R-parity violation assisted thermal leptogenesis in the seesaw mechanism

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Successful leptogenesis within the simplest type I supersymmetric seesaw mechanism requires the lightest of the three right-handed neutrino supermultiplets to be heavier than $\sim 10^6$ GeV. Thermal production of such (s)neutrinos requires very high reheating temperatures which result in an overproduction of gravitinos with catastrophic consequences for the evolution of the universe. In this letter, we let R-parity be violated through a $\lambda_i N_i H_u H_d$ term in the superpotential, where $N_i$ are right-handed neutrino supermultiplets. We show that in the presence of this term, the produced lepton-antilepton asymmetry can be enhanced. As a result, even for $N_i$ masses as low as $10^6$ GeV or less, we can obtain the observed baryon asymmetry of the universe without gravitino overproduction.

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The recent neutrino data [1, 2, 3] shows that neutrinos are massive [4]. One of the most popular ways to generate tiny nonzero neutrino masses is the seesaw mechanism [5, 6], which adds three right-handed neutrinos to the standard model with very heavy masses, $M_3 > M_2 > M_1 > m_{weak}$.

The seesaw mechanism also provides us with a framework to obtain the observed baryon-antibaryon asymmetry of the universe through a process called leptogenesis [7]. However, in the context of supersymmetry, this process suffers from a phenomenon called gravitino overproduction [8]. If we assume that lightest right-handed (s)neutrinos are thermally produced in the early universe, the reheating temperature ($T_R$) should be higher than $\sim 10^6$ GeV [2, 10]. The high reheating temperature can lead to the overproduction of gravitinos, which has catastrophic consequences for the evolution of the universe. The upper bound on $T_R$ from gravitino overproduction considerations depends on the details of model. If the gravitino has hadronic decay modes, we expect $T_R < 10^6 - 7$ GeV [11]. In the literature, a variety of solutions for this problem has been suggested [12, 13, 14, 15, 16, 17].

In this letter, we suggest an alternative solution based on the R-parity violation. The produced lepton-antilepton asymmetry can be in the expected range, even for masses of the lightest of the right-handed (s)neutrinos lower than $10^6$ GeV, avoiding gravitino overproduction.

First recall that in the simplest type I supersymmetric seesaw mechanism the superpotential is given by

$$W = \sum_{i,j} \epsilon_{\alpha\beta} (Y_{\nu})_{ij} \tilde{N}_i \tilde{L}_j \tilde{H}_u^\beta + \frac{1}{2} \sum_{ij} M_{ij} \tilde{N}_i \tilde{N}_j$$

where $\tilde{L}_j$ is the superfield associated with the left-handed lepton doublet $(\tilde{l}_j, \tilde{\nu}_j)$ and $\tilde{H}_u$ is the Higgs doublet that gives mass to the up quark. The first term is the familiar Yukawa coupling and the second is the Majorana mass term of right-handed neutrinos.

Relaxing R-parity conservation, we can add the following term to the superpotential

$$W_{RPV} = \sum_i \epsilon_{\alpha\beta} \lambda_i \tilde{N}_i \tilde{H}_d^\alpha \tilde{H}_u^\beta$$

where $\tilde{H}_d$ is the Higgs doublet that gives mass to the down quark. The existence of this R–Parity violating (RPV) term has recently been advocated to solve the $\mu$ problem [19]. In our case its contribution to generating the $\mu$ term is negligible because $N_i$, being super-heavy, do not acquire sizeable vacuum expectation values. However, this term will play a key role in making thermal seesaw leptogenesis viable.

Note that R–Parity violation in supersymmetry has been advocated as an attractive origin for neutrino masses, alternative to the supersymmetric seesaw [19]. Neutrino masses are typically hierarchical, with the atmospheric scale arising at tree level and the solar one calculable as radiative corrections [20]. However here we propose that neutrinos acquire masses $a la$ seesaw and that, although RPV is necessary to produce the observed baryon asymmetry of the universe, it is not the dominant source of neutrino masses.

Without loss of generality, we can rotate and rephase the fields to make the mass matrix $M_{ij}$ real diagonal. In this basis, the elements of $Y_{\nu}$ and $\lambda$ can in general be complex. Introduction of the coupling $\lambda_i$ adds three extra CP-violating phases to the theory, as we shall see, with consequences for the baryon asymmetry of the universe.
Let us define the following asymmetries

\[ \epsilon_{N_1} = - \sum_i \left[ \frac{\Gamma(N_1 \rightarrow l_i \bar{H}_u) - \Gamma(N_1 \rightarrow l_i H_u)}{\Gamma_{\text{tot}}(N_1)/2} + \frac{\Gamma(N_1 \rightarrow \bar{l}_i H_u) - \Gamma(N_1 \rightarrow \bar{l}_i \bar{H}_u)}{\Gamma_{\text{tot}}(N_1)/2} \right] \] (3)

and

\[ \epsilon_{\tilde{N}_1} = - \sum_i \left[ \frac{\Gamma(\tilde{N}_1 \rightarrow l_i \bar{H}_u) - \Gamma(\tilde{N}_1 \rightarrow l_i H_u)}{\Gamma_{\text{tot}}(\tilde{N}_1)/2} + \frac{\Gamma(\tilde{N}_1 \rightarrow \bar{l}_i H_u) - \Gamma(\tilde{N}_1 \rightarrow \bar{l}_i \bar{H}_u)}{\Gamma_{\text{tot}}(\tilde{N}_1)/2} \right] \] (4)

where \( N_1 \) and \( \tilde{N}_1 \) are respectively the lightest right-handed neutrino and sneutrino and \( \Gamma_{\text{tot}}(N_1) \) and \( \Gamma_{\text{tot}}(\tilde{N}_1) \) are their total decay rates. We expect the produced lepton-antilepton asymmetry to be proportional to

\[ \epsilon \equiv \epsilon_{N_1} + \epsilon_{\tilde{N}_1}. \]

In the following, we show that the R-parity violating term that we have introduced gives a new contribution to \( \epsilon_{N_1} \) and \( \epsilon_{\tilde{N}_1} \) which for certain range of parameters can enhance the effect. We show that as a result of this enhancement, even for \( M_1 \) as low as \( 10^6 \) GeV, we can have successful leptogenesis and simultaneously generate tiny masses for neutrinos [i.e., in the simplest type-I seesaw down to \( \sim 10^{-5} \) GeV], where \( v \) is 245 GeV and \( \beta = \arctan([\langle H_u \rangle/\langle H_d \rangle]). \)

Therefore, thermal production of \( N_1 \) and \( \tilde{N}_1 \) does not need too high reheating temperature and the universe would not encounter gravitino overproduction.

Fig. 1 shows the structure of the diagrams contributing to \( \epsilon_{N_1} \) and \( \epsilon_{\tilde{N}_1} \). Each line collectively represents the bosonic, fermionic or auxiliary component of the indicated superfield. The vertices marked with dots are the new Yukawa, fermionic or auxiliary coupling while those marked with \( \otimes \) are the new R-parity violating vertices given by \( \lambda_i \). Each line can be either bosonic or fermionic when appropriate. Reversing the arrows we reach the diagrams that produce antileptons instead of leptons.

Notice that both in the vertex-type diagram (b) and wave-function-type diagram (c) if we replace \( H_d \) by \( L_k \), we will arrive at the familiar diagrams of the standard leptogenesis scenario, see e.g. [21]. Diagrams (b) and (c) involve a \( \Delta L = 2 \) Majorana mass insertion in the internal \( N_k \) line (\( N_k^* C N_k \) or \( F_{N_k} \tilde{N}_k \)). There is, however, a new diagram, (d), that does not have a counterpart in the standard R-parity conserving case. Notice that, in contrast to the \( N_k \) propagator in diagram (c), the one appearing in diagram (d) is lepton number conserving.

To leading order, we have

\[ \Gamma_{\text{tot}}(N_1) = \Gamma_{\text{tot}}(\tilde{N}_1) = \frac{(Y_{\nu} Y_{\nu}^T)_{ij} + |\lambda_i|^2}{4\pi} M_i \] (5)

so that

\[ \epsilon = \frac{1}{2\pi} \sum_{k \neq 1} \left[ g \left( M^2_k / M^2_1 \right) + \frac{2|\lambda_k|^2}{M^2_1} \right] \frac{\mathcal{I}_{k1} - \frac{2J_{k1}}{M^2_1}}{\sqrt{x} \ln[(1 + x)/x]}, \] (6)

where \( g(x) = \sqrt{x} \ln[(1 + x)/x] \),

\[ \mathcal{I}_{k1} = \sum_j \text{Im}[(Y_{\nu}^* Y_{\nu}^T)_{1j}] \lambda_k (Y_{\nu} Y_{\nu}^T)_{1j} + |\lambda_i|^2 \] (7)

and

\[ \mathcal{J}_{k1} = \sum_j \text{Im}[(Y_{\nu}^* Y_{\nu}^T)_{1j}] \lambda_k (Y_{\nu} Y_{\nu}^T)_{1j} + |\lambda_i|^2 \] (8)

Notice that the term proportional to \( \mathcal{J}_{k1} \) comes from the interference of the tree-level diagram with diagram (d).

Let us suppose \( M_1 < 10^6 \) GeV so that thermal production of \( N_1 \) and \( \tilde{N}_1 \) in the early universe can take place without requiring problematic very high reheating temperatures. Moreover let us suppose \( M_2 \) is not much heavier: \( M_2^2 / M_1^2 \sim 10 \). (Since the mechanism we are describing is effective with two right-handed neutrinos, here we only concentrate on \( N_1 \) and \( \tilde{N}_1 \) dropping \( N_2 \) from the discussion. In principle \( N_3 \) can play a similar role as \( N_2 \).) For these values of \( M_1 \), to suppress the masses of left-handed neutrinos in the simplest type-I seesaw down to \( \sqrt{\Delta m_{atm}^2} \), the Yukawa couplings have to be very tiny \( (Y_{\nu} Y_{\nu}^T)_{ij} \sim 10^{-5} M_1/10^6 \) GeV, similar to that of the electron in the Standard Model. In order for \( N_1 \) and \( \tilde{N}_1 \) to decay out of equilibrium (i.e., \( \Gamma_{\text{tot}}(N_1) = \Gamma_{\text{tot}}(\tilde{N}_1) < H |_{T=M_1} \), where \( H \) is the Hubble expansion rate) \( \lambda_1 \) must be also small: \( |\lambda_1|^2 \sim (Y_{\nu} Y_{\nu}^T)_{11} \).
In this range of parameters, 
\[ \epsilon_{N_1} + \epsilon_{\bar{N}_1} \approx 10^{-6} \sqrt{\frac{M_1}{10^6 \text{ GeV}}} \lambda_2 \sin \phi \] 
(9)

where \( \phi \) is the relevant CP-violating phase which can be of order of 1.

Now, let us discuss the wash-out processes. The evolution of the numbers of the relevant particles is given by the following Boltzman equations:

\[ \frac{dN_{N_1}}{dz} = -(D + S)(N_{N_1} - N_{eq}^{N_1}), \] 
(10)

and

\[ \frac{dN_{B-L}}{dz} = -\epsilon_{N_1} D(N_{N_1} - N_{eq}^{N_1}) - WN_{B-L}, \] 
(11)

where \( N_{N_1} \) is the number of \( N_1 \) plus that of its superpartner and \( N_{B-L} \) denotes the baryon number minus the number of standard model leptons [not including the right-handed (s)neutrinos]. In the above equations, \( z = M_1/T \) and \( D \) and \( S \) respectively represent the rates of the decay and scattering of \( N_1 \) and \( \bar{N}_1 \) (\( D = \Gamma_D/H \) and \( S = \Gamma_S/H \)). \( W \equiv \Gamma_W/H \) in Eq. (10) represents the rate of processes that erase the produced \( B-L \). Here, since the rates of interactions of \( N_1 \) and its superpartner are the same, it is not necessary to consider the evolution of the number of \( N_1 \) and its superpartner, separately. Moreover, writing (11), we have used \( \epsilon_{N_1} = \epsilon_{\bar{N}_1} \).

In the R-parity conserving case, it is shown that the dependence of \( S \) and \( D \) on the seesaw parameters is through the combination

\[ \bar{m}_1 = \frac{(Y_{\nu}^\dagger Y_{\nu})_{11} v^2}{M_1}. \]

In the presence of the new interaction, there are new diagrams contributing to both decay and scattering of right-handed (s)neutrinos and the definition of \( \bar{m}_1 \) has to be modified to

\[ \bar{m}_1 = \frac{(Y_{\nu}^\dagger Y_{\nu})_{11} + |\lambda_1|^2 + v^2}{M_1}. \]

We can divide the processes that contribute to \( W \) into three categories: i) Inverse decay processes of type \( tH_u \to N_1 \) and scattering such as \( N_1 \ell \to \bar{t}q \), that involve \( N_1 \) or \( \bar{N}_1 \) along with a standard model lepton. The dependence of the rate of these processes on seesaw parameters is through the combination \( (Y_{\nu}^\dagger Y_{\nu})_{11}/M_1 \); ii) R-parity conserving \( \Delta L = 2 \) processes such as \( \ell \ell \to H_u H_u \). The rate of these processes \( (W^R) \) can be written as

\[ W^R = aM_1 \left( Y_{\nu}^\dagger Y_{\nu} \right)^2. \]

Here, \( a \) does not depend on the seesaw parameters. iii) Turning on the R-parity violating couplings, new processes will take place that contribute to wash-out of the produced \( B - L \). The new processes are of type \( \ell H \to HH \) where \( H \) collectively denotes \( H_u, H_d \) and their superpartners. Such processes take place through virtual \( N_2 \) or \( \bar{N}_2 \) exchange and their rate can be estimated as

\[ \sim aM_1 \left| \lambda_2 (Y_{\nu}^2)_{2i}/M_2 \right|^2. \]

Summing up the above discussion and remembering \( |\lambda_1|^2 \sim (Y_{\nu}^\dagger Y_{\nu})_{11} \), we conclude that, for a given value of \( \bar{m}_1 \), the effect of the wash-out processes in our scenario is similar to the R-parity conserving case provided that we replace \( M_1 \) with \( M_1 \left| \lambda_2 (Y_{\nu})_{2i} \right|^2 \). As a result, the wash-out factor for \( \lambda_2 \sim 0 \) and \( M_1 \sim 10^6 \) GeV (which results in \( \epsilon_{N_1} \sim 10^{-6} \)) will be of order of the wash-out factor for \( \lambda_2 \sim 0 \) and \( M_1 \sim 10^{15} \) GeV.

The wash-out factor for the latter case is known. In [23], it is shown that for \( \lambda_1 = 0 \), \( M_1 \sim 10^{15} \) GeV and \( \bar{m}_1 \sim 10^{-5} \) eV, an asymmetry \( \epsilon_{N_1} \) of \( 10^{-6} \) is enough to explain the baryon asymmetry of the universe provided that the initial number of the lightest right-handed neutrino is thermal. Equivalently, we conclude that for \( \lambda_2 \sim 1 \), \( |\lambda_1|^2 \sim (Y_{\nu}^\dagger Y_{\nu})_{11} \) and \( M_1 \sim 10^6 \) GeV, there is a range of parameters [corresponding to \( v^2 ((\lambda_1)^2 + (Y_{\nu})_{11}^\dagger) /M_1 \sim 10^{-5} \) eV] for which, through the scenario discussed in this letter, the observed baryon asymmetry of the universe can be produced. Notice that unlike the case \( \lambda_1 = 0 \) and \( M_1 \sim 10^{15} \) GeV, the thermal production of the right-handed (s)neutrinos in our scenario can be realized with safely low reheating temperature.

There is a subtlety here that should be noticed. As shown in [23], for small values of \( (Y_{\nu})_{11} \), which correspond to \( \bar{m}_1 \sim 10^{-5} \) eV, even if the reheating temperature is above \( M_1 \), the rates of interactions that involve \( Y_{\nu} \) will not be high enough to give rise to a thermal initial number of \( N_1 \). In order to have a thermal initial number of \( N_1 \) and \( \bar{N}_1 \), there has to be “another” mechanism for production of right-handed (s)neutrinos. In our scenario, there is a natural mechanism to create initial thermal distribution which we briefly discuss below. Since in our scenario \( \lambda_2 \sim 1 \), \( N_2 \) and \( \bar{N}_2 \) maintain their thermal equilibrium and consequently for temperatures below \( M_1 \) their numbers are negligible. As a result an interaction of type

\[ W^{N_3} = \lambda_2^{N_3} \bar{N}_2 N_2 N_1 \] 

(12)
cannot contribute to the wash-out. However, at \( T \sim M_2 \) (remember that we have assumed that \( |M_2 - M_1|/M_1 \sim 1 \) so to have \( T \sim M_2 \) the reheating temperature does not need to be far higher than \( M_1 \) the term in (12) can contribute to the production of \( N_1 \) and \( \bar{N}_1 \), giving rise to a thermal distribution of \( N_1 \) and \( \bar{N}_1 \) at \( T > M_1 \). As a result, in this framework assuming a thermal initial number of \( N_1 \) and \( \bar{N}_1 \), even for very small values of \( (Y_{\nu})_{11} \), is reasonable.

Determining the exact range of allowed parameters requires detailed numerical calculation of the wash-out ef-
fects which is beyond the scope of this letter and will be presented elsewhere.

Before concluding we note that for $\lambda_2 \sim 1$ the new term can significantly affect the renormalization group equations of the Higgs sector which may have consequences for radiative electroweak symmetry breaking. In addition, the new term also slightly shifts the vacuum expectation values of the scalars of the theory. The $N$-dependent part of the scalar potential is

$$
\sum_i |M_i \tilde{N}_i + (Y_{ij})_i \tilde{L}_j H_u + \lambda_i H_d H_u + \sum_{jk} \lambda_{ijk} \tilde{N}_j \tilde{N}_k|^2 + 
$$

$$
\sum_{ij} \left[ m_{\tilde{N}_i}^2 |\tilde{N}_i|^2 + B_i M_i (\tilde{N}_i^2 + H.c.) 
+ |A_i \tilde{N}_i H_d H_u + (A_{ij})_i \tilde{N}_j \tilde{L}_j H_u + H.c.| \right] 
$$

(13)

Because of the new term the right-handed sneutrinos develop very small vacuum expectation values:

$$
\langle \tilde{N}_i \rangle \sim \lambda_i \nu^2 \sin \beta \cos \beta \langle M_i \rangle \ll \nu.
$$

Expanding the superpotential around $\langle \tilde{N}_i \rangle$ we obtain a tiny correction to the $\mu$ term. Moreover, we obtain the following bilinear R-parity violating term

$$
\sum_i (Y_{ij})_i \epsilon \alpha \beta \lambda_i \nu^2 \sin \beta \cos \beta \langle M_i \rangle L_j \tilde{H}_u \beta.
$$

(14)

As discussed in [24], such term gives rise to the decay of lightest neutralino. For the specific parameter range studied in this letter, we can make the following estimate

$$
\Gamma(\chi \rightarrow \nu_i + e^- + e^+) \sim 10^{-11} \lambda_2^2 \times \frac{\cos^2 \beta}{0.01} \frac{\text{Max}[|Y_{ij}|^2, |Y_{ij}|^2]}{10^{-9} M_2} \left( \frac{10^6 \text{GeV}}{m_\chi} \right)^2 \left( \frac{m_\chi}{100 \text{GeV}} \right)^3.
$$

This implies that, if produced at colliders, like the Large Hadron Collider, the lightest neutralino would leave the detector before decaying, leading to the same missing energy signature as in the minimal supersymmetric standard model. However, in our model the lightest neutralino typically decays before the epoch of nucleosynthesis. Thus it cannot serve as dark matter, which needs another candidate, like the axion. The neutrino mass induced by Eq. (14) will be of order $(Y_{ij}\nu^2 \sin \beta \cos \beta/M)^2/m_{\text{nu}}$, completely negligible in comparison with the seesaw effect, $Y_{ij} \nu^2 \sin \beta \cos \beta/M$.

In conclusion we have suggested a simple variant of the supersymmetric seesaw mechanism where the thermal leptogenesis is assisted by an explicit R-parity violating term involving the heavy right-handed neutrino supermultiplets, $N_i$. In this scenario, the lightest right-handed neutrino ($\tilde{N}_1$) and its superpartner ($\tilde{N}_1$) could be as light as $10^6$ GeV or less and still account for the baryon asymmetry of the universe, avoiding overproduction of gravitinos which plagues the R-parity conserving thermal leptogenesis.

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