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TeV scale resonant leptogenesis from supersymmetry breaking

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Abstract: We propose a model of TeV-scale resonant leptogenesis based upon recent models of the generation of light neutrino masses from supersymmetry-breaking effects with TeV-scale right-handed (rhd) neutrinos, N_i. The model leads to large cosmological lepton asymmetries via the resonant behaviour of the one-loop self-energy contribution to N_i decay. Our model addresses the primary problems of previous phenomenological studies of low-energy leptogenesis: a rational for TeV-scale rhd neutrinos with small Yukawa couplings so that the out-of equilibrium condition for N_i decay is satisfied; the origin of the tiny, but non-zero mass splitting required between at least two N_i masses; and the necessary non-trivial breaking of flavour symmetries in the rhd neutrino sector. The low mass-scale of the rhd neutrinos and their superpartners, and the TeV-scale A-terms automatically contained within the model offer opportunities for partial direct experimental tests of this leptogenesis mechanism at future colliders.

Keywords: Supersymmetry Breaking, Supersymmetry Phenomenology, Baryogenesis, Neutrino Physics.

Contents

1. Introduction

The seesaw mechanism [1] and the associated mechanism of leptogenesis [2] are very attractive means to explain the origin of the small neutrino masses and the baryon asymmetry of the universe. However they are very hard to directly test, and at present we have at best only weak circumstantial evidence that they are correct. The primary difficulty is that the new particles that they involve have masses far beyond experimental reach. In particular, in the standard realisation of the seesaw, the current lower bound for successful leptogenesis on the mass of the lightest right-handed neutrino N_1 is $M_{N_1} > 5 \cdot 10^8$ GeV [3]–[8].[1] Moreover the seesaw models contain many more parameters than there are low energy observables which could constrain them. For the standard seesaw-extended standard model with three right-handed neutrinos there are 18 parameters to be compared with 7 observables in the light neutrino mass matrix. (However, in its supersymmetric version with the non-trivial assumption of soft terms universality, rare lepton flavor changing and/or CP violating processes can give access to more parameters [13, 14].)

Given this lack of direct evidence in favour of the standard mechanisms, it is important to consider possible low-energy, TeV-scale alternatives to, or variations of the standard seesaw leptogenesis mechanism, especially if they are testable. Additionally, such a mechanism could have the advantage of not requiring assumptions about the thermal history of the universe up to energy scales as high as 10^{10} GeV, as is usually the case, far above

[1]This bound holds for hierarchical rhd neutrinos assumed to be in thermal equilibrium before decaying. It can be relaxed to $M_{N_1} > 2 \cdot 10^7$ GeV if the lightest rhd neutrino is assumed to be the dominant species populating the universe at the end of inflation from the decay of the inflaton [7]. It can also be relaxed if the spectrum of rhd neutrinos is not highly hierarchical [8], or if they are quasi-degenerate in mass [6] and [8]–[12].

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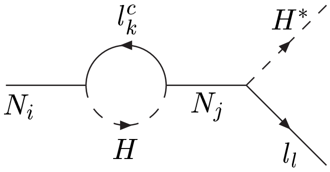

Figure 1: Self-energy diagram for the right-handed neutrino decay.

the temperature of the last epoch that has been tested, the epoch of nucleosynthesis at the MeV scale. Moreover this would allow one to avoid the potential problems of the creation of dangerous relics, such as the gravitino, in too large a number at higher temperature.

To build such a low-energy leptogenesis model is however a difficult task, essentially for the following reasons (for more details see [6]). First, low scale seesaw neutrino masses require tiny couplings and therefore generically induce too small a CP asymmetry. Second, low scale means small Hubble constant which also requires tiny couplings in order that the decay of the particle at the origin of the asymmetry is not in thermal equilibrium. Third, a small Hubble constant requires in addition that the various scatterings which can suppress the asymmetry be under control. In particular, at such a low scale, the very fast gauge scatterings strongly prefer that the decaying particle be a singlet of all low energy gauge symmetries. To avoid these problems, and limiting ourselves to standard thermal leptogenesis, one can think about three possibilities with decaying singlet particles: a large degeneracy of masses between the decaying particles [6] and [8]–[12]; a hierarchy between the couplings of real and virtual heavy particles in the one loop leptogenesis diagrams (see [6]); or three body decays of the heavy particles with suppressed two body decays [6].

In this letter we will consider the first possibility that two or more particles are quasi degenerate.[2] In this case the asymmetry can be significantly enhanced through a resonant behaviour of the propagator of the virtual heavy particle in the leptogenesis self-energy diagram, figure 1. As far as we are aware this is the only mechanism which can work at low scale within the standard leptogenesis model. However this framework still suffers from various significant difficulties: 1) As already mentioned above, at a scale as low as $1 - 10\,\text{TeV}$, neutrino mass constraints and the out-of equilibrium condition on the decay width require tiny Yukawa couplings of order $\sim 10^{-6} - 10^{-7}$, and such small couplings need explanation. 2) In order to compensate the large suppression of the asymmetry induced by these tiny couplings, an extremely tiny mass splitting is required between two right-handed neutrino masses. The degree of degeneracy required has to be of order the value of the decay width to mass ratio, which implies that $(M_{N_1} - M_{N_2})/(M_{N_1} + M_{N_2}) < 10^{-10}$ [12]. Such a tiny splitting, if not physically motivated, can be considered as a fine-tuning. 3) In a generic seesaw model there is no explanation why the right-handed neutrinos would have such a small mass ($M_N \sim$ TeV), just around the present experimentally reachable mass scale. 4) Finally, the tiny Yukawa couplings imply that the right-handed neutrino production cross

[2] Also in the context of broken supersymmetric theories, ref. [15] considers an alternative TeV-scale leptogenesis mechanism based upon hierarchical soft A-terms.

sections are very suppressed, the right-handed neutrinos are not observable, and the model is not testable even if is located at a low scale. Here we construct a resonant leptogenesis model which possesses a natural explanation of the first three problems and is in addition partially testable at future collider facilities such as the LHC. It arises in the framework of the MSSM with L-violating soft supersymmetry breaking.

Finally we note that recently a phenomenological Froggat-Nielsen type model has been proposed in [12] which gives splittings naturally of order the decay width and therefore leads naturally to a large asymmetry too. Our model differs considerably from the one studied in ref. [12].

2. The model

The most attractive suggested solution to the μ problem of supersymmetric theories is to invoke a symmetry to forbid the supersymmetry-preserving $\mu H_u H_d$ superpotential term; the presence of the μ term in the low-energy MSSM effective lagrangian is explained as the result of an intermediate scale (m_I) supersymmetry breaking expectation value within a higher-dimensional, $1/M_{\rm Planck}$-suppressed operator [16]. It has been recently emphasized [17, 19] that such a symmetry can also suppress the masses and interactions of the right-handed neutrinos in the same way (for related work see [20]). This mechanism leads to TeV scale right-handed neutrinos with tiny Yukawa couplings which, in addition naturally lead to phenomenologically successful light Majorana neutrino masses [17, 19], and an explicit and simple realization of this idea has been recently constructed [19].

The mechanism of ref. [19] is based on the existence of two heavy standard model singlet fields: X_{ij} which carries a flavour structure ($i, j = 1, 2, 3$), and, Y, which is flavor blind.[3] In this model the various particles have the following R-charge pattern: X and Y have charges 4/3, the N_i have charge 2/3, the right-handed charged lepton superfields E_i have charge 2 and H_u, H_d and L_i all have R-charge equal to 0, and the usual R-parity is assumed, with in addition $R_p(X) = R_p(Y) = +1$. This leads to the following allowed N_i-dependent interactions in the superpotential

$$\mathcal{L}_N^W = \int d^2\theta \left(g \frac{X_{ij}}{M_P} L_i N_j H_u + g' \frac{Y}{M_P} L_i N_i H_u + \cdots \right), \tag{2.1}$$

while the set of Kähler terms involving the right-handed N_i fields are

$$\mathcal{L}_N^K = \int d^4\theta \left(h \frac{Y^\dagger}{M_P} N_i N_i + h_B \frac{Y^\dagger Y X_{ij}^\dagger}{M_P^3} N_i N_j + h_x \frac{X_{ik} X_{kj}^\dagger Y^\dagger}{M_P^3} N_i N_j + \cdots \right), \tag{2.2}$$

with $M_P = 1/\sqrt{8\pi G_N} \sim 2 \cdot 10^{18}$ GeV the reduced Planck mass. We will take both types of dimensionless couplings (g's and h's) to be of order one. In addition, we assume that after supersymmetry breaking in the hidden sector, the field $Y(X)$ acquires a $F(A)$ component vacuum expectation value but no $A(F)$ component:

$$\langle Y \rangle_F = F_Y = f_Y m_I^2 , \qquad \langle X_{ij} \rangle_A = A_{X_{ij}} = a_{X_{ij}} m_I , \qquad \langle Y \rangle_F = \langle X_{ij} \rangle_A = 0 . \tag{2.3}$$

[3] For simplicity in this paper we simply assume flavour-blind Y couplings. This can result from a flavour symmetry. We leave the discussion of the various symmetry patterns to a future publication [21].

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This then results in the following interactions of the softly broken right-handed neutrino-extended-MSSM (see e.g. [22, 23]):

$$\mathcal{L} = \int d^2\theta \, (\lambda_{ij} L_i N_j H_u + M_N N_i N_i + \Delta M_{Nij} N_i N_i) + A \tilde{L}_i \tilde{n}_i h_u + B_{ij}^2 \tilde{n}_i \tilde{n}_j + \cdots \,, \tag{2.4}$$

where $\tilde{n}_i$ are the rhd sneutrino fields, $\tilde{L}_i$ is the lhd slepton doublet, h_u is the up-type Higgs scalar doublet, and the omitted terms include the usual soft scalar mass terms. In eq. (2.4) we have from eqs. (2.1)–(2.3)

$$\begin{aligned}
\lambda_{ij} &= g a_{Xij} \left(\frac{m_{3/2}}{M_P}\right)^{1/2} && \sim 10^{-7} - 10^{-8}, \\
M_N &= h f_Y m_{3/2} && \sim 1\,\text{TeV}\,, \\
\Delta M_{Nij} &= h_x a_{Xik} a^{\dagger}_{Xkj} \frac{f_Y m_{3/2}^2}{M_P} && \sim 10^{-3}\,\text{eV}\,, \\
A &= g' f_Y m_{3/2} && \sim 1\,\text{TeV}\,, \\
B_{ij}^2 &= h_B f_Y^2 a^s_{Xij} \left(\frac{m_{3/2}^5}{M_P}\right)^{1/2} && \sim 10^{-3}\,\text{GeV}^2\,,
\end{aligned} \tag{2.5}$$

where the relations hold assuming g, h, a_X and f_Y to be of order unity with $m_{3/2} \simeq 1\,\text{TeV}$, that is to say taking $m_I = \sqrt{m_{3/2} M_P} \sim 10^{11}$ GeV as we expect in hidden-sector models.[4]

In this lagrangian since the Yukawa couplings λ are of order 10^{-7}–10^{-8} for M_N of order TeV, the seesaw induced neutrino masses will be in general of order $m_\nu \sim \lambda^2 v^2 / M_N \sim 10^{-5}$-$10^{-3}$ eV. As argued in ref. [19] this could eventually explain the solar data but is to small to explain the atmospheric data. However at the one loop level the B term together with the large A terms and gauge interactions induce neutrino masses naturally of order the atmospheric lower bound on neutrino masses for the heaviest neutrino ν_3, $m_{\nu_3} > \sqrt{\delta m^2_{\text{atm}}} \sim 0.05\,\text{eV}$. This one loop contribution to the neutrino masses is of size

$$m_\nu^{\text{loop}} \sim \frac{\alpha_w}{96\pi} \frac{m_I^9 v^2}{M^5 m_{\text{susy}}^5} \simeq 10^{-2}\text{eV} - 10^{-1}\,\text{eV}\,, \tag{2.6}$$

and has flavour structure set by the lepton-number violating B-term mass, B_{ij}^2, for the sneutrinos as given in eq. (2.5), see refs. [17, 19]. Since the resulting light neutrino mass matrix is of exactly the same symmetric Majorana structure as in usual supersymmetric see-saw models, the counting of the Dirac and Majorana phases in m_ν^{loop} is identical.

[4] Note that there is no automatic vacuum stability problem implied by the A-terms. The metastability condition [18] requires that $|A|^2$ is smaller than the sum of lhd and rhd sneutrino and up-like Higgs soft mass-squareds, and this can easily be accommodated in our model by a modest suppression of A. Successful resonant leptogenesis puts no lower bound on A, while the required neutrino masses can be accommodated by, eg, Suppressing A and enhancing B by factors of O(few) from the order-of-magnitude values given in eq. (2.5). The probability during the lifetime of the Universe of a single tunneling event to the true vacuum can be estimated to be P $\sim H^{-4} M_{3/2}^4 e^{-S_4}$, where $S_4 \sim 46(M_{3/2}/A)^2$, see ref. [24, section 5.3]. For reasonable values of the ratio $M_{3/2}/A$ this probability is vanishingly tiny.

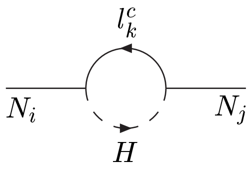

Figure 2: Self-energy contribution to eq. (2.11).

Note that since the Y field is flavour blind the spectrum of right-handed neutrinos is degenerate at leading order, with only tiny splittings of order $10^{-15} - 10^{-17}$ generated by the higher-order term $X_{ik}X^{\dagger}_{kj}Y^{\dagger}N_iN_j$. This term comes from two contributions at the same order in $1/M_P$: the tree level h_x contribution of eq. (2.2) and the one-loop contribution of figure 2 induced by the Yukawa couplings, λ_{ij}. These contributions give[5]

$$M_{Nij}^{\rm corrected} = M_N\left[\mathbb{I}_{ij} + \beta(a_X a_X^{\dagger} + a_X^* a_X^T)_{ij}\right] \tag{2.7}$$

$$\text{with}\;\; \beta \sim \frac{m_{3/2}}{hM_P}\left(h_x + \frac{g^2}{16\pi^2}\log\frac{M_P}{M_N}\right) \sim 10^{-15}\,. \tag{2.8}$$

It is also important to stress that if there exists any other hidden sector flavour non-singlet field, Z_{ij}, with symmetry properties different from X_{ij}, then irrespective of these symmetry properties, the additional Kähler term

$$\mathcal{L} \ni \int d^4\theta \frac{1}{M_{\rm P}^3} h_z Z_{ik} Z^{\dagger}_{kj} Y^{\dagger} N_i N_j \,+\, \mathcal{O}\left(\frac{1}{M^4_{\rm Planck}}\right) \tag{2.9}$$

cannot be forbidden. The addition of such a term does not in any way disturb the nice properties of the above model for neutrino masses. However, if Z_{ij} (similarly to X_{ij}) has no F-term but a A term, $\langle Z_{ij}\rangle_A = a_{Z_{ij}} m_I$, an additional low-energy effective interaction emerges:

$$\mathcal{L} \ni \frac{m_{3/2}^2}{M_P} h_z a_{Z_{ik}} a^{\dagger}_{Z_{kj}} f_Y^{\dagger} \int d^2\theta N_i N_j + \mathcal{O}\left(\frac{1}{M_P^2}\right). \tag{2.10}$$

This results in a total rhd neutrino mass matrix of the form

$$M_N^R = M_N\left[\mathbb{I} + \beta(a_X a_X^{\dagger} + a_X^* a_X^T) + \gamma(a_Z a_Z^{\dagger} + a_Z^* a_Z^T)\right] \tag{2.11}$$

with M_N and β given by eqs. (2.5) and (2.8), and

$$\gamma = \frac{h_z f_Y}{h}\frac{m_{3/2}}{M_P} \sim 10^{-15}\,. \tag{2.12}$$

The β- and γ-dependent terms are irrelevant for neutrino masses, but as we will show in the next section can be relevant for leptogenesis.

[5]For simplicity, the RG resummation of the logarithms in eq. (2.8) have not been performed. For an example of such a procedure in the context of resonant leptogenesis, see ref. [27].

3. Resonant leptogenesis

From the lagrangian of eqs. (2.4) and (2.10) we can now analyze the mechanism of leptogenesis that results. First, since the A terms are of order $m_{3/2}$, that is to say of order unity at the $M_{\tilde{N}_i}$ scale, they will put the right-handed sneutrinos in deep thermal equilibrium. Therefore the decay of the sneutrinos cannot lead to the creation of a large asymmetry. The right-handed neutrinos on the other hand have tiny effective Yukawa couplings of order $10^{-7} - 10^{-8}$ and therefore will be naturally out of equilibrium, independent of the situation for the sneutrinos. In addition, the small Yukawa couplings of the right-handed neutrinos implies that the usual leptogenesis vertex diagram leads to a far too small asymmetry (i.e. $\varepsilon_{N_i} \sim \lambda^2/8\pi < 10^{-15}$), and therefore it can be neglected. On the other hand, the self-energy diagram of Fig.1 for the rhd neutrinos, although also suppressed by Yukawa couplings, can be enhanced by a resonance effect if the mass splittings are naturally tiny. The N_i asymmetry in this case is [11, 12, 8]:

$$\varepsilon_i = -\sum_{j\neq i} \frac{M_i}{M_j}\frac{\Gamma_j}{M_j} I_{ij} S_{ij} \,, \tag{3.1}$$

where

$$I_{ij} = \frac{\mathrm{Im}\,[(\lambda\lambda^\dagger)^2_{ij}]}{|\lambda\lambda^\dagger|_{ii}|\lambda\lambda^\dagger|_{jj}}\,, \qquad S_{ij} = \frac{M_j^2\,\Delta M_{ij}^2}{(\Delta M_{ij}^2)^2 + M_i^2\Gamma_j^2}\,, \qquad \Gamma_j = \frac{|\lambda\lambda^\dagger|_{jj}}{8\pi} M_j\,. \tag{3.2}$$

In the model of eqs. (2.4) and (2.10), the lowest order $1/M_P$ contribution to the right-handed neutrino masses is flavor blind, eq. (2.5), so the right-handed neutrino splittings vanish, which in turn leads to a vanishing asymmetry. However, as we explained above, at the next order in $1/M_{\rm Planck}$, there are two sources of mass degeneracy breaking, eq. (2.11), one from the term $X_{ik}X^\dagger_{kj}Y^\dagger N_iN_j/M_P^3$, and a second contribution of the same, small size, but with in general different flavour structure, from the Z_{ij}-dependent term of eq. (2.10).

It is important that the tiny mass splittings thus induced among M_{N_i} mass eigenstates are of the same parametric size as the Yukawa-coupling-induced decay width of these massive states. Thus the propagator of the virtual rhd neutrinos in the self energy diagrams, eq. (3.2), will be naturally at the resonance or close to it. In more detail both contributions lead generically to mass splittings and decay widths of order

$$\Delta M_{ij}^2 \sim m_{3/2}^2 \frac{m_{3/2}}{M_P} \qquad \text{and} \qquad \Gamma_i \sim \frac{m_{3/2}}{8\pi}\frac{m_{3/2}}{M_P} \tag{3.3}$$

which therefore leads to

$$S_{ij} \sim \frac{M_P}{m_{3/2}} \sim \frac{M_i}{8\pi\Gamma_i} \tag{3.4}$$

and therefore to a possible total asymmetry of size $\varepsilon_i \sim I_{ij}/8\pi$ which can be as large as $1/8\pi$ assuming there are non non-trivial cancellations within I_{ij} (see [8]). To address this last issue it is necessary to diagonalise the M_N^R mass matrix and to calculate the corresponding asymmetry in the mass eigenstates basis. To this end let us consider the two cases, $\gamma = 0$, and $\gamma \neq 0$.

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4. Single source of flavour breaking: The $\gamma = 0$ case

This pattern gives mass splittings naturally at the resonance as explained above but nevertheless turns out to result in $I_{ij} = 0$ leading to a vanishing asymmetry. This can be seen in the following way: Since the mass matrix of eq. (2.11) is real and symmetric it can be diagonalised by a real orthogonal matrix O giving $O[a_X a_X^\dagger + (a_X a_X^\dagger)^T]O^T = D$ where D is a real diagonal matrix. Furthermore writing $O(a_X a_X^\dagger)O^T \equiv D/2 + C$ it is simple to check that C is a purely imaginary antisymmetric matrix. The one-loop mass eigenstate Yukawa couplings, $\lambda^{(1)}$, are therefore related to the tree level couplings, λ, via a flavour rotation $\lambda^{(1)} = O\lambda$. Upon substitution of the mass-eigenstate basis Yukawa's into the I_{ij} of eq. (3.1) we obtain $I_{ij} \sim \mathrm{Im}[(\lambda^{(1)}\lambda^{(1)\dagger})_{ij}^2] \propto \mathrm{Im}[(Oa_X a_X^\dagger O^T)_{ij}^2] \propto \mathrm{Im}[(D/2 + C)_{ij}^2] = 0$, for $i \neq j$. Thus we learn that a single source of flavour breaking as contained in the Yukawa couplings can lead to non-trivial mass splittings among the N_i, but does *not* lead to any CP-violation. Intuitively this is not too surprising since the self energy diagram of figure 1 is the same as the one in the leptogenesis self-energy diagram, figure 2.

Note that, in the context of the ordinary see-saw extended SM, the possibility of one-loop induced resonant leptogenesis has been considered in ref. [25], for a phenomenological case with two right-handed neutrinos with $M_{N_1} = M_{N_2}$ but different Yukawa couplings, and it was claimed that this leads to a naturally large resonantly enhanced asymmetry. However in this calculation it appears that the numerator of the asymmetry has been calculated with the tree level Yukawa couplings. Since the full, one-loop corrected masses differ from the tree level ones only by a small amount, one might expect that it is a good approximation not to include the full one-loop correction that arises from going to the mass eigenstate basis. However since the tree level masses are exactly degenerate, even a small off-diagonal contribution leads to large mixing and therefore it is necessary to renormalise as we have done here. We then see that in fact this system gives a vanishing asymmetry.

We conclude from the above study of the $\gamma = 0$ case that degenerate rhd neutrinos with different Yukawa couplings do not lead to a one-loop self-energy leptogenesis contribution. For the resonant leptogenesis mechanism to work it is essential to have at least two sources of flavor breaking, so we now turn to the analysis with $\gamma \neq 0$.

5. Two sources of flavour breaking: the $\gamma \neq 0$ case

When $\gamma \neq 0$ in eq. (2.11) a second source of flavour structure is introduced, and as we will argue in this Section, it can give I_{ij} factors which can be as large as 1 and S_{ij} factors naturally close to the resonance, that is to say of order $M_{N_i}/2\Gamma_i$, leading to CP asymmetries of order one.

To convince oneself that the CP-asymmetry can now be non-zero, it is useful to choose a special form for the three by three $\langle Z \rangle = a_{Z_{ij}} m_I$ matrix in eq. (2.11) such that the off-diagonal terms in the total mass matrix of eq. (2.11) cancel between the β- and γ-dependent terms, but leaving non-trivial diagonal entries $M_N^R = M_N \mathbb{I} + \mathrm{diag}(\delta M_{N_1}, \delta M_{N_2}, \delta M_{N_3})$ where $\delta M_{N_1} < \delta M_{N_2} < \delta M_{N_3}$. The δM_{N_i} are naturally of size $m_{3/2}(m_{3/2}/M_{\mathrm{P}})$, and can be chosen to be positive, leading to a mass splitting of the same order.[6] Since the total mass matrix is a diagonal matrix the Yukawa couplings in this case do not need to be rotated to the mass eigenstate basis and I_{ij} is therefore just a function of the lowest order Yukawa coupling matrix λ_{ij}. In addition, one can re-express the total asymmetry as

$$\varepsilon_{\rm tot} = -A_{21}I_{21} - A_{31}I_{31} - A_{32}I_{32}\,, \tag{5.1}$$

where $\varepsilon_{\rm tot} = \sum_i \varepsilon_i$, the I_{ij}'s are defined in eq. (3.2), and with all A_{ij}'s in eq. (5.1) being the positive quantities:

$$A_{ij} = \frac{M_i M_j}{8\pi}\Delta M^2_{ij}\left[\frac{|\lambda\lambda^\dagger|_{jj}}{(\Delta M^2_{ij})^2 + M_i^2\Gamma_j^2} + \frac{|\lambda\lambda^\dagger|_{ii}}{(\Delta M^2_{ij})^2 + M_j^2\Gamma_i^2}\right]. \tag{5.2}$$

It is now easy to see that one can choose a form for the Yukawa couplings such that, eg, $I_{21} = I_{31} = 0$ while $I_{32} \neq 0$ and $A_{32} \neq 0$. Thus $\varepsilon_{\rm tot}$ is non zero in this case. This completes the existence proof that $\gamma \neq 0$ can lead to non-zero asymmetries as claimed.

Moreover, general forms for the Yukawas and $\langle Z_{ij}\rangle_A$'s will produce independent $A_{ij}I_{ij}$ terms and consequently give rise to a non zero asymmetry, and applying the relations displayed in eq. (3.3) we get,

$$\varepsilon_{\rm total} \sim -\frac{1}{8\pi}(aI_{21} + bI_{31} + cI_{32})\,. \tag{5.3}$$

where a, b, c are $\mathcal{O}(1)$ coefficients arising from the self-energy terms in eq. (5.2). A straightforward numerical investigation of the dependence of eq. (5.3) on the forms and phases of the effective Yukawa matrix λ_{ij}, together with the α and β terms of the mass matrix eq. (2.11), shows that it is simple to find cases where $aI_{21} + bI_{31} + cI_{32} \sim 1$.[7]

Furthermore, for natural values of the model parameters, the produced lepton asymmetry is not suppressed by any wash-out effect. The Yukawa couplings which are of order $10^{-7} - 10^{-8}$ give a decay width smaller than the Hubble constant and therefore will not induce any wash-out effect via decay or via scatterings, neither in the Boltzmann equation of the N_i number density nor in the Boltzmann equation of the lepton number density. Moreover the A-terms, even if large, $A \sim$ TeV, cannot change this result because they can effect the N_i number densities only via scatterings which are also suppressed by Yukawa couplings, and they do not affect the lepton number density because they can break lepton number only when accompanied by Yukawa couplings or by a B term which is also very suppressed. In particular, it can be checked that the potentially dangerous $\tilde{L} + H \leftrightarrow \tilde{N} \leftrightarrow \tilde{L}^* + H^*$ process induced by two A terms and a B term has in fact a rate smaller than the Hubble constant. Thus the produced n_L/s can be naturally of order $\varepsilon_{N_i}/g_\star \sim \varepsilon_{N_i}/100$ and

[6] This example requires a tuning between both contributions of eq. (2.11) but it is convenient to demonstrate that a large asymmetry can be obtained.

[7] Note, that unlike the usual see-saw model based leptogenesis case, one cannot re-write the I_{ij}'s in terms of the solar and atmospheric neutrino mass-squared differences, mixing angles, and phases in the conventional way. The reason for this is that the light neutrino spectrum is dominantly set by the one-loop contribution of eq. (2.6) which depends upon the B-term susy- and lepton-number-breaking mass of the sneutrinos, with only small corrections (possibly leading to the small $\Delta m^2_{\rm solar}/\Delta m^2_{\rm atm}$ light neutrino hierarchy) arising from a tree-level see-saw contribution depending upon λ in the usual way, see ref. [19].

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therefore from the above discussion can be naturally as large as $1/(100 \cdot 8\pi)$ for order one couplings! To our knowledge, this is the only model of thermal leptogenesis which can lead naturally to such a large asymmetry. Note however that the asymmetry can be rapidly suppressed if we allow the phase to be non-maximal and the couplings to be not all of order unity. Any deviation of the constants h, a_X, or a_Z by one order of magnitude in eq. (5.2) can lead to several orders of magnitude suppression of the asymmetry. This can lead easily to the CMBR-determined experimental value: $n_B/n_\gamma = 6.1^{+0.3}_{-0.2} \cdot 10^{-10}$ [26].

What about the testability of our model? The main attractive point here is that since we have large A terms the sneutrinos could be observed quite easily and allow one to test to a large extent the one-loop diagrams at the origin of the neutrino mass. From the additional discovery of supersymmetry we could in further conclude that TeV-scale right-handed neutrinos also must exist. Although the later could not be observed because they can be produced only by the tiny Yukawa couplings, the primary ingredients of our TeV-scale leptogenesis mechanism could thus be tested.

6. Conclusions

In summary, in the framework of softly broken supersymmetric theories we have proposed a natural model of resonant leptogenesis utilizing TeV scale rhd neutrinos, N_i. Our discussion is based upon recent models of the generation of light neutrino masses from supersymmetry-breaking effects, which provide a natural explanation for the presence of rhd neutrinos at the TeV scale, and for the small Yukawa couplings and B-terms necessary for the correct resulting light neutrino spectrum. We show that these properties also lead to large cosmological lepton asymmetries, via the resonant behavior of the one-loop self-energy contribution to N_i decay. The model addresses the primary problems of previous phenomenological studies of low-energy leptogenesis: a rational for small Yukawa couplings so that the out-of equilibrium condition on the N_i decay is satisfied; the origin of tiny, but non-zero mass splitting required between at least two N_i masses; and the necessary non-trivial breaking of flavour symmetries in the rhd neutrino sector. The low mass-scale of the rhd neutrinos and their superpartners, and the weak-scale A-terms automatically contained within the model, offer opportunities for partial direct experimental tests of this leptogenesis mechanism at future colliders.

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