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A consistent model for leptogenesis, dark matter and the IceCube signal

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Abstract: We discuss a left-right symmetric extension of the Standard Model in which the three additional right-handed neutrinos play a central role in explaining the baryon asymmetry of the Universe, the dark matter abundance and the ultra energetic signal detected by the IceCube experiment. The energy spectrum and neutrino flux measured by IceCube are ascribed to the decays of the lightest right-handed neutrino N_1, thus fixing its mass and lifetime, while the production of N_1 in the primordial thermal bath occurs via a freeze-in mechanism driven by the additional $\mathrm{SU}(2)_R$ interactions. The constraints imposed by IceCube and the dark matter abundance allow nonetheless the heavier right-handed neutrinos to realize a standard type-I seesaw leptogenesis, with the $B-L$ asymmetry dominantly produced by the next-to-lightest neutrino N_2. Further consequences and predictions of the model are that: the N_1 production implies a specific power-law relation between the reheating temperature of the Universe and the vacuum expectation value of the $\mathrm{SU}(2)_R$ triplet; leptogenesis imposes a lower bound on the reheating temperature of the Universe at $7 \times 10^9\,\mathrm{GeV}$. Additionally, the model requires a vanishing absolute neutrino mass scale $m_1 \simeq 0$.

Keywords: Beyond Standard Model, Cosmology of Theories beyond the SM, Neutrino Physics

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1 Introduction

The IceCube experiment has so far reported evidence for extraterrestrial high-energy neutrinos [1–4], which cannot be explained by the atmospheric and prompt neutrino components [5–7]. It is possible to relate this signal to astrophysical sources, see for example refs. [8–10] for a general discussion, and references therein,[1] but it is certainly intriguing to speculate about a new-physics origin of these events. Some models have been proposed, explaining the high-energy IceCube signal as the decay of a long-lived particle [11, 12], that can also constitute a viable dark matter (DM) candidate [13–32].

A model of right handed neutrino DM predicting a signal in high energy neutrinos and able to reproduce the matter-antimatter asymmetry with leptogenesis was presented in ref. [23]. Ref. [24] applied a similar framework to interpret the IceCube signal. These are also the aims of this work: as in ref. [23, 24], we will make use of additional heavy right-handed (RH) neutrinos which are responsible for neutrino masses and mixing, via the seesaw mechanism [33–40], and leptogenesis [41]. However, while in ref. [24] the heavy neutrinos are produced by means of inflaton decay, here we will seek for a production mechanism directly from the thermal bath.

We therefore consider a left-right symmetric model (LRSM) [42–45] (see also [46–48] for reviews), with gauge group $\mathrm{SU}(2)_L \otimes \mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{\tilde{Y}}$ in which three RH neutrinos N_i are naturally accommodated into three RH doublets of $\mathrm{SU}(2)_R$. Among the three heavy neutrinos, we choose the lightest one, N_1, to be long-lived: the N_1 lifetime is required to be

[1] A lot of work has been done in trying to explain the IceCube events in term of astrophysical sources. We refer for example to the references contained in [8–10] for specific discussions.

long enough to make N_1 a viable dark matter candidate, and at the same time the N_1 decay width into neutrinos needs to be of the right size to make N_1 responsible for the IceCube signal. The energy spectrum of the IceCube events (in particular, its end point) and the measured flux therefore determine the mass and the decay width into neutrinos of N_1. We will show that hadronic decays of N_1 induce an additional decay channel that can be "fast" enough to be in contradiction with, among others, gamma-ray bounds: this implies that our model needs to assume a "hadrophobic" structure [49, 50], in order to stabilize this decay channel. The IceCube flux implies a severe suppression of the coupling to left-handed (LH) lepton and Higgs doublets, preventing N_1 production in the early Universe by means of this type of interactions. Instead, the additional $\mathrm{SU}(2)_R$ gauge interactions in the LRSM can provide a viable way to produce N_1 without spoiling its stability. This is accomplished by means of a freeze-in mechanism [51], tuned to reproduce the present-day DM abundance $\Omega_{N_1}h^2 = \Omega_{\mathrm{DM}}h^2 = 0.1198 \pm 0.0015$ [52]. In fact, we found that the freeze-out production mechanism in the context of a LRSM does not lead to the correct relic abundance for PeV-scale DM, since it would require an extremely large entropy dilution. This mechanism has been instead successfully used for keV sterile neutrino DM [53], when sufficient entropy dilution can be achieved [54, 55]. See instead ref. [56] for a discussion of TeV-scale DM in the context of the LRSM and ref. [57, 58] for the freeze-in mechanism in the context of a keV sterile neutrino DM.

The next-to-lightest N_2 and the heaviest N_3 neutrinos in our model are responsible, in turn, for the generation of the matter-antimatter asymmetry of the Universe, via standard thermal leptogenesis. Specifically, we rely on the N_2-dominated scenario of leptogenesis [59], in which N_2's dynamics is able to produce the correct final baryon asymmetry, measured by the baryon-to-photon ratio $\eta_B^{\mathrm{CMB}} = (6.1 \pm 0.1) \times 10^{-10}$ [52]. In this case, the thermal production of N_2 relies on the Yukawa couplings to Higgs and LH doublets, which are not suppressed.

We therefore propose a LRSM in which the lightest heavy neutrino will play the role of the DM particle, while being also responsible, through its decays, for the IceCube signal. At the same time, the other two heavy neutrinos will generate the baryon asymmetry of the Universe via standard thermal leptogenesis.

The paper is organized as follows: in section 2 we present the details of the LRSM model considered in the analysis, in section 3 we focus on the DM particle of the model, the right-handed neutrino N_1, considering in particular the constraints from IceCube and the relic abundance. In section 4 we present the calculation for the baryon asymmetry and in section 5 we draw our conclusions.

2 The model

Considering the standard minimal LRSM [46–48], the RH leptons are fitted into $\mathrm{SU}(2)_R$ doublets:

$$R_i = \begin{pmatrix} N_{Ri} \\ \ell_{Ri} \end{pmatrix}. \tag{2.1}$$

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To ensure the right spontaneous symmetry breaking pattern and a Majorana mass term for the RH neutrinos, we consider a scalar field Δ_R, triplet of $\mathrm{SU}(2)_R$, which is responsible for the breaking to the Standard Model (SM) gauge group. The left-right symmetry implies the existence of an $\mathrm{SU}(2)_L$ triplet Δ_L. The electroweak symmetry breaking (EWSB) is then obtained by exploiting a bi-doublet scalar field Φ. The Yukawa sector then reads:

$$\mathcal{L}_Y = -Y^{(1)}_{ij}\overline{L}_i\Phi R_j - Y^{(2)}_{ij}\overline{L}_i\tilde{\Phi}R_j - Y^{\Delta}_{ij}\left(L_i^T C i\tau_2 \Delta_L L_j + R_i^T C i\tau_2 \Delta_R R_j\right) + \text{h.c}, \tag{2.2}$$

where

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \qquad \Delta_{L,R} = \begin{pmatrix} \frac{1}{\sqrt{2}}\delta^+ & \delta^{++} \\ \delta^0 & -\frac{1}{\sqrt{2}}\delta^+ \end{pmatrix}_{L,R}, \tag{2.3}$$

and $\tilde{\Phi} \equiv \tau_2 \Phi^* \tau_2$ (τ_2 being the second Pauli matrix).

In order to avoid unwanted low-energy effects due to the $\mathrm{SU}(2)_R$ gauge interactions, the $\mathrm{SU}(2)_R \otimes \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_{\tilde{Y}} \to \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ breaking must take place at very high energies. Therefore, the triplet Δ_R acquires a vacuum expectation value (VEV) $\langle \Delta_R \rangle \equiv v_R$ which is suitably large. This implies a Majorana mass matrix for the RH neutrinos given by:

$$M_{ij} = 2Y^{\Delta}_{ij}\, v_R, \tag{2.4}$$

which can then be diagonalised to $D_M = \mathrm{diag}(M_1,\, M_2,\, M_3)$. We shall assume $M_{\Delta_R} \gg M_i$. The VEV v_R also sets the mass of the $\mathrm{SU}(2)_R$ gauge bosons, since $m_{Z_R},\, m_{W_R} \propto v_R$.

Labelling the LH lepton doublet with the flavour index $\alpha = e, \mu, \tau$, the Yukawa couplings of the RH neutrinos with the SM Higgs doublet H are given by:

$$\mathcal{L}_{Y^\nu} = -Y^{\nu}_{\alpha i}\, \overline{L}_\alpha H N_{Ri} + \text{h.c}, \tag{2.5}$$

with

$$Y^\nu \equiv \frac{Y^{(1)}v_1 + Y^{(2)}v_2}{\sqrt{v_1^2 + v_2^2}}, \tag{2.6}$$

where $\langle \phi_i^0 \rangle = v_i$. The SM Higgs VEV is obtained as: $\langle H \rangle = \sqrt{v_1^2 + v_2^2} \equiv v \simeq 174\,\mathrm{GeV}$. Finally, we may have that also Δ_L gets a VEV $v_L \ll v$. In general, we shall also assume $M_{\Delta_L} \gg M_i$.

With this symmetry-breaking pattern, the light neutrino masses are generally given by the combination of a type-I and a type-II seesaw terms. We will assume that the type-II seesaw contribution to the light neutrino masses is negligible, or even vanishing if $v_L = 0$, so that the light neutrino mass matrix is given by:

$$m_\nu = -Y^\nu D_M^{-1} Y^{\nu T}\, v^2. \tag{2.7}$$

The light neutrino masses are then obtained through the PMNS mixing matrix U as:

$$D_m = -U^\dagger\, m_\nu\, U^*, \tag{2.8}$$

where $D_m = \mathrm{diag}(m_1,\, m_2,\, m_3)$.

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The most relevant interaction of the RH neutrinos are then given by the Yukawa coupling to Higgs and lepton doublets, in eq. (2.5), and by the $SU(2)_R$ gauge interactions. We will comment in subsection 3.3 on the possible interactions of the RH neutrinos with the RH triplet Δ_R. If we assume the standard inflationary picture of the early Universe, charged leptons and LH neutrinos are part of the thermal bath, hence we can assume that $SU(2)_R$ interactions are able to produce the RH neutrinos. In the absence of $SU(2)_R$ interactions, the production of the RH neutrinos is possible only through the Yukawa interactions, which become effective at temperatures around the RH neutrino mass.

3 N_1 as the dark matter particle

3.1 Constraints from IceCube

In our model, we assume that the signal detected by IceCube is originated from DM decays. Given the mass pattern, the suitable DM candidate is the lightest heavy neutrino. Therefore, N_1 will be bound to constitute the whole DM content of the Universe and at the same time produce the IceCube signal.

Let us first analyse the constraints obtained from the IceCube data. N_1 decays only through the Yukawa couplings of eq. (2.5). The decay channels are:

$$N_1 \longrightarrow l_\alpha^\mp W^\pm; \qquad N_1 \longrightarrow \nu_\alpha Z, \bar{\nu}_\alpha Z; \qquad N_1 \longrightarrow \nu_\alpha h, \bar{\nu}_\alpha h. \tag{3.1}$$

For $M_1 \gg m_Z$, m_h, we have monochromatic neutrinos with energy $E_\nu \simeq M_1/2$. This will cause a sharp peak and a cutoff in the neutrino energy spectrum, while neutrino cascades will provide a soft tail in the spectrum. From the highest event detected by IceCube [2, 60] the DM mass scale can be directly determined [16, 17]: $M_1 \simeq 4\,\mathrm{PeV}$.

It is then possible to estimate the neutrino flux on Earth from the decay of N_1, assuming it constitutes the total amount of DM in the Universe. By comparing the theoretical prediction to the flux observed by IceCube, once the mass M_1 is set, the N_1's lifetime τ_{N_1} can be directly derived. From ref. [16, 17, 24], the required lifetime of the DM particles is:

$$\tau_{N_1} \simeq 10^{28}\,\mathrm{s}. \tag{3.2}$$

Although these are just approximate determinations of the mass and lifetime of N_1, as inferred from the IceCube data, for the purposes of this paper they can be regarded as a sufficiently good estimation: slight changes in these values will not make any noticeable difference.

The total decay rate $\Gamma_{D1} = \tau_{N_1}^{-1}$ at tree level is given as a function of the Yukawa parameters $Y^\nu_{\alpha 1}$ by the expression [61, 62]:

$$\Gamma_{D1} = \frac{M_1}{8\pi} \sum_\alpha |Y^\nu_{\alpha 1}|^2 \,. \tag{3.3}$$

Eq. (3.2) implies a constraint on the Yukawa couplings $Y^\nu_{\alpha 1}$:[2]

$$\sum_\alpha |Y^\nu_{\alpha 1}|^2 = \frac{8\pi}{M_1\,\tau_{N_1}} \ll 1. \tag{3.4}$$

[2]From this result, it is evident that this setup cannot realize strong thermal leptogenesis [63, 64].

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Due to the seesaw relation, this constraint will be reflected onto the other Yukawa couplings and the light neutrinos spectrum as well. It is then convenient to introduce the complex orthogonal matrix Ω parameterisation [65]:

$$Y^{\nu} = \frac{1}{v} U D_m^{1/2} \Omega D_M^{1/2}, \tag{3.5}$$

such that:

$$\sum_{\alpha} |Y^{\nu}_{\alpha 1}|^2 = \frac{M_1}{v^2} \sum_i m_i |\Omega_{i1}|^2 . \tag{3.6}$$

Hence we obtain:

$$\sum_i m_i |\Omega_{i1}|^2 = \frac{8\pi v^2}{M_1^2 \tau_{N_1}} \simeq 10^{-52}\,\text{eV}. \tag{3.7}$$

Given the light neutrinos mass spectrum with nonzero m_2 and m_3, it is clear that, in order to have a vanishing $\sum_i m_i |\Omega_{i1}|^2$, we must necessarily have $m_1 \simeq 0$ and a complex orthogonal matrix of the form:

$$\Omega \simeq \begin{pmatrix} 1 & \beta \sin\theta - \alpha \cos\theta & \beta \cos\theta + \alpha \sin\theta \\ \alpha & \cos\theta & -\sin\theta \\ -\beta & \sin\theta & \cos\theta \end{pmatrix}, \tag{3.8}$$

with α, β, θ complex and $|\alpha|, |\beta| \simeq 0$. All the Yukawa couplings, and hence all the quantities related to the other heavy neutrinos, can then be derived from a complex orthogonal matrix with the form in eq. (3.8) and a fully hierarchical light neutrino spectrum. We also notice that the requirement $m_1 \simeq 0$ (in particular $m_1 \ll 10^{-4}\,\text{eV}$) implies that this setup cannot be realised within the so-called SO(10)-inspired leptogenesis models [66–73].

This special form of the complex orthogonal matrix Ω was pointed out in ref. [23], where a model with one vanishing eigenvalue in the Yukawa matrix was presented. Thereby, the production of the decoupled heavy neutrinos was obtained through active-sterile neutrino oscillations. Alternatively, it can be achieved through inflaton decay [23, 24].

To summarize, by imposing the bound on N_1's lifetime from the IceCube flux has important consequences on the general setup of the model, due to the seesaw relation.

3.2 Constraints on the model from W_R-mediated decays of N_1

In addition to the decay mode into neutrinos listed in eq. (3.1) that allow an interpretation of the IceCube events, in the simplest realization of the LRSM where both right-handed leptons and right-handed quarks are accommodated in $\text{SU}(2)_R$ doublets, N_1 possesses also an hadronic decay through the mediation of the W_R gauge boson into right-handed charged leptons and quarks: $N_1 \to l_R q_R \bar{q}'_R$ [74, 75]. The decay rate of this process is:[3]

$$\Gamma(N_1 \to l_R q_R \bar{q}'_R) \simeq \frac{3 g_R^4}{2^{10}\pi^3 M_1^3} \int_0^{M_1^2} ds \frac{M_1^6 - 3M_1^2 s^2 + 2s^3}{\left(s - M_{W_R}^2\right)^2 + M_{W_R}^4 \frac{g_R^4}{(4\pi)^2}}. \tag{3.9}$$

[3]Note that we found a decay width $\Gamma(N_1 \to l_R q_R \bar{q}'_R)$, eq. (3.9), that is a factor 2 smaller than the one reported in ref. [74]. Our total decay width, eq. (3.10), is instead a factor of 3 smaller than the one reported in ref. [75].

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	$SU(3)_C$	$SU(2)_L$	$SU(2)_R$	$U(1)_{\tilde{Y}}$
L	**1**	**2**	**1**	$-1/2$
Q	**3**	**2**	**1**	$1/6$
R	**1**	**1**	**2**	$-1/2$
u	**3**	**1**	**1**	$2/3$
d	**3**	**1**	**1**	$-1/3$

Table 1. Multiplet assignment for the hadrophobic LRSM. L and R are fermionic doublets, Q is a quark doublet and u and d stand for the up quark and down quark singlets.

Considering $M_1 \ll M_{W_R}$ and the usual condition $M_{W_R} = g_R v_R$, the total decay width can be cast in this form:

$$\Gamma(N_1 \to l_R q_R \bar{q}'_R) + \Gamma(N_1 \to \bar{l}_R \bar{q}_R q'_R) \simeq \frac{3\, M_1^5}{2^{10}\pi^3\, v_R^4} \tag{3.10}$$

where we have also considered $g_R \ll 4\pi$.

This decay rate implies a lifetime for N_1 larger than the age of the Universe for $v_R > 5 \times 10^{17}\,\text{GeV}$. This would reflect into an additional bound on the results we will present in the next sections, nevertheless leaving the standard LRSM viable. However, extrapolations to the PeV mass range of antiproton bounds on heavy-DM decays [76] and bounds from gamma-rays [22, 77] set much stronger constraints on this decay channel. Even though extrapolations of the knowledge of hadronization processes at such large energies and astrophysical uncertainties are likely present, nevertheless the lifetime associated to this decay channel is plausibly larger than $10^{26}\,\text{s} - 10^{27}\,\text{s}$. This pushes v_R in a trans-planckian regime, complemented by a suitable requirement $g_R < g_L$, in order to keep at least the mass of W_R below the Planck scale. This solution makes the standard LRSM quite contrived.

Instead, an "hadrophobic" LR choice of the representation for the right-handed quarks prevents the decay of N_1 through W_R. This model accommodates the right-handed quarks into singlets of $SU(2)_R$, with a suitable choice of their $\tilde{Y}$ quantum number in order to satisfy the condition $Q = T_{3L} + T_{3R} + \tilde{Y}$. The assignments are summarised in table 1. This model has been considered in the literature in the past, even though for different purposes, see e.g. refs. [49, 50]. A caveat is that the model should be then embedded in a more complete and final theory to cure the problem of anomalies that are present in the hadrophobic LR model. With the assignments of table 1, the triangle diagram involving three $U(1)_Y$ gauge bosons and the one involving two $SU(2)_R$ and one $U(1)_Y$ gauge bosons are anomalous. A way to cure these anomalies is to add an additional right handed doublet of $SU(2)_R$, R', with $Y'_R = 1/2$ and a right handed singlet, X, with $Y_X = -1$ (plus suitable scalar fields to give mass to these additional fermions). Note that the hadrophobic LR model requires also a doublet to give mass to the quarks in the singlet representation, that cannot couple to the bidoublet. These addition however can be safely applied without affecting the results of our analysis.

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3.3 Relic abundance

As mentioned before, we are requiring N_1 to be the DM particle. We must therefore be able to produce the correct abundance. More specifically, given the current values of $\Omega_{\rm DM}$, of the critical density ρ_c and entropy s_0, the current DM abundance $Y^0_{\rm DM} \equiv n_{\rm DM}/s_0$ is:

$$Y^0_{\rm DM} = \frac{\Omega_{\rm DM}\,\rho_c}{M_1 s_0} \simeq 3.82 \times 10^{-10} \left(\frac{\rm GeV}{M_1}\right) \simeq 9.5 \times 10^{-17}, \tag{3.11}$$

for $M_1 = 4$ PeV.

In general, N_1 can be produced from the thermal bath through its interactions given by the $\mathrm{SU}(2)_R$ gauge bosons and the Yukawa couplings in eq. (2.5). We can safely neglect the contribution due to the coupling with Δ_R. Indeed, assuming a very high scale $v_R \gtrsim 10^{14}\,\mathrm{GeV}$, while having $M_1 \simeq 4\,\mathrm{PeV}$, implies that the Yukawa couplings Y^Δ for N_1 are extremely small and the interactions with Δ_R are strongly suppressed. Also the mixing of N_1 with the other heavy neutrinos gives a negligible contribution [23], as well as the decays $N_{2,3} \longrightarrow N_1 l_R l'_R$, mediated by W_R.

The Yukawa interactions with Higgs and lepton doublets give a decay rate at temperature T [79]:

$$\Gamma_{D1}(T) = \Gamma_{D1} \frac{\mathcal{K}_1(T)}{\mathcal{K}_2(T)}, \tag{3.12}$$

where $\mathcal{K}_i(T)$ are the modified Bessel functions of index i.

At the same time, the RH neutrinos are subject to the $\mathrm{SU}(2)_R$ gauge interactions, whose scattering rate is given by:

$$\Gamma_S(T) \equiv n^{\rm eq}_{N_1}(T) \langle \sigma\,|v|\rangle(T), \tag{3.13}$$

where $n^{\rm eq}_{N_1}$ is the equilibrium number density of N_1 and $\langle \sigma\,|v|\rangle$ is the thermally averaged cross section times velocity. The latter can be estimated via the usual neutrino scattering cross section, as [55]:

$$\langle \sigma\,|v|\rangle(T) \simeq G_F^2 T^2 \left(\frac{m_W}{m_{W_R}}\right)^4 \sim G_F^2 T^2 \left(\frac{m_W}{v_R}\right)^4, \tag{3.14}$$

where m_W is the W boson mass and G_F the Fermi constant. Both Γ_{D1} and Γ_S enter the Boltzmann equations that describe the production of N_1. To this aim we can consider the variable $z \equiv M_1/T$ and the number density N_{N_1} of RH neutrinos N_1, computed in a comoving volume containing one heavy neutrino N_1 in ultra-relativistic equilibrium. This can be easily related to the abundance Y_{N_1}. From the definition $N_{N_1} \equiv n_{N_1}(T)/n^{\rm eq}_{N_1}(T \gg M_i)$, we have $N_{N_1} = 4/3\, n_{N_1}(T)/n^{\rm eq}_\gamma(T)$, where $n^{\rm eq}_\gamma$ is the equilibrium number density of photons. This is strictly connected to the entropy $s(T) = \pi^4 g^s_*(T)\, n^{\rm eq}_\gamma(T)/45\zeta(3)$, such that:

$$N_{N_1} = \frac{4}{135} \frac{\pi^4 g^s_*(T)}{\zeta(3)} \frac{n_{N_1}(T)}{s(T)} \simeq 2.40 g^s_*(T)\, Y_{N_1}(T). \tag{3.15}$$

In the temperature range of our interest $g^s_* = g_* = 112$, considering the three heavy neutrinos as relativistic.

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Using z and N_{N_1} we can write [79]:

$$\frac{dN_{N_1}}{dz} = -(D_1(z) + S(z)) \left[N_{N_1}(z) - N_{N_1}^{\rm eq}(z) \right], \tag{3.16}$$

where $D(z)$ accounts for the decay/inverse-decay processes:

$$D_1(z) \equiv \frac{\Gamma_{D1}(z)}{H(z)z}, \tag{3.17}$$

and $S(z)$ for the scattering:

$$S(z) \equiv \frac{\Gamma_S(z)}{H(z)z}. \tag{3.18}$$

Given eqs. (3.3) and (3.4), N_1's decay rate is strongly suppressed, therefore we can just consider $S(z)$ in the equation for the evolution for N_{N_1}. Assuming vanishing initial abundance for N_1 at the end of inflation, i.e. $N_{N_1}(z_{\rm RH}) = 0$ where $z_{\rm RH} = M_1/T_{\rm RH}$ corresponds to the reheating temperature, we can write the Boltzmann equation for N_1 as:

$$\frac{dN_{N_1}}{dz} = S(z) N_{N_1}^{\rm eq}(z), \tag{3.19}$$

as long as $z < z_{\rm eq}$, calling $z_{\rm eq}$ the moment at which N_1 reaches the equilibrium distribution. Following eq. (3.18), and adopting the expression of the Hubble rate in the radiation-dominated epoch [80]:

$$H(z) = 1.66 g_*^{1/2} \frac{M_1^2}{M_{\rm Pl}\, z^2}, \tag{3.20}$$

with $M_{\rm Pl}$ being the Planck mass, we can easily find a solution:

$$N_{N_1}(z < z_{\rm eq}) = \frac{1}{4} \frac{\zeta(3) G_F^2\, M_{\rm Pl} M_1^3}{1.66 \pi^2 \sqrt{g_*}} \left(\frac{m_W}{v_R} \right)^4 \left(\frac{1}{z_{\rm RH}^3} - \frac{1}{z^3} \right), \tag{3.21}$$

where we used the expression of $S(z)$:

$$S(z) = \frac{3}{2} \frac{\zeta(3) G_F^2\, M_{\rm Pl} M_1^3}{1.66 \pi^2 \sqrt{g_*}} \left(\frac{m_W}{v_R} \right)^4 z^{-4}. \tag{3.22}$$

Given the expression of $S(z)$, it is clear that the scattering rate quickly decreases with the temperature, therefore we may expect an intermediate value $\bar{z}$, such that $z_{\rm RH} < \bar{z} < z_{\rm eq}$ at which $S(z)$ becomes negligible and the abundance of N_1 *freezes-in*, without reaching its equilibrium value. In this way, the current number density of N_1 is given by:

$$N_{N_1}^0 \equiv N_{N_1}(\bar{z}) \simeq \frac{1}{4} \frac{\zeta(3) G_F^2\, M_{\rm Pl} M_1^3}{1.66 \pi^2 \sqrt{g_*}} \left(\frac{m_W}{v_R} \right)^4 \frac{1}{z_{\rm RH}^3}. \tag{3.23}$$

In the left panel of figure 1 we show the ratio $\Gamma_S(z)/H(z)$, that measures the efficiency of the scattering reaction. As can be seen from the figure, this ratio is well below one even at the reheating temperature $T_{\rm RH}$. The abundance of N_1 as a function of temperature T is shown in the right panel of figure 1.

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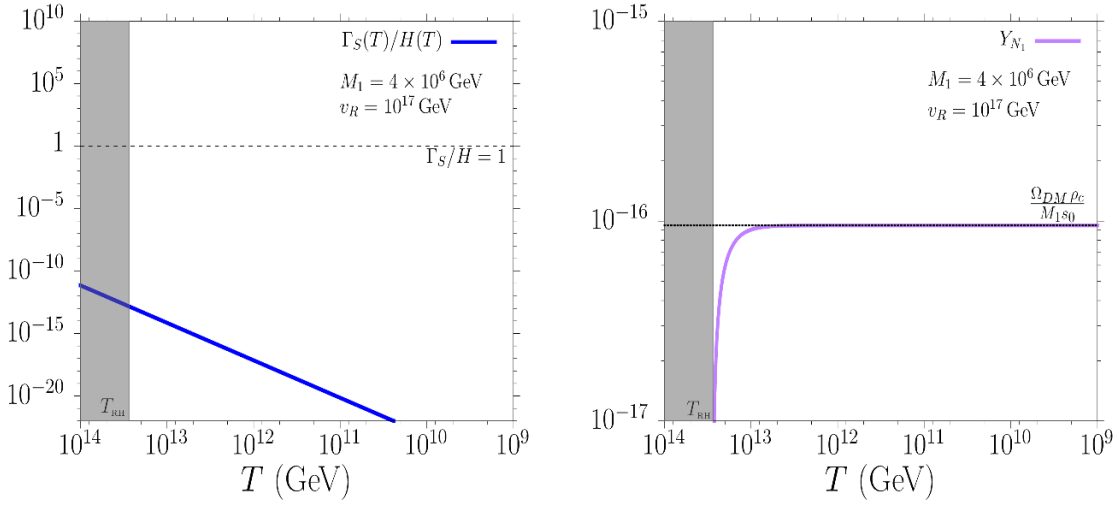

Figure 1. Left panel: evolution in temperature of the ratio Γ_S/H between scattering rate, eq. (3.13), and the Hubble parameter, eq. (3.20). Right panel: evolution of N_1 abundance as a function of temperature. We consider $M_1 = 4\times10^6$ GeV, $T_{\rm RH} = 2.7\times10^{13}$ GeV and $v_R = 10^{17}$ GeV.

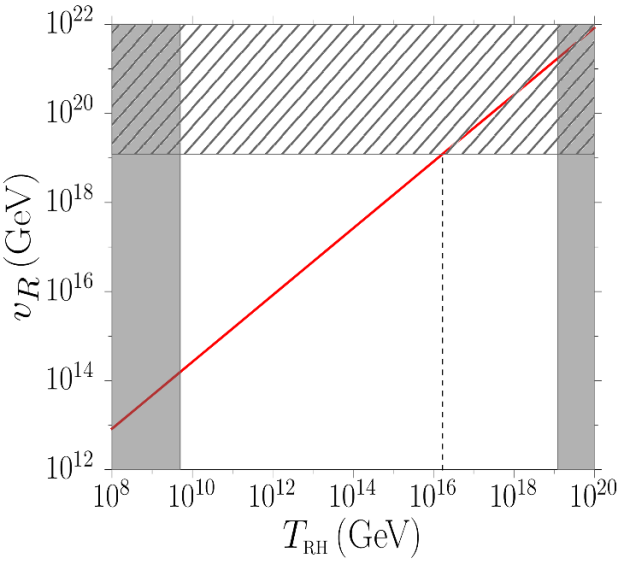

Figure 2. Relation between v_R and $T_{\rm RH}$ that give the correct final abundance $Y^0_{N_1} = Y^0_{\rm DM}$. The vertical shadowed region on the right corresponds to temperatures above the Planck scale. The upper hatched region is cut out if transplanckian values of v_R are excluded. The vertical shaded region on the left shows the lower bound given by leptogenesis, see the discussion in section 4.2 and in particular eq. (4.13).

This model allows for the production of a relatively small abundance of RH neutrinos, in particular of N_1, exploiting the freeze-in mechanism. The scattering processes mediated by $\mathrm{SU}(2)_R$ gauge bosons never become efficient after inflation, since their freeze-out temperature is always higher than $T_{\rm RH}$. Thus, after reheating, the abundance of N_1 can only reach a low final value, which is then preserved due to the absence of any other interaction.

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From eq. (3.11), we can obtain a relation between v_R and the reheating temperature $T_{\rm RH}$ such that the final abundance of N_1 is equal to the current DM abundance. We have:

$$Y^0_{N_1} \simeq \frac{1}{4\times 2.40 g_*}\frac{\zeta(3) G_F^2 M_{\rm Pl}}{1.66\pi^2\sqrt{g_*}}\left(\frac{m_W}{v_R}\right)^4 T^3_{\rm RH}, \tag{3.24}$$

and imposing the value in eq. (3.11) we get:

$$v_R = \left[\frac{1}{4\times 2.40 g_*}\frac{\zeta(3) G_F^2 M_{\rm Pl}}{1.66\pi^2 g_*^{1/2} Y^0_{\rm DM}}\right]^{\frac{1}{4}} m_W\, T^{3/4}_{\rm RH} \simeq 8.3\times 10^6\, T^{3/4}_{\rm RH}, \tag{3.25}$$

This relation is a prediction of our model, and it is shown in figure 2.

4 The baryon asymmetry

In our model the other two heavy neutrinos N_2 and N_3 are responsible for the production of the correct amount of baryon asymmetry of the Universe. Since we assume that Δ_L does not contribute to the light neutrino masses and that $M_{\Delta_L} \gg M_i$, our model reduces to ordinary RH neutrino leptogenesis, without any effect from the triplet [62]. For a discussion of the situations in which the scalar triplet could affect leptogenesis we refer to refs. [62, 81].

Given the constraint on N_1's Yukawa couplings in eq. (3.4), that completely decouples the lightest neutrino N_1, and the fact that $M_1 \ll 10^9\,$GeV [82], the asymmetry must be produced by N_2. Since N_1 does not play any role here, we can already expect on M_2 the same lower bound that applies on M_1 in N_1-dominated leptogenesis, i.e. $M_2 \gtrsim 5\times 10^8\,$GeV [82, 83]. For definiteness, we will always consider hierarchical leptogenesis, obtained by imposing $M_3 \geq 3M_2$.[4] Assuming $M_2 \gtrsim 5\times 10^8\,$GeV, the asymmetry can be produced in two different regimes according to the mass of N_2 [85–87]: for $5\times 10^8\,\text{GeV} \lesssim M_2 \lesssim 5\times 10^{11}\,$GeV leptogenesis will take place in a two fully-flavoured regime, while for $M_2 \gtrsim 5\times 10^{11}\,$GeV leptogenesis will be unflavoured.

4.1 Asymmetry production

Case I: $5\times 10^8\,\text{GeV} \lesssim M_2 \lesssim 5\times 10^{11}\,\text{GeV}$: at temperatures $T \lesssim 5\times 10^{11}\,$GeV the charged τ-Yukawa interactions are in equilibrium and more efficient than the LH-RH neutrino interactions [83, 88]. Therefore, the relevant quantities will be the asymmetries $\Delta_\tau \equiv B/3 - L_\tau$ and $\Delta_{\tau_2^\perp} \equiv B/3 - L_{\tau_2^\perp}$, where $\tau_2^\perp$ defines the flavour component produced by N_2, orthogonal to the τ flavour direction. Neglecting[5] flavour coupling [61, 86, 87, 90–93], scattering terms [79, 83, 94], thermal effects [95] and quantum corrections [89, 96, 97], the $B-L$ asymmetry produced by N_2, $N^{\rm lep,2}_{B-L}$, is given by the sum of the produced Δ_τ and $\Delta_{\tau_2^\perp}$ asymmetries as [87]

$$\begin{aligned} N^{\rm lep,2}_{B-L} &= N^{\rm lep,2}_{\Delta_\tau} + N^{\rm lep,2}_{\Delta_{\tau_2^\perp}} \\ &\simeq \varepsilon_{2\tau}\kappa_f(K_{2\tau}) + \varepsilon_{2\tau^\perp}\kappa_f(K_{2\tau_2^\perp}). \end{aligned} \tag{4.1}$$

[4] Another possibility is provided by resonant leptogenesis (see e.g. ref. [84]), in which at least two heavy neutrinos have a mass splitting comparable to their decay widths. In our situation, this could be applied to N_2 and N_3.

[5] The impacts of the corrections we neglect have been estimated not to be larger that about 20% [83, 89].

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Here, $\varepsilon_{2\tau^\perp} = \varepsilon_{2e} + \varepsilon_{2\mu}$, where $\varepsilon_{2\alpha}$ ($\alpha = e,\, \mu,\, \tau$) are the N_2 *CP* asymmetries in flavour α defined as:

$$\varepsilon_{2\alpha} \equiv -\frac{\Gamma_{2\alpha} - \overline{\Gamma}_{2\alpha}}{\Gamma_2 + \overline{\Gamma}_2}, \tag{4.2}$$

where $\Gamma_{2\alpha}$ and $\overline{\Gamma}_{2\alpha}$ are the rates of the decays $N_2 \to l_\alpha H$ and $N_2 \to \bar{l}_\alpha H$ respectively and $\Gamma_2 = \sum_\alpha \Gamma_{2\alpha}$, $\overline{\Gamma}_2 = \sum_\alpha \overline{\Gamma}_{2\alpha}$. We also introduced the flavoured decay parameters:

$$\begin{aligned} K_{i\alpha} &\equiv \frac{\Gamma_{i\alpha} + \overline{\Gamma}_{i\alpha}}{H(T = M_i)} \\ &= \frac{|Y^\nu_{\alpha i}|^2\, v^2}{m_* M_i} = \frac{1}{m_*} \left| \sum_j \sqrt{m_j}\, U_{\alpha j} \Omega_{ji} \right|, \end{aligned} \tag{4.3}$$

where m_* is the equilibrium neutrino mass [79], defined as

$$m_* = \frac{8\pi v^2}{H(T = M_i)} = \frac{16\pi^{5/2} g_*^{1/2}}{3\sqrt{5}} \frac{v^2}{M_{\rm Pl}} \simeq 1.08 \times 10^{-3}\,{\rm eV}. \tag{4.4}$$

In eq. (4.1) we have $K_{2\tau_2^\perp} = K_{2e} + K_{2\mu}$, while the $\kappa_f(K_{2\alpha})$ are the efficiency factors, whose expressions can be found in the literature [79, 83, 98].

Case II: $M_2 > 5 \times 10^{11}$ GeV: in this case, the flavour interactions are not in equilibrium, therefore the coherence of the lepton and anti-lepton states produced by N_2 is not broken. Therefore, the evolution of the full $B-L$ asymmetry and N_2's abundance is tracked by the Boltzmann equations, whose solution gives

$$N^{{\rm lep},2}_{B-L} \simeq \varepsilon_2 \kappa_f(K_2), \tag{4.5}$$

where $K_2 = \sum_\alpha K_{2\alpha}$ is the total N_2's decay parameter.

We now have an expression for the $B - L$ asymmetry produced by N_2's dynamics, both for Case I and Case II. In order to obtain the final value of the asymmetry we still have to take into account the impact of the processes involving N_1.

4.2 Expressions for the final asymmetry

Below $T \sim M_2$, the asymmetry stays constant. However, for temperatures $T \lesssim 5 \times 10^8$ GeV the μ Yukawa interactions are in equilibrium.

In Case I, this implies that at temperatures $M_1 < T' \lesssim 5 \times 10^8$ GeV the asymmetry $N^{{\rm lep},2}_{\Delta_{\tau_2^\perp}}$ gets projected onto the e and μ flavour directions. Neglecting phantom terms [85, 87], we obtain the asymmetries in $\Delta_\delta \equiv B/3 - L_\delta$ ($\delta = e,\, \mu$), at temperature T', simply by

$$N_{\Delta_e}(T') = \frac{K_{2e}}{K_{2\tau_2^\perp}} N^{{\rm lep},2}_{\Delta_{\tau_2^\perp}}, \qquad N_{\Delta_\mu}(T') = \frac{K_{2\mu}}{K_{2\tau_2^\perp}} N^{{\rm lep},2}_{\Delta_{\tau_2^\perp}}. \tag{4.6}$$

The action of N_1 will then take place along the three flavour directions e, μ, τ. Considering that the asymmetry produced by N_1 can be safely neglected, neglecting again flavour

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coupling and taking as initial conditions N_{Δ_e}, N_{Δ_μ} in eq. (4.6) and N_{Δ_τ} in eq. (4.1), the Boltzmann equations can be solved, for Case I, giving the total final asymmetry as the sum of the final asymmetries in Δ_α, $\alpha = e,\, \mu,\, \tau$ [71, 83, 85, 87, 98]:

$$\begin{aligned} N_{B-L}^{\rm lep,f} &= \sum_{\alpha=e,\,\mu,\,\tau} N_{\Delta_\alpha}^{\rm lep,f} \\ &\simeq \frac{K_{2e}}{K_{2\tau_2^\perp}} \varepsilon_{2\tau^\perp} \kappa_f(K_{2\tau_2^\perp}) e^{-\frac{3\pi}{8} K_{1e}} + \frac{K_{2\mu}}{K_{2\tau_2^\perp}} \varepsilon_{2\tau^\perp} \kappa_f(K_{2\tau_2^\perp}) e^{-\frac{3\pi}{8} K_{1\mu}} \\ &\quad + \varepsilon_{2\tau} \kappa_f(K_{2\tau}) e^{-\frac{3\pi}{8} K_{1\tau}}. \end{aligned} \tag{4.7}$$

As for Case II, the asymmetry in eq. (4.5) gets projected as:

$$N_{\Delta_e}(T') = \frac{K_{2e}}{K_2} N_{B-L}^{\rm lep,2}, \qquad N_{\Delta_\mu}(T') = \frac{K_{2\mu}}{K_2} N_{B-L}^{\rm lep,2}, \qquad N_{\Delta_\tau}(T') = \frac{K_{2\tau}}{K_2} N_{B-L}^{\rm lep,2}, \tag{4.8}$$

and, similarly, the final asymmetry is given by

$$\begin{aligned} N_{B-L}^{\rm lep,f} &= \sum_{\alpha=e,\,\mu,\,\tau} N_{\Delta_\alpha}^{\rm lep,f} \\ &\simeq \frac{K_{2e}}{K_2} \varepsilon_2 \kappa_f(K_2) e^{-\frac{3\pi}{8} K_{1e}} + \frac{K_{2\mu}}{K_2} \varepsilon_2 \kappa_f(K_2) e^{-\frac{3\pi}{8} K_{1\mu}} \\ &\quad + \frac{K_{2\tau}}{K_2} \varepsilon_2 \kappa_f(K_2) e^{-\frac{3\pi}{8} K_{1\tau}}. \end{aligned} \tag{4.9}$$

As already pointed out, in our specific model, N_1's Yukawa couplings are suppressed. Therefore its washout is negligible and the final asymmetry expressions can be simplified. Indeed, from eqs. (4.3) and (3.4) we have $K_{1\alpha} \simeq 0$. For this reason, from eq. (4.7), we obtain for Case I:

$$N_{B-L}^{\rm lep,f} \simeq \varepsilon_{2\tau^\perp} \kappa_f(K_{2\tau_2^\perp}) + \varepsilon_{2\tau} \kappa_f(K_{2\tau}), \tag{4.10}$$

while for Case II, from eq. (4.9), we get:

$$N_{B-L}^{\rm lep,f} \simeq \varepsilon_2 \kappa_f(K_2). \tag{4.11}$$

We notice that, due to the negligible washout by N_1, the phantom terms would anyway cancel out, therefore these final expressions are not affected by this correction.

As an example, in figure 3 we consider the case with $M_2 = 10^{11}$ GeV and a particular choice of the matrix Ω in eq. (3.8). In the left panel, we show the ratios $\Gamma_{D2}(z)/H(z)$ and $\Gamma_{ID2}(z)/H(z)$. These ratios reach a value equal to two at a temperatures between $10^{11} - 10^{12}$ GeV. In the right panel, we show the evolution of N_2 and $B-L$ asymmetry abundances. As can be seen, the N_2's abundance (red line) can reach and track the equilibrium distribution (dashed line), as in the strong-washout regime [79, 87]. The blue line marks the evolution of the asymmetry abundance in modulus. The asymmetry produced at early stages, $T > M_2$, is quickly erased at around $T \simeq M_2$ when decay processes enter equilibrium. Successively, a new asymmetry is produced with opposite sign and, due to the departure from equilibrium of inverse decay processes, it freezes out to a final value above the experimental bound $Y_{B-L}^{\rm CMB}$. Therefore, we can conclude that it is possible, in

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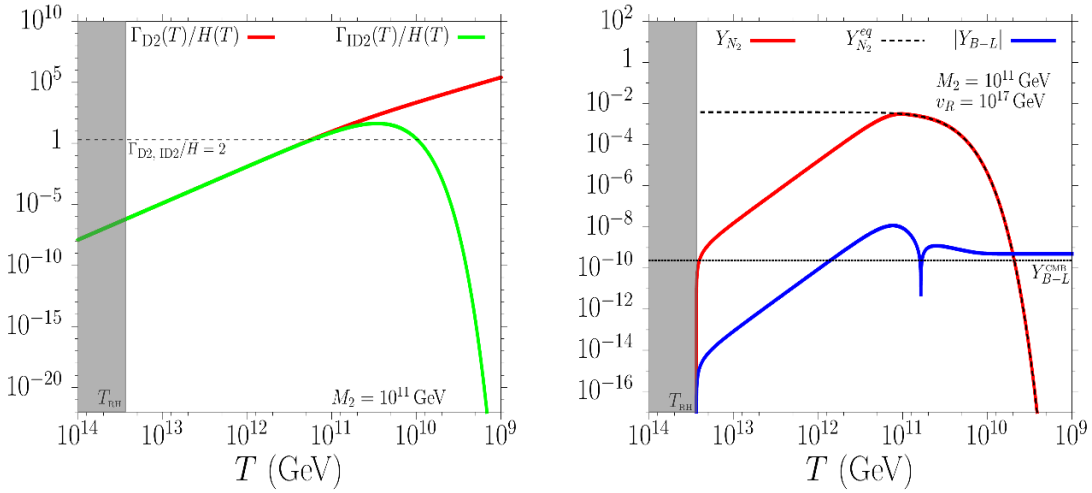

Figure 3. Left panel: evolution in temperature of the ratio between N_2's decay and inverse decay rates and the Hubble parameter. Right panel: evolution of N_2 and $B-L$ asymmetry abundances as a function of temperature. We consider $M_2 = 10^{11}$ GeV, $T_{\rm RH} = 2.7\times10^{13}$ GeV and $v_R = 10^{17}$ GeV. The final $B-L$ asymmetry abundance is obtained as $Y^{\rm CMB}_{B-L} = \eta^{\rm CMB}_B/(2.40 g_*\, 0.96\times10^{-2})$, accounting for the sphaleron conversion rate and the dilution factor.

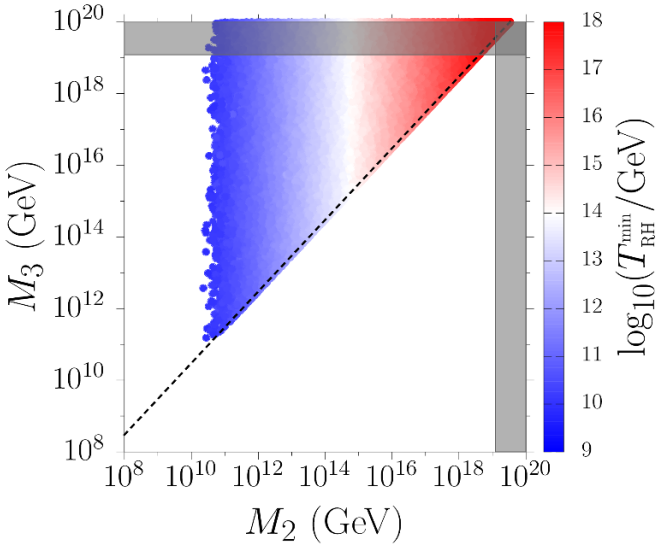

Figure 4. Scatter plot of points in the plane M_2-M_3 that give successful leptogenesis. The colour gives the corresponding $\log_{10}(T^{\rm min}_{\rm RH}/{\rm GeV})$. The dashed line shows the hierarchical limit $M_3 = 3M_2$. The shaded regions exclude masses $M_i \geq M_{\rm Pl}$.

our model, to obtain successful leptogenesis from N_2's decays. Note that in the case shown in the figure, the asymmetry is produced in the two-flavour regime, i.e. Case I eq. (4.10).

In figure 4 a general scan on M_2 and M_3, with $M_3 \geq 3M_2$ is performed and Case I and Case II (eqs. (4.10) and (4.11) respectively) are adopted according to the value of M_2. Here we assume Normal Ordering (NO) of the light neutrino masses $m_2^2 = m_1^2 + m_{\rm sol}^2$,

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$m_3^2 = m_1^2 + m_{\rm atm}^2$ with $m_{\rm sol} = 0.0087\,{\rm eV}$ and $m_{\rm atm} = 0.0496\,{\rm eV}$ [99].[6] According to our previous discussion, we set $m_1 = 0$. We scan on the neutrino mixing angles uniformly extracting them in their experimental 3σ ranges as in [99], while the Dirac and Majorana phases are extracted on their full variability range. We also scan on the complex angle θ in Ω, while setting $\alpha = \beta = 0$ for simplicity. Moreover, we require $|\Omega_{ij}|^2 \leq 2$.

We assume here an instantaneous transition from the two fully-flavoured to the unflavoured regime, at $M_2 = 5 \times 10^{11}\,{\rm GeV}$. A more accurate description should employ a full density matrix formalism [85] to describe leptogenesis in the transition region, however only a small impact is expected. It is also possible to find a lower bound on M_2 as we expected from considerations similar to N_1-dominated leptogenesis. In this case the bound appears to be slightly higher: $M_2 \gtrsim 10^{10}\,{\rm GeV}$.

In figure 4, the colours encode the minimal reheating temperature $T_{\rm RH}^{\rm min}$ needed to produce the correct final asymmetry for each value of $(M_2,\, M_3)$. This is obtained as [79]:

$$T_{\rm RH}^{\rm min} \simeq \frac{M_2}{z_L(\overline{K}_2) - 2e^{-3/\overline{K}_2}}, \tag{4.12}$$

where $\overline{K}_2 = K_2$ in Case II, while $\overline{K}_2 = K_{2\tau}, K_{2\tau_2^\perp}$ in Case I, if the asymmetry in the τ or $\tau_2^\perp$ flavour dominates respectively. From figure 4 it is possible to notice that the lowest values are obtained for M_2 around the lower bound. We obtained a lower bound entirely given by leptogenesis:

$$T_{\rm RH}^{\rm min} \gtrsim 7 \times 10^9\,{\rm GeV}. \tag{4.13}$$

Using this lower bound on $T_{\rm RH}^{\rm min}$ and eq. (3.25), we can find a range of allowed values of v_R, see also figure 2. We have $2 \times 10^{14}\,{\rm GeV} \lesssim v_R \lesssim 2 \times 10^{21}\,{\rm GeV}$, or the more restrictive case $2 \times 10^{14}\,{\rm GeV} \lesssim v_R \lesssim 10^{19}\,{\rm GeV}$ if we require $v_R < M_{\rm Pl}$. We notice that the lower bound on v_R agrees with our assumptions of a very high symmetry breaking scale.

5 Conclusions

In this paper, we have considered a left-right symmetric model, where a DM particle is produced through a freeze-in process. The model is able to produce the correct final abundance of DM and baryon asymmetry, while at the same time the DM candidate is suitably heavy, $M_1 = 4\,{\rm PeV}$, and long-lived, $\tau_{N_1} \simeq 10^{28}\,{\rm s}$, to produce high-energy neutrinos consistent with the IceCube signal. It is interesting to point out that this can be realised only with vanishing absolute neutrino mass scale $m_1 \simeq 0$, which allows for a long lifetime τ_{N_1} through tiny values of N_1's Yukawa couplings. As a consequence, this model predicts a power-law relation between the reheating temperature of the Universe, $T_{\rm RH}$, and the vacuum expectation value of the $\mathrm{SU}(2)_R$ triplet. To obtain successful leptogenesis, a lower bound $T_{\rm RH} \gtrsim 7 \times 10^9\,{\rm GeV}$ should be satisfied.

In conclusion, the model presented in this paper successfully provides a consistent solution to both the DM problem and the generation of the matter-antimatter asymmetry of the Universe through leptogenesis, while producing at the same time a viable interpretation

[6] See also ref. [100] for a recent review on neutrino oscillation parameters.

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of the highest energy IceCube neutrino events. The basis is a left-right symmetric model, where the right-handed neutrino fields are all involved and mutually necessary in the generation of the different mechanisms at hand (correct DM relic abundance, baryon asymmetry and IceCube neutrino flux). The results can therefore be accommodated naturally in the LRSM scheme.

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