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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 $\sim 10^9$ 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 $\lambda_i \hat{N}_i \hat{H}_u \hat{H}_d$ term in the superpotential, where $\hat{N}_i$ 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 $\hat{N}_1$ 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, $M_3 > M_2 > M_1 \gg m_{weak}$.

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 $\sim 10^9$ GeV [9, 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

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

where $\hat{L}_j$ is the superfield associated with the left-handed lepton doublet $(\hat{\nu}_j, \hat{l}_j)$ and $\hat{H}_u$ 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

$$W_{\rm RPV} = \sum_i \epsilon_{\alpha\beta} \lambda_i \hat{N}_i \hat{H}_d^\alpha \hat{H}_u^\beta \quad (2)$$

where $\hat{H}_d$ 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 μ problem [18]. In our case its contribution to generating the μ term is negligible because $\tilde{N}_i$, 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_ν and λ can in general be complex. Introduction of the coupling λ_i adds three extra CP-violating phases to the theory, as we shall see, with consequences for the baryon asymmetry of the universe.

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Let us define the following asymmetries

$$\begin{aligned}\epsilon_{N_1} &= -\sum_i \left[\frac{\Gamma(N_1 \to \bar{l}_i \bar{H}_u) - \Gamma(N_1 \to l_i H_u)}{\Gamma_{\rm tot}(N_1)/2}\right. \\ &+ \left.\frac{\Gamma(N_1 \to \bar{\tilde{l}}_i \bar{\tilde{H}}_u) - \Gamma(N_1 \to \tilde{l}_i \tilde{H}_u)}{\Gamma_{\rm tot}(N_1)/2}\right] \end{aligned} \quad (3)$$

and

$$\begin{aligned}\epsilon_{\tilde{N}_1} &= -\sum_i \left[\frac{\Gamma(\tilde{N}_1 \to \bar{l}_i \bar{\tilde{H}}_u) - \Gamma(\tilde{N}_1^* \to l_i \tilde{H}_u)}{\Gamma_{\rm tot}(\tilde{N}_1)/2}\right. \\ &+ \left.\frac{\Gamma(\tilde{N}_1^* \to \bar{\tilde{l}}_i \bar{H}_u) - \Gamma(\tilde{N}_1 \to \tilde{l}_i H_u)}{\Gamma_{\rm tot}(\tilde{N}_1)/2}\right] \end{aligned} \quad (4)$$

where N_1 and $\tilde{N}_1$ are respectively the lightest right-handed neutrino and sneutrino and $\Gamma_{\rm tot}(N_1)$ and $\Gamma_{\rm tot}(\tilde{N}_1)$ 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 ϵ_{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 $(Y_\nu)_{ij} \lesssim \sqrt{(\Delta m^2_{atm})^{1/2} M_i/(v^2 \sin^2\beta)} \sim 10^{-5}\sqrt{M_i/(10^6 \text{ GeV})}$, where $v = 245$ 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 ϵ_{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 Yukawa vertices while those marked with $\otimes$ are the new R-parity violating vertices given by λ_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 leptogensis scenario, see e.g. [21]. Diagrams (b) and (c) involve a $\Delta L = 2$ Majorana mass insertion in the internal N_k line ($N_k^T C N_k$ 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

$$\Gamma_{\rm tot}(N_i) = \Gamma_{\rm tot}(\tilde{N}_i) = \frac{(Y_\nu Y_\nu^\dagger)_{ii} + |\lambda_i|^2}{4\pi} M_i \quad (5)$$

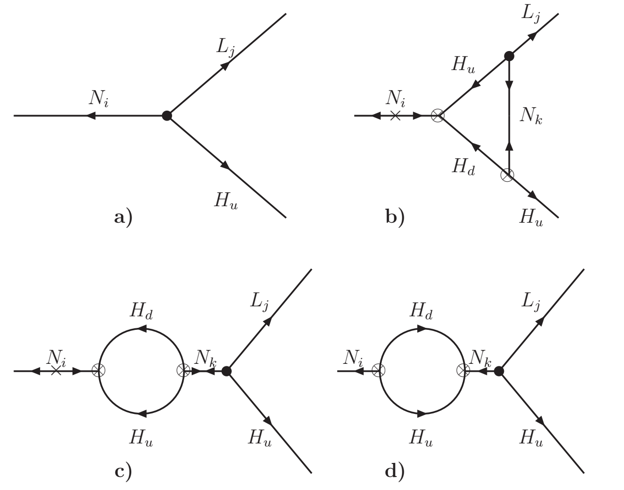

FIG. 1: Diagrams contributing to lepton-antilepton asymmetry. Vertices marked with dots and $\otimes$ denote Yukawa (Y_ν) and R-parity violating (λ) couplings, respectively. The $\times$ indicates mass insertion.

so that

$$\epsilon = \frac{1}{2\pi}\sum_{k\neq 1}\left[[g(\frac{M_k^2}{M_1^2}) + \frac{2\frac{M_k}{M_1}}{\frac{M_k^2}{M_1^2}-1}]\mathcal{I}_{k1} - \frac{2\mathcal{J}_{k1}}{\frac{M_k^2}{M_1^2}-1}\right], \quad (6)$$

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

$$\mathcal{I}_{k1} = \frac{\sum_j {\rm Im}[(Y_\nu^*)_{1j}\lambda_1^*\lambda_k (Y_\nu)_{kj}]}{(Y_\nu Y_\nu^\dagger)_{11} + |\lambda_1|^2} \quad (7)$$

and

$$\mathcal{J}_{k1} = \frac{\sum_j {\rm Im}[(Y_\nu^*)_{1j}\lambda_1\lambda_k^* (Y_\nu)_{kj}]}{(Y_\nu Y_\nu^\dagger)_{11} + |\lambda_1|^2}. \quad (8)$$

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 [11]. 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 N_2 dropping N_3 from the discussion. In principle N_3 can play a similar role as N_2.) For these values of M_i, to suppress the masses of left-handed neutrinos in the simplest type-I seesaw down to $\sqrt{\Delta m^2_{atm}}$, the Yukawa couplings have to be very tiny $(Y_\nu)_{ij} \lesssim 10^{-5}\sqrt{M_i/10^6 \text{ GeV}}$, 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., $\Gamma_{\rm tot}(N_1) = \Gamma_{\rm tot}(\tilde{N}_1) < H|_{T=M_1}$, where H is the Hubble expansion rate) λ_1 must be also small: $|\lambda_1|^2 \sim (Y_\nu Y_\nu^\dagger)_{11}$.

However the decay of the heavier (s)neutrinos does not need to be out of equilibrium, so that $\lambda_2 \sim 1$ is allowed. In this range of parameters,

$$\epsilon_{N_1} + \epsilon_{\tilde{N}_1} \approx 10^{-6}\sqrt{\frac{M_1}{10^6 \text{ GeV}}}\lambda_2 \sin\phi \qquad (9)$$

where ϕ 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:

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

and

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

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 $\tilde{N}_1^{(*)}$ ($D = \Gamma_D/Hz$ and $S = \Gamma_S/Hz$). $W \equiv \Gamma_W/Hz$ in Eq. (11) 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 [22]. Moreover, writing (11), we have used $\epsilon_{N_1} = \epsilon_{\tilde{N}_1}$.

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

$$\tilde{m}_1 = \frac{(Y_\nu Y_\nu^\dagger)_{11} v^2}{M_1}.$$

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 $\tilde{m}_1$ has to be modified to

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

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

$$W^R = aM_1 \left(Y_\nu^\dagger \frac{1}{M} Y_\nu \right)^2.$$

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 $\tilde{N}_2$ exchange and their rate can be estimated as

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

Summing up the above discussion and remembering $|\lambda_1|^2 \sim (Y_\nu Y_\nu^\dagger)_{11}$, we conclude that, for a given value of $\tilde{m}_1$, 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 $M_1[\lambda_2/(Y_\nu)_{2i}]^2$. As a result, the wash-out factor for $\lambda_1 \sim (Y_\nu)_{1i}$, $\lambda_2 \sim 1$ 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 $\lambda_2 = 0$ and $M_1 \sim 10^{15}$ GeV. The wash-out factor for the latter case is known. In [23], it is shown that for $\lambda_i = 0$, $M_1 \sim 10^{15}$ GeV and $\tilde{m}_1 \sim 10^{-5}$ eV, an asymmetry (ϵ_{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 $\lambda_2 \sim 1$, $|\lambda_1|^2 \sim (Y_\nu Y_\nu^\dagger)_{11}$ and $M_1 \sim 10^6$ GeV, there is a range of parameters [corresponding to $v^2 \left(|\lambda_1|^2 + (Y_\nu Y_\nu^\dagger)_{11} \right)/M_1 \sim 10^{-5}$ 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 $\lambda_i = 0$ and $M_1 = 10^{15}$ 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 $(Y_\nu)_{1i}$ which correspond to $\tilde{m}_1 \lesssim 10^{-5}$ eV, even if the reheating temperature is above M_1, the rates of interactions that involve Y_ν 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 $\tilde{N}_1$, 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 $\tilde{N}_2$ maintain their thermal equilibrium and consequently for temperatures below M_1 their numbers are negligible. As a result an interaction of type

$$W^{N^3} = \lambda_{221}^{N^3} \hat{N}_2 \hat{N}_2 \hat{N}_1 \qquad (12)$$

cannot contribute to the wash-out. However, at $T \gtrsim M_2$ (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 $\tilde{N}_1$, giving rise to a thermal distribution of N_1 and $\tilde{N}_1$ at $T > M_1$. As a result, in this framework assuming a thermal initial number of N_1 and $\tilde{N}_1$, even for very small values of $(Y_\nu)_{1i}$ [23], is reasonable.

Determining the exact range of allowed parameters requires detailed numerical calculation of the wash-out effects 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 $\tilde{N}$-dependant part of the scalar potential is

$$\sum_i |M_i \tilde{N}_i + (Y_\nu)_{ij} \tilde{L}_j H_u + \lambda_i H_d H_u + \sum_{jk} \lambda^{N^3}_{ijk} \tilde{N}_j \tilde{N}_k|^2 + \sum_{ij} \Big[m_0^2 |\tilde{N}_i|^2 + B_\nu M_i (\tilde{N}_i^2 + \text{H.c.}) + [A^i_\lambda \tilde{N}_i H_d H_u + (A_\nu)_{ij} \tilde{N}_i \tilde{L}_j H_u + \text{H.c.}] \Big] \quad (13)$$

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

$$\left| \left\langle \tilde{N}_i \right\rangle \right| \simeq \frac{\lambda_i v^2 \sin\beta \cos\beta}{M_i} \ll v.$$

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

$$\sum_i (Y_\nu)_{ij} \epsilon_{\alpha\beta} \frac{\lambda_i v^2 \sin\beta \cos\beta}{M_i} \hat{L}^\alpha_j \hat{H}^\beta_u. \quad (14)$$

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

$$\Gamma(\chi \to \nu_i + e^- + e^+) \sim 10 \text{ sec}^{-1}\ \lambda_2^2 \times \frac{\cos^2\beta}{0.01} \frac{\text{Max}[(Y_\nu)^2_{2i}, (Y_\nu)^2_{2e}]}{10^{-9}} \left(\frac{10^6 \text{GeV}}{M_2} \right)^2 \left(\frac{m_\chi}{100\text{GeV}} \right)^3.$$

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 $(Y_\nu \lambda v^2 \sin\beta \cos\beta / M)^2 / m_{susy}$[19], completely negligible in comparison with the seesaw effect, $Y_\nu^2 v^2 \sin^2\beta / M$.

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 (N_1) and its superpartner ($\tilde{N}_1$) could be as light as $\sim 10^6$ 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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