Low-scale leptogenesis and soft supersymmetry breaking

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We investigate the possibility of low-scale leptogenesis in the minimal supersymmetric standard model extended with right-handed (rhd) neutrinos. We demonstrate that successful leptogenesis can be easily achieved at a scale as low as \( \sim \) TeV where lepton number and CP violation comes from soft supersymmetry breaking terms. The scenario is shown to be compatible with neutrino masses data.

A. Introduction. The experimental observations of neutrinos oscillations gave overwhelming evidence for small neutrino masses. The see-saw mechanism [1] can explain elegantly such small masses from the existence of right-handed (rhd) neutrinos. Furthermore, in the leptogenesis scenario [2], the out-of-equilibrium decay of these rhd neutrinos can lead to a lepton asymmetry, that is partly converted to a baryon number through sphalerons, providing in this way a simple and attractive explanation of the baryon asymmetry of the universe.

In the standard thermal leptogenesis scenario the mass of rhd neutrinos must lie above \( 10^9 \) GeV or so [3–8]. (here, we are not considering the case where rhd neutrinos are quasi-degenerate [5,8–13]). In supergravity, this implies the well known gravitino problem [14]. To avoid this gravitino problem and also, independently of it, in order to be as close to experiment as possible, it would be nice to have leptogenesis at the lowest possible scale, i.e. near the Fermi scale. Low rhd neutrino masses can occur naturally in realistic supersymmetric theories such as the minimal Pati-Salam model [15] or if the rhd neutrino masses themselves come from supersymmetry breaking [16].

Building such a low energy leptogenesis model is however difficult for a number of reasons (see [5] for a detailed discussion). The main reason is that if all \( \bar{L} \)-violating interactions come from the see-saw then the asymmetry is proportional to Yukawa couplings which explains the small neutrino masses have to be tiny, leading to a far too small asymmetry. We need therefore other sources of \( L \)-violation which do not give rise to see-saw neutrino masses. The most natural and simple framework leading to such interactions is low-energy supersymmetry. By transferring the notion of lepton number to scalar partners, supersymmetry introduces new sources of lepton number violation through soft supersymmetry breaking [17]. Being pure scalar, these interactions are less constrained by the neutrino masses (since they lead to neutrino masses only at one loop, as we will see) and therefore allow to get much larger asymmetries at low scale, leading to successful leptogenesis. This is the central point of this letter.

B. Soft SUSY breaking terms. Let us consider the \( R \)-conserving MSSM extended by a singlet rhd neutrino for each generation \( N_i \). The model is described by the usual SUSY see-saw superpotential

\[
W = W_{\text{MSSM}} + Y_{ij} L_i H_U N_j + \frac{1}{2} M_{ij} N^2_i, \tag{1}
\]

where we have rotated the \( N_i \)'s into the basis where the rhd neutrino mass matrix is real and diagonal. We are interested in the situation where the mass of rhd (s)neutrinos is above but not too far from the scale of the supersymmetry breaking. Following a bottom-up approach, we consider the most general soft SUSY breaking terms compatible with gauge invariance and \( R \)-parity conservation. The relevant \( L \) and \( CP \) violating terms in the Lagrangian are given by

\[
\mathcal{L}^- = (m^2_N)_{ij} \tilde{N}_i^* \tilde{N}_j + B_{ij} \tilde{N}_i \tilde{N}_j + A^U_{ij} \tilde{L}_i H_U \tilde{N}_j + A^D_{ij} \tilde{D}_i H_D \tilde{N}_j + A^D_{ij} \tilde{D}_i H_D \tilde{N}_j + \text{h.c.} \tag{2}
\]

The first line of Eq. (2) represents the usual soft masses, \( B \)-term and holomorphic \( A \)-terms, generally present in gravity mediated scenarios. The additional terms are the so-called non-holomorphic \( A \)-terms, and they are highly suppressed in supergravity. Although they are not essential for our discussion, we include them for the sake of completeness.

Note the important role \( R \)-parity is playing here. In general, \( R \)-parity is invoked in order to prevent a too fast proton decay. It also provides a natural dark matter candidate (LSP). In our case, \( R \)-parity makes Eq. (2) the most general renormalizable, \( B-L \) violating superpotential with this field content. Furthermore, and due to the presence of a singlet in the model, \( R \)-parity prevents the occurrence of dangerous tadpoles that induce quadratic divergences. Indeed, if we relax \( R \)-parity, we would have \( \lambda_{ijk} \tilde{N}_i \tilde{N}_j \tilde{N}_k \) as a soft term, that would induce a tadpole for the operator \( \tilde{L}_i H_U \).

It is remarkable that the \( B-L \) symmetry leads automatically to \( R \)-parity conservation [18]. After the subsequent spontaneous breaking of \( B-L \), which leads to non-vanishing rhd neutrino masses, exact \( R \)-parity survives as a discrete \( Z_2 \) symmetry. This is true at all energy scales [19]. In other words, \( R \)-parity is inherent in this picture of the see-saw mechanism and leptogenesis through the
spontaneous breaking of $B-L$ symmetry. It is quite natural to expect that it survives the spontaneous breaking of supersymmetry.

Regarding the lower limit on the mass of rhd neutrinos in the standard leptogenesis scenario, one could imagine that the situation could change dramatically due to a natural presence of the $SU(2)_L$ triplet superfields (necessarily present in the LR symmetric theories). It turns out that the situation is very similar to the standard one [20], and thus the soft supersymmetry breaking terms are really indispensable.

Going to the sneutrino mass basis $\tilde{N}_I$ ($I = 1, \ldots, 6$) resulting from the diagonalization of the 6 by 6 mass matrix containing the three types of mass term in Eqs. (1) and (2), and rephasing the $\tilde{N}_I$ so that they are real fields, the Lagrangian reduces to the compact form

$$L_N = \tilde{M}_N^2 \tilde{N}_I + \nu_{ij} \tilde{N}_I \phi_\alpha + \nu_{ij}^* \tilde{N}_I \phi_\alpha^*, \quad (3)$$

where $\phi_{1,2} \equiv H_U, H_D^*$. The $\nu_{ij}$ are related to the initial soft parameter $A^U_{ij}, A^D_{ij}, A^D_{ij}$ and $A^D_{ij}$ through the rhd sneutrino mixing matrix.

**C. Leptogenesis.** Since, due to neutrino mass constraints, low scale rhd neutrinos must generally have tiny Yukawa couplings the rhd neutrino asymmetries will be highly suppressed. One possibility to compensate this suppression is to consider a highly degenerate spectrum of rhd neutrinos. In this case the asymmetry can be highly resonantly enhanced. We will not consider this possibility here and assume that the various rhd neutrinos and sneutrinos have a hierarchical mass spectrum (for a low energy model based on degeneracy also in the framework of supersymmetry breaking theories see [13]).

An other possibility we could think of is to invoke a hierarchy between the couplings of the virtual and real particles entering in the leptogenesis diagrams. This consists in taking small couplings for the particle decaying in order to satisfy the out-of-equilibrium condition $\Gamma_D < H$ and to take larger couplings for the (heavier) virtual particle, since those couplings are not constrained by this condition. This at a scale as low as few TeV doesn’t work for the rhd neutrinos due to the neutrino constraints (for more details see [5] and [3,4]). However for the sneutrinos this simple possibility could work because they are not inducing neutrino masses directly through the see-saw mechanism but only at the one loop level.

The diagrams for the decay of the sneutrinos which can lead to successful leptogenesis in this way are given in Fig 1. They involve only scalar fields. From Eq. (3) these diagrams give the following $CP$ asymmetry:

$$\varepsilon_I = \frac{\Gamma(\tilde{N}_I \rightarrow \tilde{L}_j \phi_\alpha) - \Gamma(\tilde{N}_I \rightarrow \tilde{L}_j^* \phi_\alpha^*)}{\Gamma(N_I \rightarrow L_j \phi_\alpha) + \Gamma(N_I \rightarrow L_j^* \phi_\alpha^*)} = \varepsilon_I^V + \varepsilon_I^S, \quad (4)$$

where as $\varepsilon_I^V$ and $\varepsilon_I^S$ (self-energy and vertex diagrams respectively) are given by

$$\varepsilon_I^V = \frac{-1}{8\pi M_N^2} \frac{1}{|\nu_{ij}|^2} \sum_{K \neq I} \text{Im} \left[ \nu_{1m} \nu_{Km} \nu_{1j} \nu_{Kj} \right] F_V(x_K), \quad (5)$$

$$\varepsilon_I^S = \frac{-1}{4\pi M_N^2} \frac{1}{|\nu_{ij}|^2} \sum_{K \neq I} \text{Im} \left[ \nu_{1m} \nu_{Km} \nu_{1j} \nu_{Kj} \right] F_S(x_K), \quad (6)$$

with $x_K = M^2_{N_I}/M^2_{N_\alpha}$. The loop functions $F_V,S(x)$ can be calculated easily. They are given by

$$F_V(x) = \ln(1 + x), \quad F_S(x) = x/(1 - x). \quad (7)$$

As we did already in the denominator of the $CP$ asymmetry, in the following we will assume for simplicity that the Yukawa couplings are negligible. These couplings are not essential in our scenario, neither for leptogenesis, nor for the neutrino masses. The suppression effects they can induce if they are not negligible will be studied in a further publication. Now, since we have assumed a hierarchical spectrum of sneutrino masses in Eq. (1), it is a very good approximation to neglect the asymmetry produced from the decay of the 4 heaviest eigenstates. Furthermore for simplicity, in order to show that sufficient leptogenesis can be created easily, without loss of generality, one can consider only the asymmetry produced by the lightest eigenstate $N_1$ which at lowest order in $M_{N_1}^2/M_{N_\alpha}^2$ is given by:

$$\varepsilon_1 \simeq -\frac{3}{8\pi} \frac{1}{M_{N_\alpha}^2} \frac{1}{|\nu_{ij}|^2} \sum_{K \neq 1} \text{Im} \left[ (\nu_{1m} \nu_{1j})^2 \right]_{1K}, \quad (8)$$

where we have neglected the terms where $\alpha \neq \beta$ again for the sake of simplicity. To have successful leptogenesis from Eq. (8) there are essentially three constraints:

- The out of equilibrium condition for $\tilde{N}_1$ eigenstate

$$\Gamma_{\tilde{N}_1} = \frac{1}{4\pi} \frac{|\nu_{ij}|^2}{M_{N_1}} < H(T = M_{N_1}) \quad (9)$$

translates into a bound on its couplings $|\nu_{ij}|^2 \lesssim 10^{-7} M_{N_1}$, if $M_{N_1} \sim 1$ TeV or $|\nu_{ij}|^2 \lesssim 4 \cdot 10^{-7} M_{N_1}$ if $M_{N_1} \sim 10$ TeV. In the following, we will assume that this condition is satisfied in order to avoid wash-out suppressions coming from these couplings.

- The couplings of the virtual sneutrino eigenstate must be large enough to give sufficient $CP$-asymmetry. In order to reproduce the experimental value from CMB,
that the potentially dangerous scattering is of order $10^{-6}$, which is the strongest assumption we have to make here in order that this mechanism work. This might seem a large hierarchy, but after all it is of order the ratio of tau to electron Yukawa couplings.

- In order to avoid that the soft interactions of the virtual $N_2$ could wash out the asymmetry it is necessary that the potentially dangerous scattering $L + H \leftrightarrow N_2 \leftrightarrow L^* + H^*$ be under control. This scattering is not present in the Boltzmann equation for the $N_1$ number density but rather in the one for the lepton number. For $T \sim M_{N_2}$ the on-shell contribution to this scattering is quite fast because with $\mu_{2j} \sim 10^{-3} \langle M_{N_2} \rangle$ and $M_{N_2} \sim$ few TeV we have $\Gamma_{N_2} \gg H(T = M_{N_2})$. For $M_{N_2} \sim$ few $M_{N_1}$ this contribution, even if Boltzmann suppressed, remains fast down to a temperature of order $M_{N_1}$ or few times less. If $\Gamma_{N_1} \sim H(T = M_{N_1})$ this can induce a sizable wash-out suppression of the asymmetry. However for $\Gamma_{N_1}$ a few times smaller than $H(M_{N_1})$ the asymmetry will be produced at smaller temperature when this suppression is further Boltzmann suppressed and negligible. Similarly the off-shell contribution to this scattering can have an effect, especially for low temperatures. However this effect will be fastly Boltzmann suppressed and negligible at temperatures below the threshold $s_0 = (m_H + m_{N_1})^2$. It is therefore easy to avoid large wash-out from this scattering as we have checked by considering explicitly the corresponding Boltzmann equations.

Note that if our scenario is to be embedded into a theory where $B - L$ is gauged, such as in say Pati-Salam theory or $SO(10)$, wash-out constraints require that the corresponding gauge boson mass to be much heavier than $M_{N_1}$. This happens naturally if the Yukawa couplings giving mass to rh neutrinos are small.

Altogether for example with $M_{N_1} \sim 2$ TeV, $M_{N_2} \sim 6$ TeV, $(\mu_{ij})^\text{max} \sim 5 \cdot 10^{-8} M_{N_1}$, $(\mu_{2j})^\text{max} \sim 10^{-3} M_{N_2}$, $m_\mu + m_L = 700$ GeV, one can check that a large enough asymmetry can be created. It gives $\varepsilon_1 \sim 10^{-7}$, and from the Boltzmann equations we get $n_B / n_\gamma \sim 6 \cdot 10^{-10}$ (in agreement with data [21]). The neutrino mass constraints can be easily accommodated with this set of values (see below). There is a large range of parameters in the parameter space which leads to successful leptogenesis. Note however that it appears to be difficult to generate a large enough asymmetry before the sphalerons gets out-of-equilibrium around $T \sim 100 - 200$ GeV for $M_{N_1}$ below one TeV and $M_{N_2}$ below 3-4 TeV. Finite temperature effects can change the produced asymmetry by effects of order unity [7] which we didn’t take into account here. Note that all constraints can be relaxed by scaling up all masses.

The impact of our results on the original basis in Eq. (2) is worth studying in well defined theories of supersymmetry breaking, and we plan to return to this issue elsewhere. This requires typically similar hierarchies between some of the couplings of different generations of rh neutrinos. In addition, their mixings have not to be larger that $\sim 10^{-4}$ in order that the decay rate of the lightest sneutrino remains sufficiently suppressed. This implies in particular an alignment of the $B$ terms, i.e. they should be almost diagonal.

### D. Neutrino Masses

In our scenario, neutrinos masses originate from two sources.

1. **See-saw contribution.** The first one, which occurs at tree level, is the usual see-saw given by:

   $$m_{\nu}^{\text{tree}} \sim -Y^T M_{\text{SM}}^{-1} Y (H_U)^2.$$  \hspace{1cm} (10)

In order that this doesn’t induce too large neutrino masses with rh neutrino masses of order TeV the Yukawa couplings have to be tiny, i.e. $Y \lesssim 10^{-7} - 10^{-6}$. As said above, here for simplicity we assume all effects of Yukawa couplings negligible.

2. **Radiative soft contribution.** The second source of neutrino masses is the radiative one-loop contribution of Fig. 2 coming from the sneutrinos soft term sector (see also Ref. [16]). The resulting radiative neutrino mass in the limit $m_{N_i} \gg m_{\tilde{\nu}}$, $m_\chi$ is

   $$(m_{\nu}^{\text{rad}})_{jk} \approx \frac{\alpha}{4\pi} \frac{\mu_{ij} \mu_{jk}}{M_{N_i}^2} \frac{m_X}{m_{\tilde{\nu}_j}^2 - m_{\tilde{\nu}_j}^2} \langle \phi_a \rangle \langle \phi_b \rangle \times \left\{ \frac{m_{\tilde{\nu}_j}^2}{m_{\tilde{\nu}_j}^2 - m_\chi^2} \ln \frac{m_{\tilde{\nu}_j}^2}{m_\chi^2} - j \to k \right\}.$$ \hspace{1cm} (11)

The estimate of the above contribution depends crucially on the masses of rh sneutrinos $m_{\tilde{\nu}_j}$, and for $m_{\tilde{\nu}_j}$ large enough it is negligible. However we have seen in the previous section that $m_{\tilde{\nu}_j}$ can be as low as TeV from the leptogenesis discussion, and for such value the induced masses turns out to be not negligible. Plugging in the values of $\mu_{ij}$ and $m_{\tilde{\nu}_j}$ (i.e. $\mu_{ij} \sim 10^{-3} M_{N_2}$ and $M_{N_2} \sim$ 6 TeV and a typical value for $m_{\nu_i} \approx 500$ GeV

![Fig. 2: Diagram contributing to neutrino masses.](image-url)
and a somewhat smaller $m_{\chi} \approx 100$ GeV as in the numerical example above), it is easy to show that $m_{\nu}^{\text{rad}} \approx 1$ eV! A nice feature of our scenario is therefore the fact that the sets of parameters leading to successful leptogenesis also lead to a maximal neutrino mass in agreement with data within one or two orders of magnitude (although close to the upper experimental limit). The correct neutrino flavor structure can be then obtained by considering flavor hierarchies between the $\mu_{IJ}$ couplings.

Taken at face value, this would imply degenerate neutrino masses (especially for large gaugino masses), which interestingly enough is in the sensitivity region of present experiments [22]. However, there can be accidental (unnatural) cancellations with the see-saw contribution. Moreover, one could use easily the freedom of choosing appropriately tan$\beta$ (and taking $\mu_{1j}^{\text{rad}}$ couplings slightly smaller than $\mu_{1j}^3$ couplings) and thus easily getting the appropriate suppression in the case of hierarchical neutrino masses. Barring these mild fine tunings for a low $B - L$ scale, we generically expect degenerate neutrinos. Although not a hard prediction, it would indicate a possible low scale leptogenesis scenario discussed here.

E. Lepton flavor violation. As in any low-energy supersymmetric framework with soft supersymmetric breaking terms, in our scenario we expect an appreciable amount of lepton flavor violation. In fact, rare processes like $\mu \rightarrow e\gamma$ are used to set stringent limits on slepton masses and mixings. In our case, due to the smallness of $\mu_{1j}^{\text{rad}}/m_{\tilde{g}}^2$, and to the alignment of the $B$-terms, it is easy to see that the contributions to these processes coming from Eq. (3) are well below experimental limits.

F. Collider signatures. The real test of this scenario would be of course the production of rhd (s)neutrinos in the 1 - 10 TeV region. Due to their tiny couplings, this will be a very hard task for the rhd neutrinos. But it is worth noting that one expects to produce rhd sneutrino much more efficiently than their fermionic partners, since the coefficients of the soft terms (i.e. $\mu_{1j}^3$ in the example given above) are necessarily much bigger than the Dirac Yukawa couplings. This provides a unique opportunity to test directly the origin of both neutrino masses and leptogenesis at the same time.

In conclusion, in addition of providing a simple solution to the hierarchy problem and protecting the flatness of the inflationary potential, supersymmetry could also lead to the generation of the baryon asymmetry at a scale of order the Fermi scale.

We thank B. Bajc, J. March-Russell and S. West for useful conversations and comments. The work of L.B. is supported by PPARC, of T.H. by EU Marie Curie contract HPMF-CT-01765 and of G.S. partially supported by EEC, under the TMR contracts ERBFMRX-CT960090 and HPRN-CT-2000-00152. L.B. thanks the high energy section of the Abdus Salam ICTP for kind hospitality during the last stage of this work.

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