Leptogenesis from R Parity Nonconservation

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Abstract

It is known that realistic neutrino masses for neutrino oscillations may be obtained from R parity nonconserving supersymmetry. It is also known that such interactions would erase any preexisting lepton or baryon asymmetry of the Universe because of the inevitable intervention of the electroweak sphalerons. We now show how a crucial subset of these R parity nonconserving terms may in fact create its own successful leptogenesis.

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To obtain naturally small Majorana neutrino masses[1], the canonical approach is to have heavy singlet right-handed neutrinos in the guise of the famous seesaw mechanism[2]. An equally simple alternative is to use a heavy scalar Higgs triplet[3]. Both offer the important additional benefit of generating a lepton asymmetry[3, 4] which gets converted into the present observed baryon asymmetry of the Universe by virtue of the electroweak sphalerons[5, 6].

In the past several years, a third way has come under active consideration, namely that of R parity nonconserving supersymmetry. Whereas realistic neutrino masses for neutrino oscillations may be generated[7], the lepton-number nonconserving interactions involved would certainly erase any preexisting lepton asymmetry of the Universe[8]. This negative aspect of R parity nonconserving supersymmetry has prevented it from being considered as a truly competitive alternative to the canonical seesaw or the triplet Higgs mechanism for neutrino masses. To remedy this situation, we now show for the first time how a crucial subset of these R parity nonconserving terms may in fact create its own successful leptogenesis, through the suppressed decay of the lightest neutralino into a charged Higgs boson and a lepton.

In a supersymmetric extension of the minimal standard model of fundamental interactions, the R parity of a particle is conventionally defined as

$$\mathrm{R} \equiv (-1)^{3B+L+2J}, \tag{1}$$

where B is its baryon number, L its lepton number, and J its spin angular momentum. Hence the standard-model particles have R $= +1$ and their supersymmetric partners have R $= -1$. Using the common notation where all chiral superfields are considered left-handed, the three families of leptons and quarks are given by

$$L_i = (\nu_i, e_i) \sim (1, 2, -1/2), \quad e_i^c \sim (1, 1, 1), \tag{2}$$

$$Q_i = (u_i, d_i) \sim (3, 2, 1/6), \tag{3}$$

$$u_i^c \sim (3^*, 1, -2/3), \quad d_i^c \sim (3^*, 1, 1/3), \tag{4}$$

where i is the family index, and the two Higgs doublets are given by

$$H_1 = (h_1^0, h_1^-) \sim (1, 2, -1/2), \tag{5}$$

$$H_2 = (h_2^+, h_2^0) \sim (1, 2, 1/2), \tag{6}$$

where the $SU(3)_C \times SU(2)_L \times U(1)_Y$ content of each superfield is also indicated. If R parity is conserved, the superpotential is restricted to have only the terms

$$\begin{aligned} W &= \mu H_1 H_2 + f_{ij}^e H_1 L_i e_j^c + f_{ij}^d H_1 Q_i d_j^c \\ &\quad + f_{ij}^u H_2 Q_i u_j^c. \end{aligned} \tag{7}$$

If R parity is violated but not baryon number, then the superpotential contains the additional terms

$$W' = \mu_i L_i H_2 + \lambda_{ijk} L_i L_j e_k^c + \lambda'_{ijk} L_i Q_j d_k^c, \tag{8}$$

resulting in nonzero neutrino masses either from mixing with the neutralino mass matrix or in one-loop order[7].

Consider the simplified case[9, 10] of having only L_τ and τ^c as effective $L = 0$ superfields. The neutralino mass matrix now spans five fields: the U(1) gaugino ($\tilde{B}$), the SU(2) gaugino ($\tilde{W}_3$), the two Higgsinos ($\tilde{h}_1^0$, $\tilde{h}_2^0$), and ν_τ.

$$\mathcal{M} = \begin{bmatrix} M_1 & 0 & -sm_3 & sm_4 & -sm_5 \\ 0 & M_2 & cm_3 & -cm_4 & cm_5 \\ -sm_3 & cm_3 & 0 & -\mu & 0 \\ sm_4 & -cm_4 & -\mu & 0 & -\mu_\tau \\ -sm_5 & cm_5 & 0 & -\mu_\tau & 0 \end{bmatrix}, \tag{9}$$

where $s = \sin\theta_W$, $c = \cos\theta_W$, $m_3 = M_Z \cos\beta \cos\theta_\tau$, $m_4 = M_Z \sin\beta$, $m_5 = M_Z \cos\beta \sin\theta_\tau$, with $\tan\beta = \langle h_2^0 \rangle / [\langle h_1^0 \rangle^2 + \langle \tilde{\nu}_\tau \rangle^2]^{1/2}$ and $\tan\theta_\tau = \langle \tilde{\nu}_\tau \rangle / \langle h_1^0 \rangle$.

To understand the structure of the above 5×5 mass matrix, let us assume that μ is the dominant term, then $\tilde{h}^0_{1,2}$ form a heavy Dirac particle of mass μ and the reduced 3×3 matrix in the basis $(\tilde{B}, \tilde{W}_3, \nu_\tau)$ becomes

$$\mathcal{M} = \begin{bmatrix} M_1 - s^2\delta & sc\delta & -s\epsilon \\ sc\delta & M_2 - c^2\delta & c\epsilon \\ -s\epsilon & c\epsilon & 0 \end{bmatrix}, \tag{10}$$

where

$$\delta = 2m_3 m_4/\mu = M_Z^2 \sin 2\beta \cos\theta_\tau/\mu, \tag{11}$$

$$\epsilon = m_5 - m_3\mu_\tau/\mu$$

$$= M_Z \cos\beta \cos\theta_\tau(\tan\theta_\tau - \mu_\tau/\mu). \tag{12}$$

From the above, ν_τ gets a seesaw mass, i.e.

$$m_{\nu_\tau} = -\epsilon^2\left(\frac{s^2}{M_1} + \frac{c^2}{M_2}\right), \tag{13}$$

corresponding to the mass eigenstate

$$\nu'_\tau = \nu_\tau + \frac{s\epsilon}{M_1}\tilde{B} - \frac{c\epsilon}{M_2}\tilde{W}_3, \tag{14}$$

whereas the other two mass eigenstates are

$$\tilde{B}' = \tilde{B} + \frac{sc\delta}{M_1 - M_2}\tilde{W}_3 - \frac{s\epsilon}{M_1}\nu_\tau, \tag{15}$$

$$\tilde{W}'_3 = \tilde{W}_3 - \frac{sc\delta}{M_1 - M_2}\tilde{B} + \frac{c\epsilon}{M_2}\nu_\tau. \tag{16}$$

Because $\tilde{B}'$ and $\tilde{W}'_3$ contain ν_τ, they are not stable, but would decay into $\tau^\mp W^\pm$ pairs. This may generate a lepton asymmetry, but it is several orders of magnitude too small because it has to be much less than $(\epsilon/M_{1,2})^2$, which is of order $m_{\nu_\tau}/M_{1,2}$, i.e. $< 5 \times 10^{-13}$ if $m_{\nu_\tau} < 0.05$ eV and $M_{1,2} > 100$ GeV.

Consider now the couplings of $\tilde{B}$ and $\tilde{W}_3$. The former couples to both $\bar{\tau}_L\tilde{\tau}_L$ and $\bar{\tau}^c_L\tilde{\tau}^c_L$, but the latter only to $\bar{\tau}_L\tilde{\tau}_L$, because τ^c_L is a singlet under $SU(2)_L$. Since R parity is violated,

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both $\tilde{\tau}_L$ and $\tilde{\tau}_L^c$ mix with the charged Higgs boson of the supersymmetric standard model: $h^\pm = h_2^\pm \cos\beta + h_1^\pm \sin\beta$. Hence $\tilde{B}'$ and $\tilde{W}_3'$ may decay into $\tau^\mp h^\pm$. We now assume that the $\tilde{\tau}_L$ mixing is negligible, so that the only relevant coupling is that of $\tilde{B}$ to $\bar{\tau}_L^c h^+$. Hence $\tilde{W}_3'$ decay (into $\tau^\mp h^\pm$) is suppressed because it may only do so through the small component of $\tilde{B}$ that it contains, assuming of course that $\tilde{\tau}_L$ and $\tilde{\tau}_L^c$ are heavier than $\tilde{B}$ or $\tilde{W}_3$.

We now envisage the following leptogenesis scenario. At temperatures well above $T = M_{SUSY}$, the presence of sphalerons and R parity nonconserving interactions together ensure that there is no L or B asymmetry. As the Universe cools down to a temperature below the masses of all supersymmetric particles except $\tilde{B}'$ and $\tilde{W}_3'$, we need only consider their interactions. In Figure 1 we show the lepton number violating processes (a) $\tilde{B}' \leftrightarrow \tau_R^\pm h^\mp$, where we have adopted the more conventional notation of an outgoing τ_R in place of an incoming τ_L^c. These are still certainly fast and there can be no L asymmetry. Let us assume now that $M_1 > M_2$, then for $T < M_1$, the $\tilde{B}'$ interactions are suppressed and we need only consider $\tilde{W}_3'$. However, as we have explained in the previous paragraph, the lepton-number violating processes (b) $\tilde{W}_3' \leftrightarrow \tau_R^\pm h^\mp$ are slow and can satisfy the out-of-equilibrium condition for generating a lepton asymmetry of the Universe, resulting from the interference of this tree-level diagram with the one-loop (c) self-energy and (d) vertex diagrams. Since the unsuppressed lepton number violating couplings of $\tilde{B}'$ are involved, a realistic lepton asymmetry may be generated. It is then converted by the still active sphalerons into the present observed baryon asymmetry of the Universe.

We start with the well-known interaction of $\tilde{B}$ with τ and $\tilde{\tau}_R$ given by[11]

$$ -\frac{e\sqrt{2}}{\cos\theta_W} \left[\bar{\tau} \left(\frac{1-\gamma_5}{2} \right) \tilde{B}\tilde{\tau}_R + H.c. \right]. \tag{17} $$

We then allow $\tilde{\tau}_R$ to mix with h^-, and $\tilde{B}$ to mix with $\tilde{W}_3$, so that the interaction of the

physical state $\tilde{W}_3'$ of Eq. (16) with τ and $h^\pm$ is given by

$$\left(\frac{sc\xi}{M_1 - M_2}\right)\left(\frac{e\sqrt{2}}{\cos\theta_W}\right)\left[\delta\bar{\tau}\left(\frac{1-\gamma_5}{2}\right)\tilde{W}_3'h^- + H.c.\right], \tag{18}$$

where ξ represents the $\tilde{\tau}_R - h^-$ mixing and is assumed real, but the parameter δ of Eq. (10) is now assumed complex. The origin of this nontrivial CP phase is from the 2×2 Majorana mass matrix spanning $\tilde{B}$ and $\tilde{W}_3$. It is well-known that the off-diagonal entry is in general complex here if the diagonal entries are chosen to be real.

To satisfy the out-of-equilibrium condition, we require[12]

$$\Gamma(\tilde{W}_3' \to \tau^\pm h^\mp) < H = 1.7\sqrt{g_*}(T^2/M_{Pl}) \tag{19}$$

at the temperature $T \sim M_2$, where H is the Hubble expansion rate of the Universe with g_* the effective number of massless degrees of freedom and M_{Pl} the Planck mass. This implies

$$\left(\frac{\xi|\delta|r}{M_1 - M_2}\right)^2 \frac{(M_2^2 - m_h^2)^2}{M_2^5} < 1.9 \times 10^{-14}\ \text{GeV}^{-1}, \tag{20}$$

where we have used $g_* = 10^2$ and $M_{Pl} = 10^{18}$ GeV, and $r = (1 + M_2/\mu \sin 2\beta)/(1 - M_2^2/\mu^2)$ is a correction factor for finite M_2/μ. To make sure that at $T \sim M_2$, the lepton number violating processes $\tau^\pm h^\mp \leftrightarrow \tau^\mp h^\pm$ through $\tilde{B}'$ exchange do not erase the lepton asymmetry created by the decay of $\tilde{W}_3'$, we require

$$\left(\frac{2e^2\xi^2}{\cos^2\theta_W}\right)^2 \frac{1}{M_1^2} \frac{T^3}{32\pi} \frac{f(x)}{(1-x)^2} < H \tag{21}$$

at $T \sim M_2$, where $x = M_2^2/M_1^2$ and

$$f(x) = 1 + \frac{2(1-x)}{x^2}[(1+3x)\ln(1+x) - x(1+x)], \tag{22}$$

which implies

$$\frac{\xi^4}{M_2}\frac{xf(x)}{(1-x)^2} < 2.6 \times 10^{-14}\ \text{GeV}^{-1}. \tag{23}$$

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Both conditions may be satisfied for example with the following choice of parameters:

$$\mu = 5 \text{ TeV}, \quad \sin 2\beta = 0.5, \quad \xi = 2 \times 10^{-3}, \tag{24}$$

$$M_1 = 3 \text{ TeV}, \quad M_2 = 2 \text{ TeV}, \quad m_h = 200 \text{ GeV}, \tag{25}$$

for which the left-hand sides of Eqs. (20) and (23) are 0.6×10^{-14} GeV^{-1} and 2.6×10^{-14} GeV^{-1} respectively.

The interference between the tree-level and self-energy + vertex diagrams results in the following lepton asymmetry in the decay of $\tilde{W}_3'$:

$$\epsilon = \frac{\alpha \xi^2}{2 \cos^2 \theta_W} \frac{Im\delta^2}{|\delta|^2} \left(1 - \frac{m_h^2}{M_2^2}\right)^2 \frac{x^{1/2} g(x)}{(1-x)}, \tag{26}$$

where

$$g(x) = 1 + \frac{2(1-x)}{x} \left[\left(\frac{1+x}{x}\right) \ln(1+x) - 1\right], \tag{27}$$

which yields 3.6×10^{-8} if $Im\delta^2/|\delta|^2 = 1$. Dividing by $3g_*$, we then obtain a realistic baryon-to-photon ratio of the order 10^{-10}.

In the above, we have assumed $M_{SUSY} \sim$ few TeV. This value is naturally suited for our scenario because (1) it allows ϵ in Eq. (26) to be large enough without contradicting the condition given by Eq. (21), and (2) it allows the lepton asymmetry to be converted at a temperature when the sphalerons are still very active[13, 14], into the present observed baryon asymmetry of the Universe. We have also assumed in our scenario that $\tilde{W}_3'$ only couples to $\tau^\pm h^\mp$ through its $\tilde{B}$ content. This means that $\tilde{\tau}_R - h^-$ mixing is significant, but $\tilde{\tau}_L - h^-$ mixing is negligible. The latter is related to the mixing of the neutrino with the neutralinos as given by Eq. (9) which is indeed small and may even be set to zero, because realistic neutrino masses may be obtained through R parity nonconserving trilinear couplings instead[7]. The former is usually taken[15] from the soft supersymmetry breaking term $h_1^- \tilde{\nu}_\tau \tilde{\tau}_L^c$ and the $\tilde{\tau}_L h_2^+$ term through $\tilde{\tau}_L^* - \tilde{\tau}_L^c$ mixing. In that case, the $\tilde{\tau}_R - h^-$ mixing

would be proportional to a linear combination of $\langle \tilde{\nu}_\tau \rangle$ and μ_τ, and be constrained to be very small[16]. For our scenario to be successful, we need to introduce the new term

$$H_2^\dagger H_1 \tilde{\tau}_L^c = (\bar{h}_2^+ h_1^0 + \bar{h}_2^0 h_1^-)\tilde{\tau}_L^c, \tag{28}$$

which is unconstrained and may have the desired mixing of $\xi = 2 \times 10^{-3}$ shown. Such a soft supersymmetry breaking term is nonholomorphic[17] and whereas this class of couplings has been studied previously[18], we are the first to make use of this particular term. We should also mention that although we have used only τ to represent a lepton, replacing it with e or μ is just as appropriate.

In conclusion, we have shown in this paper how R parity nonconservation may create its own successful leptogenesis through the suppressed decay of $\tilde{W}_3'$ into $l^\pm h^\mp$. We require $M_{SUSY} \sim$ few TeV and the existence of the nonholomorphic soft supersymmetry breaking term $H_2^\dagger H_1 \tilde{l}_L^c$. Whereas this mechanism depends on R parity nonconservation, it comes from a sector distinct from that responsible for nonzero neutrino masses[19]. We also require the soft supersymmetry breaking gaugino masses M_1 and M_2 to be a few TeV with $M_1 > M_2$.

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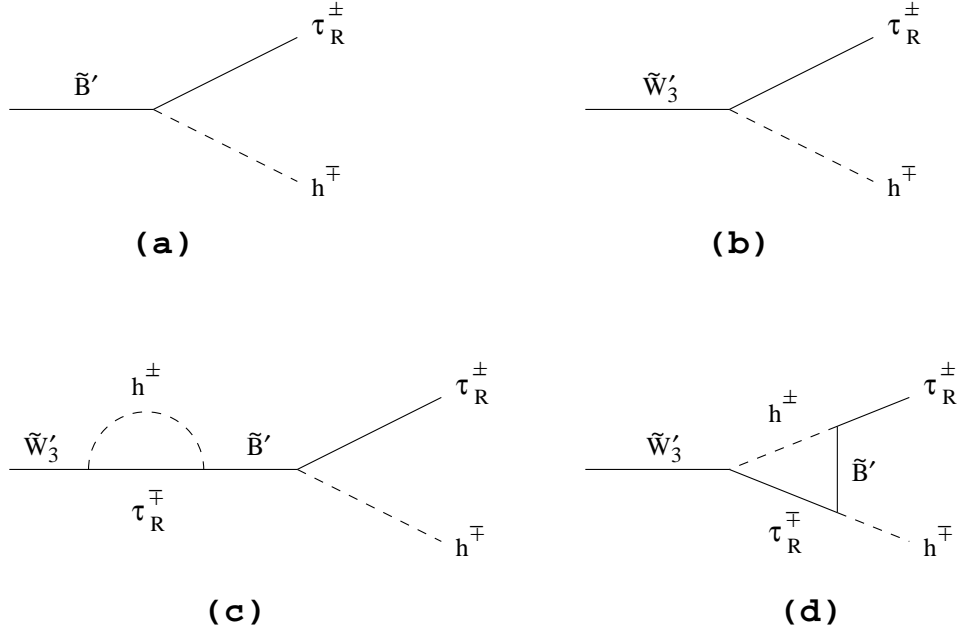
Figure 1: Tree-level diagrams for (a) $\tilde{B}'$ decay and (b) $\tilde{W}'_3$ decay (through their $\tilde{B}$ content), and the one-loop (c) self-energy and (d) vertex diagrams for $\tilde{W}'_3$ decay which have absorptive parts of opposite lepton number.