not the goal of this work to provide a fit to all neutrino oscillation data. Instead we have asked whether successful leptogenesis is possible with the contribution of N_2 to LNM being of order 0.05 eV. If that is the case, it is to expect that the measured mixing angles and squared mass differences can be easily accommodated, because leptogenesis is not very sensitive to the exact structure of the Yukawa coupling matrix. We have found that the answer to that question depends on the initial condition for the density of the heavy neutrinos:

- Zero initial density of N_1: in this case the N_1 are only produced by the Yukawa interactions. We have verified that the BAU, DM, and LNM can be explained all together only if a fine tuning among the phases of the Yukawa couplings of N_2 and N_3 is allowed. To be more concrete, let us define the contribution of N_2 to the LNM as the quantity

$$m_\nu(M_2) = \frac{\lambda_5 v^2}{8\pi^2}(\lambda^\dagger\lambda)_{22}\frac{M_2}{M_2^2 - M_{H_2}^2}\left(\frac{M_2^2}{M_2^2 - M_{H_2}^2}\ln\frac{M_2^2}{M_{H_2}^2} - 1\right) \ . \tag{4.3}$$

Then if e.g. $m_\nu(M_2)$ is allowed to take values as large as 5, 1, and 0.5 eV, the lower bound on M_1 for successful leptogenesis is, respectively, $\sim 1700\ \mathrm{GeV}, \sim 2000\ \mathrm{GeV}$, and

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[5]Actually, if there are coannihilations between the DM particles and some of the singlet neutrinos N_i, the observed abundance of DM can be obtained also for $100\ \mathrm{GeV} \lesssim M_{\mathrm{LSP}} \lesssim 500\,\mathrm{GeV}$ [58]. However this requires that the N_i be almost degenerate with the lightest scalar, resulting in M_i values that are too low for successful leptogenesis (unless a couple of the heavy neutrinos are almost degenerate, allowing for resonant leptogenesis).

[6]Resorting to large fine tunings of this type might demand extra care due to the running of the parameters [62].

above 10 TeV. As an example consider

$$M_1 = 2000 \, \text{GeV}, \qquad \frac{M_2}{M_1} = 3, \qquad \frac{M_{H_2}}{M_1} = 0.6,$$
$$(\lambda^\dagger \lambda)_{11} = 3 \times 10^{-15}, \qquad (\lambda^\dagger \lambda)_{22} = 6 \times 10^{-4}, \qquad K_{\alpha i} = 1/3 \,, \tag{4.4}$$

where $m_\nu(M_2) \sim 1 \, \text{eV}$ when the smallest value of λ_5 compatible with the inequality (4.2) is chosen. As we explained above, in the minimal scenario to explain LNM it is necessary to add another singlet, N_3, and in this case it must give a similar contribution to LNM as N_2. This allows to obtain LNM of the order of $\sqrt{\Delta m_{\text{atm}}^2}$ or $\sqrt{\Delta m_{\text{sun}}^2}$ by fine tuning the phases of the Yukawa couplings (see e.g. [12] for the type I seesaw). These results hold for $T_{\text{sfo}} = 80 \, \text{GeV}$. Instead, if $T_{\text{sfo}} = 140 \, \text{GeV}$, the lower bound on M_1 is 2800 (3500) GeV when the contribution of N_2 to the LNM is allowed to take values as large as ~ 5 (1) eV.

- Initial thermal population of N_1: if the N_1 have a thermal distribution at the onset of leptogenesis, the DM, BAU, and LNM can be simultaneously explained in the IDM, without a fine tuning of the Yukawa coupling phases, for $M_1 \gtrsim 750 \, \text{GeV}$. The exact lower bound has a mild dependence on M_2/M_1, but it stays in the range $750 - 850 \, \text{GeV}$ for $3 \lesssim M_2/M_1 \lesssim 50$. As an example, consider

$$M_1 = 850 \, \text{GeV}, \qquad \frac{M_2}{M_1} = 3, \qquad \frac{M_{H_2}}{M_1} = 0.6,$$
$$(\lambda^\dagger \lambda)_{11} = 1.4 \times 10^{-15}, \qquad (\lambda^\dagger \lambda)_{22} = 2 \times 10^{-5}, \qquad K_{\alpha i} = 1/3 \,. \tag{4.5}$$

In this example the contribution of N_2 to the LNM is of the order of $\sqrt{\Delta m_{\text{atm}}^2}$ (again taking for λ_5 the smallest value that respects the condition (4.2)). As already explained, since at least two light neutrinos have mass, it is necessary to add a third singlet N_3, whose contribution to leptogenesis will be $O(1)$ for a inverted hierarchy of LNM and negligible for a normal one (in the latter case the Yukawa couplings of N_3 would be typically quite smaller than those of N_2). These results hold for $T_{\text{sfo}} = 80 \, \text{GeV}$. Instead, if $T_{\text{sfo}} = 140 \, \text{GeV}$, the lower bound on M_1 rises to $\sim 1200 \, \text{GeV}$.

5 Conclusions and outlook

We have studied different mechanisms for thermal BarPaDe in the framework of the IDM. For non-degenerate heavy neutrinos we have found the following: (i) If the neutrinos that generate the asymmetry, N_1, can only be produced by their — CP violating — Yukawa interactions and the masses of the decay products are negligible (i.e. $M_{H_2} \ll M_1$), successful leptogenesis requires $M_1 \gtrsim 80 \, \text{TeV}$. The lower bound on M_1, $M_{1\text{min}}$, is given in figure 2 as a function of M_2/M_1. (ii) If the masses of the decay products are negligible, but the N_1 have a thermal density at the start of leptogenesis, i.e. $Y_N(T \gg M_1) = Y_N^{\text{eq}}(T \gg M_1)$, then it is possible to generate the BAU at the TeV scale, with the lower bound on M_1 depending on the freeze out temperature of the sphalerons, T_{sfo}. For $T_{\text{sfo}} = 80 \, (140) \, \text{GeV}$, $M_{1\text{min}} \sim 2000 \, (2500) \, \text{GeV}$. For models with perturbative violation of B, baryogenesis is possible, at least in principle, at much lower energies (see section 2.3). (iii) If the N_1 are only produced through the Yukawa interactions, but the mass of H_2 is not so tiny with respect to M_1, it is also possible to have baryogenesis at the TeV scale. This is shown in figure 5, where $M_{1\text{min}}$ is plotted as a function of M_{H_2}/M_1 for two different values of T_{sfo}. Also in this case baryogenesis could be possible at much lower energies if B is violated at the perturbative level.

In turn, for degenerate neutrinos, we have derived an upper bound on the degeneracy parameter $\delta \equiv (M_2 - M_1)/M_1$ as a function of M_1, for different constraints on the hierarchy among the Yukawa couplings of N_1 and N_2. This is shown in figure 3. The hierarchy among the couplings can be parametrized in terms of r, defined as the minimum non-null quantity in the list $\left\{ \sqrt{(\lambda^\dagger \lambda)_{11}/(\lambda^\dagger \lambda)_{22}}, \sqrt{K_{\alpha i}} \, (i = 1, 2, \ \alpha = e, \mu, \tau) \right\}$. Then the relation between the amount of mass degeneracy and coupling hierarchy that are required can be summarized in the following rules: $\delta \times r \lesssim 10^{-8}$ for $M_1 \sim 4\,\text{TeV}$ and $\delta \times r \lesssim 2 \times 10^{-9}$ for $250\,\text{GeV} \lesssim M_1 \lesssim 1\,\text{TeV}$. To derive trustable results for masses lower than $\sim 250\,\text{GeV}$ a more involved treatment including scatterings and thermal effects should be performed.

The conclusions stated above are expected to be generic of models for thermal BarPaDe, apart from not very significant quantitative modifications. This is because the key point behind those results, namely the washouts associated to the kinematical phase of the CP asymmetry (see figure 1), is common to all these models. In section 3 we have studied more quantitatively this issue, giving a recipe to extrapolate the mass bounds from the IDM to other models. We have shown that the bound for case (i) is the most sensitive to model differences (although it is to expect that it always be well above $10\,\text{TeV}$), while the bound for (iii) is very insensitive (it depends mainly on the value of T_{sfo}).

Finally, in section 4 we revisited the issue of whether or not LNM, DM, and the BAU can be explained simultaneously in the IDM, at low energies, and without resorting to the resonant leptogenesis mechanism. We found that it is indeed possible for an initial thermal density of N_1 and $M_1 \gtrsim 750\,\text{GeV}$ (with M_{H_2} in the high mass window, $M_{H_2} \gtrsim 500\,\text{GeV}$). Furthermore, if there are only three singlet neutrinos, one of the light neutrinos has to be — almost — massless. In case that the N_1 can only be produced by their Yukawa interactions, it is still possible to encompass LNM, DM, and the BAU, for $M_1 \gtrsim 1700\,\text{GeV}$, but quite a lot of fine tuning among the Yukawa phases of N_2 and N_3 is required.

We note that in [63] a similar analysis to that of point (i) was performed, but for models with small violation of $B - L$, so that the most important contribution to the CP asymmetries comes from the part that conserves total lepton number. The lower bound found was $M_1 \gtrsim 10^6\,\text{GeV}$.[7] Another interesting possibility not analyzed in this work was proposed in [64], namely to generate the BAU in three-body decays of a heavy particle (with two body decays forbidden). The basic idea is that washout processes involving three particles in the initial or final state are naturally phase-space suppressed with respect to $1 \leftrightarrow 2$ and $2 \leftrightarrow 2$ interactions, while the CP asymmetry could be still sizeable. It would be worth to explore more this possibility and confirm that actually all washout processes can be suppressed without reducing too much the CP asymmetry. Also, arguments for a new way to TeV scale leptogenesis have been recently presented in [65].

From an experimental point of view, the results of this work are useful because they show in detail how BarPaDe can be achieved for particle masses somewhat below $1\,\text{TeV}$ in the case of leptogenesis and well below the electroweak scale if baryon number is violated perturbatively. However in order to detect a particle — directly or indirectly — it is also necessary that its couplings to SM fields be not too tiny. This is certainly not the case for the particle producing the matter-antimatter asymmetry (N_1) whose couplings to SM fields cannot be much larger than $\sim 10^{-7}$ for baryogenesis at or below the TeV scale (this is true

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[7]Note that this bound is ~ 12 times larger than the lower bound found for case (i), $M_{1\text{min}} \sim 80\,\text{TeV}$. The main reasons for this difference are that the bound in [63] was calculated in a two — instead of three — flavor regime, the L-conserving part of the CP asymmetry is smaller than the L-violating one, and some of the projectors must be small when $\epsilon = 0$ (resulting in smaller flavored CP asymmetries). On top of this it must be taken into account that the lower bound on M_1, for massless decay products and zero initial density, is quadratically sensitive to factors modifying the final amount of baryon asymmetry.

except possibly in the resonant mechanism, see e.g. [25]). Instead, the couplings of N_2 (the particle contributing virtually to the CP asymmetry in N_1 decays) can in principle be quite sizeable. Their maximum size is very model dependent, as are the possible experimental signals. E.g. for the IDM to explain the BAU with non-degenerate singlet neutrinos, DM, and LNM without fine tuning among the phases of the Yukawa couplings, $\lambda_{\alpha 2} \lesssim 0.01$ for $M_2 \sim 2.5\,\mathrm{TeV}$ (see the example given in eq. (4.5)). Those couplings are still too small to induce signals in forthcoming experiments. However it should be noted that the bound $\lambda_{\alpha 2} \lesssim 0.01$ does not come from baryogenesis, but is a non-trivial consequence of the smallness of LNM and the lower bound on λ_5 due to direct DM searches. This illustrates the strong model dependence of the allowed range for the couplings. Moreover if some fine tuned cancellations among the contributions of N_2 and N_3 to LNM is allowed, larger Yukawa couplings are possible. Experiments searching for charged lepton flavor violation and lepton colliders would be the most promising ways to search for these neutrinos [66–68].

Acknowledgments

A   Boltzmann equations

The lagrangian for the IDM above the electroweak phase transition, on a basis in which the singlet neutrinos N_i are mass eigenstates, can be written as

$$\mathcal{L} = \mathcal{L}_{\mathrm{SM}} - V(H_1, H_2) + \frac{i}{2}\overline{N}_i \slashed{\partial} N_i - \frac{M_i}{2}\overline{N}_i N_i + (D_\mu H_2)^\dagger (D^\mu H_2) - \lambda_{\alpha i}\,\widetilde{H}_2^\dagger\,\overline{P_R N_i}\ell_\alpha + \mathrm{h.c.}\,,\ \ (\mathrm{A.1})$$

where α, i are family indices ($\alpha = e, \mu, \tau$ and $i = 1, 2, \dots$), ℓ_α are the leptonic SU(2) doublets, H_2 is the inert Higgs doublet ($\widetilde{H}_2 = i\tau_2 H_2^*$, with τ_2 Pauli's second matrix), and $P_{R,L}$ are the chirality projectors. Moreover, $V(H_1, H_2)$ is the scalar potential, containing, among others, the term $\lambda_5(H_1^\dagger H_2)^2/2$.

The BEs we have used are (see section 2 for comments on the approximations that have been made)

$$\frac{\mathrm{d}Y_{N_i}}{\mathrm{d}z} = \frac{-1}{sHz}\left(\frac{Y_{N_i}}{Y_{N_i}^{\mathrm{eq}}} - 1\right)\sum_\alpha 2\gamma_{\ell_\alpha H_2}^{N_i}\,,$$

$$\begin{aligned}
\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} = \frac{-1}{sHz}\Bigg\{ &\sum_i \left(\frac{Y_{N_i}}{Y_{N_i}^{\mathrm{eq}}} - 1\right)\epsilon_{\alpha i}\,2\gamma_{\ell_\alpha H_2}^{N_i} \\
&- \sum_i \gamma_{\ell_\alpha H_2}^{N_i}\,[y_{\ell_\alpha} + y_{H_2}] \\
&- \sum_\beta (1 + \delta_{\alpha\beta})\left(\gamma_{\bar{\ell}_\beta \bar{H}_2}^{\ell_\alpha H_2}{}' + \gamma_{\bar{H}_2 \bar{H}_2}^{\ell_\beta \ell_\alpha}\right)[y_{\ell_\alpha} + y_{\ell_\beta} + 2y_{H_2}] \\
&- \sum_{\beta\neq\alpha}\left(\gamma_{\ell_\alpha H_2}^{\ell_\beta H_2}{}' + \gamma_{\ell_\alpha \bar{H}_2}^{\ell_\beta \bar{H}_2} + \gamma_{\ell_\alpha \bar{\ell}_\beta}^{H_2 \bar{H}_2}\right)[y_{\ell_\alpha} - y_{\ell_\beta}]\Bigg\}\,,
\end{aligned}\quad (\mathrm{A.2})$$

where $z \equiv M_1/T$, $Y_X \equiv n_X/s$ is the number density of the particle specie X normalized to the entropy density, and $y_X \equiv Y_{\Delta_X}/Y_X^{\text{eq}} \equiv (Y_X - Y_{\bar{X}})/Y_X^{\text{eq}}$ is the asymmetry density normalized to the equilibrium density. For a massless particle X we use the convention that Y_X gives the density of a single degree of freedom of X. We have also defined $Y_{\Delta_\alpha} \equiv Y_B/3 - Y_{L_\alpha}$, where Y_B is the baryon asymmetry and $Y_{L_\alpha} = (2y_{\ell_\alpha} + y_{e_{R\alpha}})Y^{\text{eq}}$ is the total lepton asymmetry in the flavor α (with $Y^{\text{eq}} \equiv Y_{\ell_\alpha}^{\text{eq}} = Y_{e_{R\alpha}}^{\text{eq}}$, and $e_{R\alpha}$ the right handed charged leptons). Finally, we have also introduced the notation $\gamma_{c,d,\ldots}^{a,b,\ldots} \equiv \gamma(a, b, \ldots \to c, d, \ldots)$ for the reaction densities. The prime on some rates indicates that the on-shell contribution has to be subtracted, and the $(1 + \delta_{\alpha\beta})$ factor, with $\delta_{\alpha\beta} = 1(0)$ for $\alpha = (\neq)\beta$, takes into account that processes with $\alpha = \beta$ change Y_{Δ_α} by two units.

The CP asymmetries in N_i decays, $\epsilon_{\alpha i}$, are defined as

$$\epsilon_{\alpha i} \equiv \frac{\Gamma\left(N_i \to \ell_\alpha H_2\right) - \Gamma\left(N_i \to \bar{\ell}_\alpha \bar{H}_2\right)}{\sum_\beta \Gamma\left(N_i \to \ell_\beta H_2\right) + \Gamma\left(N_i \to \bar{\ell}_\beta \bar{H}_2\right)} \; .$$

To calculate them it is important to take into account that $M_{H_2} \neq 0$. Assuming that $M_j - M_i \gg \Gamma_{i,j}$, with Γ_i the decay width of N_i, we get (for $M_{H_2} = 0$ see [10])

$$\epsilon_{\alpha i} = \frac{-1}{8\pi(\lambda^\dagger\lambda)_{ii}} \sum_{j \neq i} \left[\frac{M_i M_j}{M_j^2 - M_i^2} - f(a_j) \right] (1 - a_{H_2})^2 \text{Im}\left[\lambda_{\alpha j}^* \lambda_{\alpha i}(\lambda^\dagger\lambda)_{ji}\right] + \epsilon_{\alpha i}^L \qquad (A.3)$$

$$= \frac{-1}{8\pi} \sum_{j \neq i} \left[\frac{M_i M_j}{M_j^2 - M_i^2} - f(a_j) \right] (1 - a_{H_2})^2 (\lambda^\dagger\lambda)_{jj} \sqrt{K_{\alpha i}}\sqrt{K_{\alpha j}} \sum_\beta \sqrt{K_{\beta i}}\sqrt{K_{\beta j}} \, p_{\alpha\beta}^{ij} + \epsilon_{\alpha i}^L \; ,$$

where $a_j \equiv M_j^2/M_i^2$, $a_{H_2} \equiv M_{H_2}^2/M_i^2$, $f(x) = \sqrt{x}(1 - (1 + x) \ln\left[(1 + x)/x\right])$, and $\epsilon_{\alpha i}^L$ is defined below. For the second line of eq. (A.3) we have used the following parametrization of the Yukawa couplings in terms of the projectors $K_{\alpha i}$ and some phases $\phi_{\alpha i}$:

$$\lambda_{\alpha i} = \sqrt{K_{\alpha i}}\sqrt{\left(\lambda^\dagger\lambda\right)_{ii}} e^{i\phi_{\alpha i}} \; , \qquad (A.4)$$

where

$$K_{\alpha i} = \frac{\lambda_{\alpha i}\lambda_{\alpha i}^*}{\left(\lambda^\dagger\lambda\right)_{ii}} \; , \qquad (A.5)$$

$$p_{\alpha\beta}^{ij} = p_{\beta\alpha}^{ij} = -p_{\alpha\beta}^{ji} = \sin\left(\phi_{\alpha i} - \phi_{\alpha j} + \phi_{\beta i} - \phi_{\beta j}\right) \; . \qquad (A.6)$$

For each pair of neutrinos, $\{N_i, N_j\}$, there are three independent combination of phases, $\phi_{\alpha i} - \phi_{\alpha j}$, $\alpha = e, \mu, \tau$. In order to maximize the CP asymmetries, we have taken $\phi_{\alpha 1} - \phi_{\alpha 2} = \pi/4 \; \forall \alpha$. There is also a contribution to the CP asymmetries that does not violate total lepton number, $\epsilon_{\alpha i}^L$, so that $\sum_\alpha \epsilon_{\alpha i}^L = 0$ [10]. However, with the choice of phases we have made, $\epsilon_{\alpha i}^L = 0$.

The BEs (A.2) are complemented by a set of relations among the density asymmetries, which is obtained considering the fast interactions not included in eqs. (A.2) and the conservation laws [69, 70]. Since the scenario considered in this work takes place at temperatures $T \ll 10^5$ GeV, the analysis is quite similar to the low mass regime described in [27, 70]. The only difference comes from the additional scalar doublet H_2. We assume that the interaction $H_2\bar{H}_1 \to \bar{H}_2 H_1$ is fast while there is a significant amount of H_2 in the thermal bath (the

reason is explained in section 2.4), so that $\mu_{H_2} = \mu_{H_1}$, with μ_X the chemical potential of X. This implies that $Y_{\Delta H_2} = g Y_{\Delta H_1}$, with g defined below. Altogether we get

$$Y_{\ell_\alpha} = \frac{4(Y_{B-L} - Y_{\Delta_\alpha})(4+g) - Y_{\Delta_\alpha}(221+35g)}{9(79+13g)} \, , \qquad Y_B = \frac{4(7+g)}{79+13g} Y_{B-L} \, ,$$

$$Y_{\Delta H} = -\frac{8}{79+13g} Y_{B-L} \, , \qquad\qquad\qquad Y_{\Delta H_2} = -\frac{8g}{79+13g} Y_{B-L} \, , \qquad \text{(A.7)}$$

with

$$
\begin{aligned}
g = g(M_{H_2}/T) &= \frac{1}{2}\left(\frac{M_{H_2}}{T}\right)^2 K_2\left(M_{H_2}/T\right) \\
&\approx \sqrt{\frac{\pi}{8}}\left(\frac{M_{H_2}}{T}\right)^{3/2} e^{-M_{H_2}/T} \qquad \text{for} \quad T \ll M_{H_2} \, ,
\end{aligned}
\tag{A.8}
$$

and K_2 the modified Bessel function of second type. Two comments are in order:

- The relations (A.7) change slightly with the temperature between T_c and T_{sfo}, but the effect on our analysis is very small, so we have neglected this variation.

- The BEs have been derived assuming kinetic equilibrium and Maxwell-Boltzmann statistics for the distribution functions appearing explicitly in the transport equations (instead for the total number of relativistic degrees of freedom we have used the standard expression that distinguishes between fermions and bosons). This is usually a good approximation, especially in the strong washout regime [71–73]. However, when spectator processes are taken into account, the — correct — use of quantum statistical distributions (including the Fermi-Dirac blocking factor and the stimulated emission factor for bosons) brings a relative factor of $1/2$ between the washout terms induced by bosons and fermions (see [74] or the appendix A of [75] for details), which is not necessarily negligible. The reason is that what really multiplies the rates in the washout part is not the difference between the density asymmetries, but the difference between the corresponding chemical potentials. This effect can be taken into account by replacing y_X by $\mathcal{y}_X$ for the massless particles in the above BEs, where $\mathcal{y}_X \equiv Y_{X-\bar{X}}/Y_f^{\mathrm{eq}}$ ($Y_{X-\bar{X}}/Y_s^{\mathrm{eq}}$) for fermions (bosons) and $Y_f^{\mathrm{eq}} \equiv \frac{1}{2} Y_s^{\mathrm{eq}} \equiv \frac{15}{8\pi^2 g_*}$. The results of this paper are only slightly modified by that replacement.

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