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Leptogenesis from Neutralino Decay with Nonholomorphic R-Parity Violation

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Abstract

In supersymmetric models with lepton-number violation, hence also R-parity violation, it is easy to have realistic neutrino masses, but then leptogenesis becomes difficult to achieve. After explaining the general problems involved, we study the details of a model which escapes these constraints and generates a lepton asymmetry, which gets converted into the present observed baryon asymmetry of the Universe through the electroweak sphalerons. This model requires the presence of certain nonholomorphic R-parity violating terms. For completeness we also present the most general R-parity violating Lagrangian with soft nonholomorphic terms and study their consequences for the charged-scalar mass matrix. New contributions to neutrino masses in this scenario are discussed.

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1 Introduction

The creation of a lepton asymmetry, i.e. leptogenesis [1, 2, 3], which gets converted into the present observed baryon asymmetry of the Universe, is closely related to the mechanism by which neutrinos obtain mass. In general, all models of Majorana neutrino masses with the same low-energy particle content as that of the standard model are equivalent in the sense that they are all characterized by the same nonrenormalizable dimension-five operator [4] $\Lambda^{-1}\nu_i\nu_j\phi^0\phi^0$. Different models of neutrino mass are merely different realizations [5] of this operator. They become distinguishable only at high energies, and since their interactions must violate lepton number, leptogenesis is a very natural possibility. For the canonical seesaw mechanism [6] and the Higgs triplet model [7], leptogenesis does indeed occur naturally [1, 7]. On the other hand, if neutrino masses are obtained radiatively [5], not only is leptogenesis difficult to achieve, the mechanism by which the former is accomplished leads naturally to the erasure of any primordial baryon asymmetry of the Universe [8, 9, 10]. This is especially true in supersymmetric models of neutrino mass [11] with R-parity violation. In a recent article [12], it was pointed out that leptogenesis is still possible in this case, provided that certain conditions regarding the R-parity violating terms are satisfied. Here we study this model in detail.

In Section 2 we write down the superpotential of the lepton-number violating (but baryon-number conserving) extension of the supersymmetric standard model, together with all possible soft supersymmetric breaking terms, including the nonholomorphic terms [13]. In Section 3 we consider bilinear R-parity violation and how leptogenesis is related to neutrino mass in this limited scenario. We find it to be negligible for realistic values of m_ν. In Section 4 we discuss how leptogenesis may occur without being constrained by neutrino mass in an expanded scenario. We assume negligible (enhanced) mixing between doublet (singlet) sleptons and charged Higgs bosons by allowing nonholomorphic soft supersymmetry breaking terms. In Section 5 we present the details of our calculations using the Boltzmann equations for obtaining the eventual lepton asymmetry. In Section 6, the complete charged-scalar mass matrix is displayed and analyzed. In Section 7, a new two-loop mechanism for neutrino mass is proposed. Finally in Section 8, there are some concluding remarks.

2 Superpotential and Soft Supersymmetry Breaking

In an unrestricted supersymmetric extension of the standard model of particle interactions, the chiral scalar superfields allow baryon-number violating terms which are not necessarily suppressed. These dangerous terms are usually avoided by assuming a conserved discrete quantum number for each particle called R-parity, which is 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. With this definition, the standard-model particles have $\mathrm{R} = +1$ and their supersymmetric partners have $\mathrm{R} = -1$. We can list the three families of leptons and quarks of the standard model using the notation where all superfields are considered left-handed:

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

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

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

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

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

In this case, both baryon and lepton numbers are conserved. However, to forbid proton decay, it is sufficient to conserve either baryon number or lepton number (because the final state of the proton decay must contain a lepton or antilepton). If only baryon number or only lepton number is violated (thus R-parity is also violated), the conservation of the other quantum number is enough to satisfy all present experimental constraints. This has motivated numerous studies of R-parity violating models.

If R-parity is violated, the superpotential contains the additional terms

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

We cannot have all of these terms because then the proton will decay very quickly. We may choose only the lepton-number violating terms or only the baryon-number violating terms. Following the overwhelming choice of many others, we consider only the former case and set $\lambda''_{ijk} = 0$. The remaining terms may now induce nonzero neutrino masses, either from mixing with the neutralino mass matrix, or in one-loop order[11]. Although these terms are allowed, we do not know if they originate from any fundamental theory, so the couplings are considered free parameters, constrained only by experiment.

Other free parameters exist in the minimal supersymmetric standard model (MSSM), i.e. the soft supersymmetry breaking terms, which do not introduce quadratic divergences to the renormalized theory. Usually only the holomorphic terms are considered which come from the chiral superpotential interacting with gravity, together with the gaugino masses. The most general such Lagrangian conserving R-parity is:

$$\begin{aligned}\mathcal{L}_{soft} &= -\tilde{L}_i^{a*}(M_L^2)_{ij}\tilde{L}_j^a - \tilde{e}_i^{c*}(M_e^2)_{ij}\tilde{e}_j^c - \tilde{Q}_i^{a*}(M_Q^2)_{ij}\tilde{Q}_j^a - \tilde{u}_i^{c*}(M_u^2)_{ij}\tilde{u}_j^c \\ &\quad -\tilde{d}_i^{c*}(M_d^2)_{ij}\tilde{d}_j^c - M_{H_1}^2 H_1^{a*}H_1^a - M_{H_2}^2 H_2^{a*}H_2^a - \varepsilon_{ab}(BH_1^aH_2^b + h.c.) \\ &\quad -\varepsilon_{ab}\Big((A_ef_e)_{ij}H_1^a\tilde{L}_i^b\tilde{e}_j^c + (A_uf_u)_{ij}H_2^b\tilde{Q}_i^a\tilde{u}_j^c + (A_df_d)_{ij}H_1^a\tilde{Q}_i^b\tilde{d}_j^c \,+\, h.c.\Big) \\ &\quad -\frac{1}{2}\Big(M_3\tilde{g}\tilde{g} + M_2\tilde{W}\tilde{W} + M_1\tilde{B}\tilde{B} \,+\, h.c.\Big)\,. \end{aligned} \tag{9}$$

If R-parity is violated, more soft terms may be present, i.e.

$$\mathcal{L}_{soft}^{\not R} = -\varepsilon_{ab}(B_i'\tilde{L}_i^aH_2^b + A_{ijk}'^e\tilde{L}_i^a\tilde{L}_j^b\tilde{e}_k^c + A_{ijk}'^d\tilde{Q}_i^a\tilde{L}_j^b\tilde{d}_k^c) - A_{ijk}'^S\tilde{u}_i^c\tilde{d}_j^c\tilde{d}_k^c \,+\, h.c.\,. \tag{10}$$

We follow the convention that the coupling constants of all the R-parity conserving soft terms are denoted without a prime, while the R-parity violating terms are denoted with a prime.

Since the soft terms may originate from gravity couplings, there is no clear reason for them to come only from the renormalizable chiral superpotential. They may also originate from nonrenormalizable terms, which can be functions of both left and right chiral superfields. Such nonholomorphic terms allow new supersymmetry-breaking soft terms [13] in the Lagrangian. The most general set of nonholomorphic soft terms conserving R-parity is:

$$\mathcal{L}_{soft}^{NH} = -N_{ij}^eH_2^{a*}\tilde{L}_i^a\tilde{e}_j^c - N_{ij}^dH_2^{a*}\tilde{Q}_i^a\tilde{d}_j^c - N_{ij}^uH_1^{a*}\tilde{Q}_i^a\tilde{u}_j^c \,+\, h.c.\,. \tag{11}$$

Similarly, nonholomorphic soft terms breaking R-parity are:

$$\begin{aligned}\mathcal{L}_{soft}^{NH\not R} &= -N_i'^BH_1^{a*}\tilde{L}_i^a - N_i'^eH_2^{a*}H_1^a\tilde{e}_i^c - N_{ijk}'^u\tilde{L}_i^{a*}\tilde{Q}_j^a\tilde{u}_k^c \\ &\quad -N_{ijk}'^S\tilde{u}_i^c\tilde{e}_j^c\tilde{d}_k^{c*} - N_{ijk}'^d\varepsilon_{ab}\tilde{Q}_i^a\tilde{Q}_j^b\tilde{d}_k^{c*} \,+\, h.c. \end{aligned} \tag{12}$$

In our convention, all "N" constants are for nonholomorphic terms. Since we assume lepton-number violation but baryon-number conservation, it implies $\lambda''_{ijk} = A_{ijk}'^S = N_{ijk}'^d = 0$.

3 Bilinear R-Parity Violation

Before coming to the explicit model in the next section, we first look at the possibility of generating a lepton asymmetry from bilinear R-parity violating terms. This case is a good illustration of the general problems involved. One immediate consequence of the violation of lepton number through the bilinear terms is the mixing of the neutrinos with the neutralinos. In the MSSM there are four neutralinos, the U(1) gaugino ($\tilde{B}$), the SU(2) gaugino ($\tilde{W}_3$), and the two Higgsinos ($\tilde{h}_1^0$, $\tilde{h}_2^0$). When lepton number is violated through the R-parity violating terms, it is possible to assign zero lepton number to what we usually regard as lepton superfields [14, 15]. Suppose only the τ neutrino mixes with the neutralinos, then both τ and ν_τ may be assigned an effective vanishing lepton number while the other leptons remain leptons. However, in the general three-family case, all three neutrinos may mix with the neutralinos, and the scalar partners of the neutrinos ($\tilde{\nu}_i$) may all acquire nonzero vacuum expectation values (VEVs).

In the most general case the neutralino mass matrix with all the seven fields in the basis $\left[\; \tilde{B}, \;\; \tilde{W}_3, \;\; \tilde{h}_1^0, \;\; \tilde{h}_2^0, \;\; \nu_1, \;\; \nu_2, \;\; \nu_3 \;\right]$ is given by

$$\mathcal{M} = \begin{bmatrix} M_1 & 0 & -sr_Zv_1 & sr_Zv_2 & -sr_Zv_{\nu_1} & -sr_Zv_{\nu_2} & -sr_Zv_{\nu_3} \\ 0 & M_2 & cr_Zv_1 & -cr_Zv_2 & cr_Zv_{\nu_1} & cr_Zv_{\nu_2} & cr_Zv_{\nu_3} \\ -sr_Zv_1 & cr_Zv_1 & 0 & -\mu & 0 & 0 & 0 \\ sr_Zv_2 & -cr_Zv_2 & -\mu & 0 & -\mu_1 & -\mu_2 & -\mu_3 \\ -sr_Zv_{\nu_1} & cr_Zv_{\nu_1} & 0 & -\mu_1 & 0 & 0 & 0 \\ -sr_Zv_{\nu_2} & cr_Zv_{\nu_2} & 0 & -\mu_2 & 0 & 0 & 0 \\ -sr_Zv_{\nu_3} & cr_Zv_{\nu_3} & 0 & -\mu_3 & 0 & 0 & 0 \end{bmatrix}, \qquad (13)$$

where $s = \sin\theta_W$, $c = \cos\theta_W$, $r_Z = M_Z/v$, and v_1, v_2, v_{ν_i} are the VEVs of h_1^0, h_2^0, and $\tilde{\nu}_i$ respectively, with $v_1^2 + v_2^2 + v_\nu^2 = v^2 \simeq (246\ \mathrm{GeV})^2$ and $v_\nu^2 = v_{\nu_1}^2 + v_{\nu_2}^2 + v_{\nu_3}^2$. We also define $\tan\beta = v_2/(v_1^2 + v_\nu^2)^{1/2}$.

To understand the structure of the above 7×7 mass matrix, let us assume that μ is the dominant term, then $\tilde{h}_{1,2}^0$ form a heavy Dirac particle of mass μ which mixes very little with the other physical fields. Removing these heavy fields will then give us the reduced 5×5 matrix in the basis $(\tilde{B}, \tilde{W}_3, \nu_1, \nu_2, \nu_3)$:

$$\mathcal{M} = \begin{bmatrix} M_1 - s^2\delta & sc\delta & -s\epsilon_1 & -s\epsilon_2 & -s\epsilon_3 \\ sc\delta & M_2 - c^2\delta & c\epsilon_1 & c\epsilon_2 & c\epsilon_3 \\ -s\epsilon_1 & c\epsilon_1 & 0 & 0 & 0 \\ -s\epsilon_2 & c\epsilon_2 & 0 & 0 & 0 \\ -s\epsilon_3 & c\epsilon_3 & 0 & 0 & 0 \end{bmatrix}, \qquad (14)$$

where

$$\delta = 2M_Z^2 \frac{v_1 v_2}{v^2} \frac{1}{\mu} = \frac{M_Z^2 \sin 2\beta}{\mu} \sqrt{1 - \frac{v_\nu^2}{v^2 \cos^2 \beta}}, \tag{15}$$

$$\epsilon_i = \frac{M_Z}{v} \Big(v_{\nu_i} - \frac{\mu_i}{\mu} v_1 \Big). \tag{16}$$

From the above, only the combination $\nu_l \equiv (\epsilon_1 \nu_1 + \epsilon_2 \nu_2 + \epsilon_3 \nu_3)/\epsilon$, with $\epsilon^2 = \epsilon_1^2 + \epsilon_2^2 + \epsilon_3^2$, mixes with the gauginos. This state will have an effective vanishing lepton number and the other two orthogonal combinations decouple from the neutralino mass matrix. In this case, only the eigenstate

$$\nu_l' = \nu_l + \frac{s\epsilon}{M_1} \tilde{B} - \frac{c\epsilon}{M_2} \tilde{W}_3, \tag{17}$$

gets a seesaw mass, i.e.

$$m_{\nu_l'} = -\epsilon^2 \left(\frac{s^2}{M_1} + \frac{c^2}{M_2} \right), \tag{18}$$

whereas the other two neutrinos remain massless. They may get masses through one-loop radiative corrections from the usual trilinear R-parity violating terms which we have not yet considered.

The two gauginos mix with the neutrino ν_l and form mass eigenstates given by

$$\tilde{B}' = \tilde{B} + \frac{sc\delta}{M_1 - M_2} \tilde{W}_3 - \frac{s\epsilon}{M_1} \nu_l, \tag{19}$$

$$\tilde{W}_3' = \tilde{W}_3 - \frac{sc\delta}{M_1 - M_2} \tilde{B} + \frac{c\epsilon}{M_2} \nu_l. \tag{20}$$

The physical states $\tilde{B}'$ and $\tilde{W}_3'$ now contain ν_l. This gives the main feature of R-parity violation, which is the decay of the lightest neutralino. By virtue of their ν_l components, both neutralinos will now decay into a lepton or an antilepton and a weak gauge boson, such as $\tilde{W}_3' \to l^- W^+$ and $l^+ W^-$, thus violating lepton number. Since the mixing of the neutralinos may also have CP violation through the complex gaugino masses (thus making δ complex), a lepton asymmetry may be generated from these decays. However, the amount of asymmetry thus generated is several orders of magnitude too small because it has to be much less than $(\epsilon/M_{1,2})^2$, which is of order $m_{\nu_l'}/M_{1,2}$, i.e. $< 5 \times 10^{-13}$ if $m_{\nu_l'} < 0.05$ eV and $M_{1,2} > 100$ GeV. In addition, the out-of-equilibrium condition on the decay width of the lightest neutralino imposes an upper bound on $(\epsilon/M_{1,2})^2$ which is independent of $m_{\nu_l'}$, and that also results in an asymmetry very much less than 10^{-10}.

We now consider the R-parity violating trilinear couplings, i.e. λ and λ' of Eq. (8). Since the particles involved should have masses at most equal to the supersymmetry breaking scale, i.e. a few TeV, their L violation together with the $B+L$ violation by sphalerons [16] would erase any primordial B or L asymmetry of the Universe [9]. To avoid such a possibility, we may reduce λ and λ' to less than about 10^{-7}, but a typical minimum value of 10^{-4} is required for realistic neutrino masses in one-loop order [11]. Hence it appears that the MSSM with R-parity violation is not only unsuitable for leptogenesis, it is also a destroyer of any lepton or baryon asymmetry which may have been created by some other means before the electroweak phase transition.

4 Leptogenesis from Neutralino Decay

From the discussion of the previous section we observe that for a leptogenesis mechanism to be successful in the MSSM with R-parity violation, two requirements have to be fulfilled. First we must use lepton-number violating terms which are not constrained by neutrino masses. Second we must satisfy the out-of-equilibrium condition for the lightest neutralino in such a way that the asymmetry is not automatically suppressed. More explicitly, we will consider the possibility that the heavier neutralino does not satisfy the out-of-equilibrium condition and decays very quickly, but the lighter neutralino decays very slowly and satisfies the out-of-equilibrium condition. Since the asymmetry comes from the interference of the one-loop CP violating contribution of the heavier neutralino, it is then unsuppressed. We will demonstrate explicitly in the following how this scenario may be realized.

We assume first that $M_1 > M_2$, so that the bino $\tilde{B}$ is heavier than the wino $\tilde{W}_3$. While the former couples to both $\bar{e}_{i_L}\tilde{e}_{i_L}$ and $\bar{e}^c_i\tilde{e}^c_i$, the latter couples only to $\bar{e}_{i_L}\tilde{e}_{i_L}$, because the e^c_i are singlets under $SU(2)_L$. Since R-parity is violated, one combination of the $\tilde{e}_{i_L}$ and another of the $\tilde{e}^c_i$ mix with the charged Higgs boson of the supersymmetric standard model: $h^\pm = h^\pm_2 \cos\beta + h^\pm_1 \sin\beta$. Let us denote them by $\tilde{l}_L$ and $\tilde{l}^c$ respectively. Their corresponding leptons are of course l_L and l^c. Hence both $\tilde{B}'$ and $\tilde{W}'_3$ may decay into $l^\mp h^\pm$.

We assume next that the $\tilde{l}_L$ mixing with h^- is negligible, so that the only relevant coupling is that of $\tilde{B}$ to $\bar{l}^c h^+$. Hence $\tilde{W}'_3$ decay (into $l^\mp h^\pm$) is suppressed because it may only do so through the small component of $\tilde{B}$ that it contains, assuming of course that all charged sleptons are heavier than $\tilde{B}$ or $\tilde{W}_3$.

With this choice that the heavier neutralino $\tilde{B}'$ decays quickly and the lighter neutralino $\tilde{W}'_3$ decays much more slowly, we now envisage the following leptogenesis scenario. At temperatures well above $T = M_{SUSY}$, there are fast lepton-number and R-parity violating interactions, which will wash out any L or B asymmetry of the Universe in the presence of sphalerons. This will be the case even at temperatures around M_1, when $\tilde{W}'_3$ interactions violate L_i as well as $B-3L_i$ for $i = e, \mu, \tau$ [9]. We assume here that all other supersymmetric particles are heavier than the neutralinos, so that at temperatures below M_1 we need only consider the interactions of $\tilde{B}'$ and $\tilde{W}_3'$. In Figure 1 we show the lepton-number violating processes (a) $\tilde{B}' \leftrightarrow l_R^\pm h^\mp$, where we have adopted the more conventional notation of an outgoing l_R in place of an incoming l^c. These processes are certainly still fast and there can be no L asymmetry. At temperatures far below the mass of the heavier neutralino, the $\tilde{B}'$ interactions are suppressed and we need only consider those of $\tilde{W}_3'$. With our assumptions, the lepton-number violating processes (b) $\tilde{W}_3' \leftrightarrow l_R^\pm h^\mp$ are slow and will satisfy the out-of-equilibrium condition for generating a lepton asymmetry of the Universe. Specifically, it comes 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. In this scenario the mass of $\tilde{W}_3'$ also has to be small enough so that the scattering processes mediated by the heavier $\tilde{B}'$ are negligible at temperature below M_2 when the asymmetry is produced.

We start with the well-known interaction of $\tilde{B}$ with l and $\tilde{l}_R$ given by [17]

$$-\frac{e\sqrt{2}}{\cos\theta_W}\left[\bar{l}\left(\frac{1-\gamma_5}{2}\right)\tilde{B}\tilde{l}_R + H.c.\right]. \tag{21}$$

We then allow $\tilde{l}_R$ 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. (20) with l and $h^\pm$ is given by

$$\left(\frac{sc\xi\delta r}{M_1-M_2}\right)\left(\frac{e\sqrt{2}}{\cos\theta_W}\right)\left[\bar{l}\left(\frac{1-\gamma_5}{2}\right)\tilde{W}_3' h^- + H.c.\right], \tag{22}$$

where ξ represents the $\tilde{l}_R - h^-$ mixing and is assumed real, but the parameter δ of Eq. (15) is complex. We have also inserted a correction factor $r = (1 + M_2/\mu \sin 2\beta)/(1 - M_2^2/\mu^2)$ for finite values of M_2/μ. The origin of a nontrivial CP phase in the above is from the 2×2 Majorana mass matrix spanning $\tilde{B}$ and $\tilde{W}_3$, with complex M_1 and M_2. It is independent of the phase of μ and contributes negligibly to the neutron electric dipole moment because the magnitude of δ is very small. [Note that the usual assumption of CP violation in supersymmetric models is that M_1 and M_2 have a common phase, in which case the phase of δ would be equal to the phase of μ.]

The decay width of the bino is then

$$\Gamma_{\tilde{B}'} = \Gamma(\tilde{B}' \to l^+h^-) + \Gamma(\tilde{B}' \to l^-h^+) = \frac{1}{4\pi}\xi^2\frac{e^2}{c^2}\frac{(M_{\tilde{B}'}^2 - m_h^2)^2}{M_{\tilde{B}'}^3}, \tag{23}$$

while that of the wino is

$$\Gamma_{\tilde{W}_3'} = \Gamma(\tilde{W}_3' \to l^+ h^-) + \Gamma(\tilde{W}_3' \to l^- h^+) = \frac{1}{4\pi}\xi^2 \left(\frac{es|\delta|r}{M_1 - M_2}\right)^2 \frac{(M_{\tilde{W}_3'}^2 - m_h^2)^2}{M_{\tilde{W}_3'}^3}. \qquad (24)$$

Using Eqs. (21) and (22), we calculate the interference between the tree-level and self-energy + vertex diagrams of Figure 1 and obtain the following asymmetry from the decay of $\tilde{W}_3'$:

$$\epsilon = \frac{\Gamma(\tilde{W}_3' \to l^+ h^-) - \Gamma(\tilde{W}_3' \to l^- h^+)}{\Gamma_{\tilde{W}_3'}} = \frac{\alpha \xi^2}{2\cos^2\theta_W} \frac{Im\delta^2}{|\delta|^2} \left(1 - \frac{m_h^2}{M_{\tilde{W}_3'}^2}\right)^2 \frac{x^{1/2} g(x)}{(1-x)}, \qquad (25)$$

where $x = M_{\tilde{W}_3'}^2 / M_{\tilde{B}'}^2$ and

$$g(x) = 1 + \frac{2(1-x)}{x} \left[\left(\frac{1+x}{x}\right) \ln(1+x) - 1\right]. \qquad (26)$$

If the $\tilde{W}_3'$ interactions satisfy the out-of-equilibrium condition, then a lepton asymmetry may be generated from the above decay asymmetry. Note that in the above expression for ϵ, the parameter δ appears only in the combination $\mathrm{Im}\delta^2/|\delta|^2$, which may be of order one. If the absolute value of δ is small, it slows down the decay rate of $\tilde{W}_3'$ and a departure from equilibrium may be achieved without affecting the amount of decay asymmetry generated in the process.

At the time this lepton asymmetry is generated, if the sphaleron interactions [16] are still in equilibrium, they will convert it into a baryon asymmetry of the Universe [3]. If the electroweak phase transition is strongly first-order, the sphaleron interactions freeze out at the critical temperature. Lattice simulations suggest that for a Higgs mass of around $m_H \approx 70$ GeV, the critical temperature is around $T_c \approx 150$ GeV [18]. Higher values of m_H will increase the critical temperature, but the increase is slower than linear. For example, for $m_H \approx 150$ GeV, the critical temperature could go up to $T_c \approx 250$ GeV.

For a second-order or weakly first-order phase transition[1], the sphaleron interactions freeze out at a temperature lower than the critical temperature. After the electroweak phase transition ($T < T_c$), the sphaleron transition rate is given by [20]

$$\Gamma_{sph}(T) = (2.2 \times 10^4\ \kappa)\ \frac{[2\ M_W(T)]^7}{[4\pi\ \alpha_W T]^3}\ e^{-E_{sph(T)}/T}, \qquad (27)$$

[1] Note that we do not require the electroweak phase transition to be first-order for satisfying the out-of-equilibrium condition. See for example Ref. [19].

where

$$M_W(T) = \frac{1}{2} g_2 \langle v(T) \rangle = \frac{1}{2} g_2 \langle v(T=0) \rangle \, \left(1 - \frac{T^2}{T_c^2}\right)^{1/2}, \tag{28}$$

and the free energy of the sphaleron is $E_{sph}(T) \approx (2 \; M_W(T)/\alpha(W)) \; B(m_h/M_W)$, with $B(0) = 1.52$, $B(\infty) = 2.72$ and $\kappa = e^{-3.6}$ [21]. In this case, the sphalerons freeze out at a temperature T_{out} which is the temperature at which their interaction strength equals the expansion rate of the universe,

$$\Gamma_{sph}(T_{out}) = H(T_{out}) = 1.7 \sqrt{g_*} \frac{T_{out}^2}{M_{Pl}}. \tag{29}$$

For a critical temperature of about $T_c \sim 250$ GeV, the freeze-out temperature comes out to be around $T_{out} \sim 200$ GeV.

These discussions indicate that as long as the lepton asymmetry is generated at a temperature above, say 200 GeV, it will be converted to a baryon asymmetry of the Universe. Since the sphaleron interactions grow exponentially fast, they can convert a lepton asymmetry to a baryon asymmetry [a $(B-L)$ asymmetry to be precise] by the time the temperature drops by only a few GeV. In the next section we discuss how the decay asymmetry of the neutralinos becomes a lepton asymmetry of the Universe.

5 Boltzmann Equations

We now solve the Boltzmann equations [24] to estimate the amount of lepton asymmetry created after the decays of the neutralinos. When the decay of the $\tilde{W}_3'$ satisfies the out-of-equilibrium condition, i.e. when the decay rate is slower than the expansion rate of the Universe, the generated asymmetry is of the order of the decay asymmetry given in Eq. (25). This argument could replace the details of solving the Boltzmann equations for an order-of-magnitude estimate of the asymmetry in many scenarios. However, in the present case there are other constraints and depleting factors, and we need to solve the Boltzmann equations explicitly for a reliable estimate.

If the $\tilde{W}_3'$ decay rate is much less than the expansion rate of the Universe, the generated lepton asymmetry is the same as the decay asymmetry. In other words, the out-of-equilibrium condition reads

$$K_{\tilde{W}_3'} = \frac{\Gamma_{\tilde{W}_3'}}{H(M_{\tilde{W}_3'})} \ll 1, \tag{30}$$

where $H(T)$ is the Hubble constant at the temperature T and is given by

$$H(T) = \sqrt{\frac{4\pi^3 g_*}{45}} \frac{T^2}{M_{Planck}}, \tag{31}$$

with g_* the number of massless degrees of freedom which we take equal to 106.75 and $M_{Pl} \sim 10^{18}$ GeV is the Planck scale. If this condition is satisfied, the lepton asymmetry is given by $n_L = n_l - n_{\bar{l}} \sim \epsilon/g_*$. But in practice, when $K \ll 1$, there is no time for the asymmetry to grow to its maximum value before the sphaleron transitions are over. So we need to study the case $K \sim 1$. Furthermore, a reasonable amount of asymmetry cannot be obtained unless the inverse decay and the scattering from bino exchange have rates lower than the expansion rate of the Universe. All these effects result in the further diminution of the lepton asymmetry and we need to solve the Boltzmann equations to take care of them properly.

At temperatures $T < M_2$, the decays of $\tilde{W}_3'$ given in Eq. (24) start generating an asymmetry. At this time there are important damping contributions coming from the inverse decays of $\tilde{W}_3'$ and $\tilde{B}'$ as well as the scattering processes $l^\pm + h^\mp \to \tilde{B}' \to l^\mp + h^\pm$. As we will see, the last two processes are especially important because $\tilde{B}'$ tends to remain in equilibrium and its presence washes out the created lepton asymmetry from the $\tilde{W}_3'$ decays. The reason is that its interactions are strong enough so that the Boltzmann exponential suppression of its number density may not be sufficient to compensate its large inverse decay and scattering cross sections. The effect of the scattering $l^\pm + h^\mp \to \tilde{W}_3' \to l^\mp + h^\pm$ is on the other hand negligible because it is suppressed by a factor of $[(sc\delta r)/(M_1 - M_2)]^2$ with respect to the scattering $l^\pm + h^\mp \to \tilde{B}' \to l^\mp + h^\pm$. Neglecting this term and defining the variable $z \equiv M_{\tilde{W}_3}/T$, the Boltzmann equations are then:

$$\frac{dX_{\tilde{W}_3'}}{dz} = -\gamma^{eq}_{\tilde{W}_3'} \frac{z}{sH(M_{\tilde{W}_3'})} \Big(\frac{X_{\tilde{W}_3'}}{X^{eq}_{\tilde{W}_3'}} - 1 \Big), \tag{32}$$

$$\begin{aligned} \frac{dX_L}{dz} &= \gamma^{eq}_{\tilde{W}_3'} \frac{z}{sH(M_{\tilde{W}_3'})} \Big[\varepsilon \Big(\frac{X_{\tilde{W}_3'}}{X^{eq}_{\tilde{W}_3'}} - 1 \Big) - \frac{1}{2} \frac{X_L}{X_\gamma} \Big] \\ &\quad - \frac{z}{sH(M_{\tilde{B}'})} \Big(\frac{M_{\tilde{B}'}}{M_{\tilde{W}_3'}} \Big)^2 \Big[\gamma^{eq}_{\tilde{B}'} \frac{1}{2} \frac{X_L}{X_\gamma} + 2 \frac{X_L}{X_\gamma} \gamma^{eq}_{scatt.} \Big], \end{aligned} \tag{33}$$

where we have defined the number densities per comoving volume $X_i = n_i/s$ in terms of the number densities of particles "i" and

$$s = g_* \frac{2\pi^2}{45} T^3 \tag{34}$$

is the entropy density. The equilibrium distributions of the number densities are given by the Maxwell-Boltzmann statistics:

$$n_{\tilde{W}_3'} = g_{\tilde{W}_3'} \frac{M^2_{\tilde{W}_3'}}{2\pi^2} T K_2(M_{\tilde{W}_3'}/T), \tag{35}$$

$$
\begin{aligned}
n_{\tilde{B}'} &= g_{\tilde{B}'}\frac{M_{\tilde{B}'}^2}{2\pi^2}TK_2(M_{\tilde{B}'}/T), && (36)\\
n_\gamma &= \frac{g_\gamma T^3}{\pi^2}, && (37)
\end{aligned}
$$

where $g_{\tilde{W}_3'} = 1$, $g_{\tilde{B}'} = 1$, and $g_\gamma = 2$ are the numbers of degrees of freedom of $\tilde{W}_3'$, $\tilde{B}'$, and the photon respectively.

The quantities $\gamma_{\tilde{W}_3'}^{eq}$ and $\gamma_{\tilde{B}'}^{eq}$ are the reaction densities for the decays and inverse decays of $\tilde{W}_3'$ and $\tilde{B}'$:

$$
\begin{aligned}
\gamma_{\tilde{W}_3'}^{eq} &= n_{\tilde{W}_3'}^{eq}\frac{K_1(M_{\tilde{W}_3'}/T)}{K_2(M_{\tilde{W}_3'}/T)}\Gamma_{\tilde{W}_3'}, && (38)\\
\gamma_{\tilde{B}'}^{eq} &= n_{\tilde{B}'}^{eq}\frac{K_1(M_{\tilde{B}'}/T)}{K_2(M_{\tilde{B}'}/T)}\Gamma_{\tilde{B}'}, && (39)
\end{aligned}
$$

K_1 and K_2 being the usual modified Bessel functions. The reaction density for the scattering is given by

$$
\gamma_{scatt.}^{eq} = \frac{T}{64\pi^4}\int_{(m_h+m_l)^2}^{\infty} ds\, \hat{\sigma}_{\tilde{B}'}(s)\,\sqrt{s}\,K_1(\sqrt{s}/T), \qquad (40)
$$

where $\hat{\sigma}_{\tilde{B}'}$ is the reduced cross section and is given by $2[s-(m_h+m_l)^2][s-(m_h-m_l)^2]\sigma_{\tilde{B}'}/s \sim 2s\sigma_{\tilde{B}'}$. The cross section $\sigma_{\tilde{B}'}$ does not contain the contribution of the on-mass-shell bino (which is already taken into account in the decay and inverse decay terms). This is achieved by replacing the usual propagator $1/(s-m^2+i\Gamma m)$ with the off-mass-shell propagator [22, 23]:

$$
D_s^{-1} = \frac{s-m^2}{(s-m^2)^2+\Gamma^2 m^2}. \qquad (41)
$$

The cross section $\sigma_{\tilde{B}'}$ contains the s- and t-channel contributions together with their interference terms and is given by

$$
\begin{aligned}
&\sigma_{\tilde{B}'} \equiv \sigma(l^\pm h^\mp \to \tilde{B}' \to l^\mp h^\pm) \;=\\
&\frac{1}{8\pi s^2}\Big(\frac{e^2\xi^2}{cos^2\theta_W}\Big)^2 m_{\tilde{B}'}^2\left[\frac{s^2}{D_s^2}+\frac{4s}{D_s}+\frac{2s}{M_{\tilde{B}'}^2}-\Big(2+4\frac{s+M_{\tilde{B}'}^2}{D_s}\Big)\ln\Big(1+\frac{s}{M_{\tilde{B}'}^2}\Big)\right] \qquad (42)
\end{aligned}
$$

In Eq. (40) the integral is dominated by the s-channel contribution in the resonance region and to a very good approximation, $\gamma_{scatt.}^{eq}$ reduces to

$$
\gamma_{scatt.}^{eq} = \frac{T}{512\pi^4}\frac{M_{\tilde{B}'}^4}{\Gamma_{\tilde{B}'}}K_1(M_{\tilde{B}'}/T)\frac{e^4\xi^4}{cos^4\theta_W}. \qquad (43)
$$

From Eqs. (34) to (43), we find the Boltzmann equations, i.e. (32) and (33), to be given by

$$
\begin{aligned}
\frac{dX_{\tilde{W}_3'}}{dz} &= -zK_{\tilde{W}_3'}\frac{K_1(z)}{K_2(z)}\Big(X_{\tilde{W}_3'}-X^{eq}_{\tilde{W}_3'}\Big), \\
\frac{dX_L}{dz} &= zK_{\tilde{W}_3'}\frac{K_1(z)}{K_2(z)}\Big[\varepsilon(X_{\tilde{W}_3'}-X^{eq}_{\tilde{W}_3'})-\frac{1}{2}\frac{X_{\tilde{W}_3}}{X_\gamma}X_L\Big] \\
&\quad -z\Big(\frac{M_{\tilde{B}'}}{M_{\tilde{W}_3'}}\Big)^2 K_{\tilde{B}'}\Big[\frac{1}{2}\frac{K_1(zM_{\tilde{B}'}/M_{\tilde{W}_3'})}{K_2(zM_{\tilde{B}'}/M_{\tilde{W}_3'})}\frac{X_{\tilde{B}'}}{X_\gamma}X_L+2\frac{X_L}{X_\gamma}\frac{\gamma^{eq}_{scatt.}}{s\Gamma_{\tilde{B}'}}\Big],
\end{aligned}
\tag{44}
$$

with

$$
K_{\tilde{B}'}=\Gamma_{\tilde{B}'}/H(M_{\tilde{B}'}), \tag{45}
$$

which gives the strength of lepton-number violation in the decays of $\tilde{B}'$.

If we now ignore the inverse decay and scattering processes, we can simplify the problem by requiring the out-of-equilibrium condition to be

$$
K_{\tilde{W}_3'}<1. \tag{46}
$$

With the terms proportional to $K_{\tilde{B}'}$ this condition is necessary but not sufficient. In the present scenario for a large asymmetry we also require $K_{\tilde{B}'}>1$. Indeed, $K_{\tilde{B}'}$ is larger than $K_{\tilde{W}_3'}$ by a factor $R_K=K_{\tilde{B}'}/K_{\tilde{W}_3'}\sim[(sc\delta r)/(M_1-M_2)]^{-2}(M_{\tilde{W}_3'}/M_{\tilde{B}'})$, which is larger than one by several orders of magnitude. Therefore $\tilde{B}'$ remains in equilibrium and the $\tilde{B}'$ damping terms, due to its inverse decay and scattering, dominate over the $\tilde{W}_3'$ inverse decay damping term as long as the Boltzmann suppression factor in the $\tilde{B}'$ equilibrium distribution has not compensated the large value of R_K. For example with the set of parameters $M_{\tilde{W}_3'}=2$ TeV, $M_{\tilde{B}'}=3$ TeV, $\sin 2\beta=0.5$, $\xi=2\times10^{-3}$, $\mu=5$ TeV used in Ref. [12], we obtain $K_{\tilde{W}_3'}=0.63$ and $K_{\tilde{B}'}=7.8\times10^5$, and the $\tilde{B}'$ damping terms dominate over those of $\tilde{W}_3'$ as long as the temperature is above $\sim$ 65 GeV (the former differs from the latter by a factor $\sim R_K(M_{\tilde{B}'}/M_{\tilde{W}_3'})^3 e^{-(M_{\tilde{B}'}-M_{\tilde{W}_3'})/T}$). In this case the inverse decay and scattering of $\tilde{B}'$ cause a considerable wash-out of the asymmetry because $\tilde{W}_3'$ has mostly decayed away already at temperatures well above $\sim$ 65 GeV. This is illustrated in Figure 2 showing the effects of various terms in the Boltzmann equations.

To avoid this wash-out, the value of $M_{\tilde{B}'}$ has to be larger in order that the $\tilde{B}'$ number density is further suppressed at temperatures below $M_{\tilde{W}_3'}$ when the asymmetry is produced. Varying the parameters of these $\tilde{B}'$ damping terms, it appears difficult to induce a sufficiently large asymmetry of order 10^{-10} for $M_{\tilde{B}'}$ below 4 TeV. Two typical situations for which a large asymmetry is produced are for example:

$$
M_{\tilde{B}'}=6\text{ TeV},\quad M_{\tilde{W}_3'}=3.5\text{ TeV},\quad \xi=5\times10^{-3},
$$

$$\mu = 10 \text{ TeV}, \quad \sin 2\beta = 0.10, \quad m_h = 200 \text{ GeV}, \tag{47}$$

and

$$\begin{aligned} M_{\tilde{B}'} &= 5 \text{ TeV}, \quad M_{\tilde{W}'_3} = 2 \text{ TeV}, \quad \xi = 5 \times 10^{-3}, \\ \mu &= 7.5 \text{ TeV}, \quad \sin 2\beta = 0.05, \quad m_h = 200 \text{ GeV}, \end{aligned} \tag{48}$$

for which we have $K_{\tilde{W}'_3} = 0.02$, $K'_{\tilde{B}'} = 2.4 \times 10^6$ and $K_{\tilde{W}'_3} = 0.02$, $K'_{\tilde{B}'} = 2.9 \times 10^6$ respectively. At $T = M_Z$ the leptonic asymmetry produced is $X_L = 1.0 \times 10^{-10}$ with the parameters of Eqs. (47) and $X_L = 1.2 \times 10^{-10}$ with those of Eq. (48). Figures 3 and 4 show the evolution of the asymmetry in these two cases. As can be seen from these figures, the damping effects of the inverse decay of $\tilde{W}'_3$ and of the scattering are small[2]. The damping effect from the inverse decay of $\tilde{B}'$ is however not small and reduces the asymmetry by a factor of 2 to 4 by washing out all the asymmetry produced above $T \sim 300 - 400$ GeV.

A large asymmetry of order 10^{-10} is produced provided $M_{\tilde{B}'}$ is of the order 4 TeV or more. A low value of $\sin 2\beta$ below ~ 0.30 is generally necessary. Values of ξ around $3 - 5 \times 10^{-3}$, of μ around $5 - 10$ TeV, and of $M_{\tilde{W}'_3}$ from 1 TeV to $2M_{\tilde{B}'}/3$ are also preferred.

6 Charged Scalar Mass Matrix

The mechanism we propose for leptogenesis requires the decay of $\tilde{B}'$ to be fast, while that of $\tilde{W}'_3$ is very slow. This is achieved by requiring $l^\pm_R h^\mp$ to be the main decay mode and $l^\pm_L h^\mp$ to be negligible. Hence $\tilde{l}_R$ must mix with h^- readily, so that $\tilde{B}'$ could decay directly, but $\tilde{W}'_3$ could decay only through its small $\tilde{B}$ component. We now consider the charged scalar mass matrix which determines this mixing. As shown in the following, our present scenario requires one more new ingredient, i.e. the presence of nonholomorphic soft terms.

The value of the $\tau_R - h^+$ mixing parameter ξ is governed by the charged scalar mass matrix which follows from the quadratic terms in the Lagrangian:

$$\mathcal{L} \ni -\Phi^\dagger \mathcal{M}^2_{S^\pm} \Phi, \tag{49}$$

with $\Phi = [h_1^{-*}, h_2^+, \tilde{e}_{i_L}^{-*}, \tilde{e}_i^c]^T$. In the case where all N constants are put to zero and for parameters satisfying the various constraints from the Boltzmann equations (see previous section), it is difficult to generate a sufficiently large value of the $\tilde{l}^c - h^-$ mixing parameter ξ (see Appendix A). Thus we will ignore the μ_i terms and the associated vaccuum

[2] Note that in the Boltzmann equations we neglected the damping contributions of the scatterings $l^\pm l^\pm \to h^\pm h^\pm$ mediated by a neutralino in the t-channel. Their effect is negligible within the ranges of the parameters we consider.

expectation values v_{ν_i} in the following. To induce a large $\tilde{e}^c - h^-$ mixing, we introduce the nonholomorphic terms of Eqs. (11) and (12). The mass matrix is then given by

$$\mathcal{M}^2_{S^\pm} = \begin{pmatrix} \mathcal{M}^2_{h_1^- - h_1^+} & \mathcal{M}^2_{h_1^- - h_2^+} & \mathcal{M}^2_{h_1^- - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{h_1^- - \tilde{e}^c_i} \\ \mathcal{M}^2_{h_2^- - h_1^+} & \mathcal{M}^2_{h_2^- - h_2^+} & \mathcal{M}^2_{h_2^- - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{h_2^- - \tilde{e}^c_i} \\ \mathcal{M}^2_{\tilde{e}_{j_L} - h_1^+} & \mathcal{M}^2_{\tilde{e}_{j_L} - h_2^+} & \mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^c_i} \\ \mathcal{M}^2_{\tilde{e}^c_j - h_1^+} & \mathcal{M}^2_{\tilde{e}^c_j - h_2^+} & \mathcal{M}^2_{\tilde{e}^c_j - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{\tilde{e}^c_j - \tilde{e}^c_i} \end{pmatrix}, \tag{50}$$

where

$$\mathcal{M}^2_{h_1^- - h_1^+} = \frac{g^2}{4} v_2^2 - B\frac{v_2}{v_1}, \tag{51}$$

$$\mathcal{M}^2_{h_1^- - h_2^+} = \frac{g^2}{4} v_1 v_2 - B, \tag{52}$$

$$\mathcal{M}^2_{h_1^- - \tilde{e}^*_{i_L}} = N_i'^B, \tag{53}$$

$$\mathcal{M}^2_{h_1^- - \tilde{e}^c_i} = \frac{1}{\sqrt{2}} N_i'^e v_2, \tag{54}$$

$$\mathcal{M}^2_{h_2^- - h_2^+} = \frac{g^2}{4} v_1^2 - B\frac{v_1}{v_2}, \tag{55}$$

$$\mathcal{M}^2_{h_2^- - \tilde{e}^*_{i_L}} = -B_i', \tag{56}$$

$$\mathcal{M}^2_{h_2^- - \tilde{e}^c_i} = \frac{1}{\sqrt{2}} N_i'^e v_1, \tag{57}$$

$$\mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^*_{i_L}} = (M_L^2)_{ij} - \frac{1}{8}(g^2 - g'^2)(v_1^2 - v_2^2)\delta_{ij} + \frac{1}{2} f^e_{jk} f^e_{ik} v_1^2, \tag{58}$$

$$\mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^c_i} = \frac{1}{\sqrt{2}} f^e_{ji}\mu v_2 + \frac{1}{\sqrt{2}}(A_e f_e)_{ji} v_1 + \frac{1}{\sqrt{2}} N^e_{ji} v_2, \tag{59}$$

$$\mathcal{M}^2_{\tilde{e}^{c*}_j - \tilde{e}^c_i} = (M^2_{\tilde{e}^c})_{ji} - \frac{g'^2}{4}(v_1^2 - v_2^2)\delta_{ij} + \frac{1}{2} f^e_{ki} f^e_{kj} v_1^2. \tag{60}$$

In the following we assume that $N_i'^B$ and B_i' are negligible so that they do not induce a large mixing of the left-handed charged sleptons with the charged Higgs boson. We are then left with the $N_i'^e$ and N^e_{ji} terms. Going to the basis of the physical charged Higgs boson h^+ and of the Goldstone boson G^+, the latter decouples and in the basis $[h^+, \tilde{e}^{-*}_{i_L}, \tilde{e}^c_i]$ we obtain the mass matrix

$$\mathcal{M}^2_{S^\pm} = \begin{pmatrix} m_W^2 - 2\frac{B}{\sin 2\beta} & 0 & \frac{1}{\sqrt{2}} N_i'^e v \\ 0 & \mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^c_i} \\ \frac{1}{\sqrt{2}} N_i'^e v & \mathcal{M}^2_{\tilde{e}^c_j - \tilde{e}^*_{i_L}} & \mathcal{M}^2_{\tilde{e}^c_j - \tilde{e}^c_i} \end{pmatrix}. \tag{61}$$

We observe that only one combination of the charged right-handed sleptons mixes with the charged Higgs boson:

$$\tilde{l}^c = \frac{N_1'^e \tilde{e}_1^c + N_2'^e \tilde{e}_2^c + N_3'^e \tilde{e}_3^c}{N'^e}, \tag{62}$$

with $(N'^e)^2 = (N_1'^e)^2 + (N_2'^e)^2 + (N_3'^e)^2$. In the Lagrangian the mass term which couples the charged Higgs boson with the sleptons reduces therefore to the following single term:

$$\mathcal{L} \ni -\frac{1}{\sqrt{2}} v N'^e h^+ \tilde{l}^c \, + \, h.c. \tag{63}$$

In the case of one family, we obtain to a very good approximation:

$$\xi \;=\; \frac{\frac{1}{\sqrt{2}} N'^e v}{m_{h^+}^2 - M_{\tilde{e}_L - \tilde{e}_L^*}^2 - \frac{\left(M_{\tilde{e}_L - \tilde{e}^c}^2\right)^2}{m_{h^+}^2 - M_{\tilde{e}_L - \tilde{e}_L^*}^2}}, \tag{64}$$

$$\xi' \;=\; \frac{M_{\tilde{e}_L - \tilde{e}^c}^2}{m_{h^+}^2 - M_{\tilde{e}_L - \tilde{e}_L^*}^2} \, \xi, \tag{65}$$

where ξ' is the $h^+ - \tilde{e}_L$ mixing and

$$m_{h^+}^2 = m_W^2 - 2 \frac{B}{\sin 2\beta}. \tag{66}$$

As required in section 3, the mixing ξ' has to be much smaller than the mixing ξ in order to avoid having the transition $\tilde{W}_3' \to \tilde{e}_L^{\pm} e_L^{\mp} \to h^{\pm} e_L^{\mp}$. This requires

$$|\xi'| \; < \left| \xi \frac{2 s^2 \delta r}{M_1 - M_2} \right| \tag{67}$$

which for the parameters of Eq. (47) implies that $\xi' < 7 \times 10^{-5} \, \xi$. Our mechanism requires also that $m_{h^+} < M_{\tilde{W}_3'}$ and that the mass of any charged slepton is larger than $M_{\tilde{W}_3'}$. We then have

$$\xi \simeq \frac{\frac{1}{\sqrt{2}} N'^e v}{m_{h^+}^2 - M_{\tilde{e}_L - \tilde{e}_L^*}^2}, \tag{68}$$

$$\left| m_e \mu \tan\beta + A_e m_e + \frac{1}{\sqrt{2}} N_e v_2 \right| < M_{\tilde{e}_L - \tilde{e}_L^*}^2 \left| \frac{2 s^2 \delta r}{M_1 - M_2} \right|, \tag{69}$$

which for the set of parameters given in Eq. (47) and for $M_{\tilde{e}_L - \tilde{e}_L^*}^2 \sim M_{\tilde{B}'}^2$ gives $N'^e \sim 1$ TeV and $|m_e \mu \tan\beta + A_e m_e + \frac{1}{\sqrt{2}} N_e v_2| < (50\,\mathrm{GeV})^2$. The latter condition requires $m_e < 0.02$ GeV, $|A_e m_e| < (50\,\mathrm{GeV})^2$, and $|N_e| < 15$ GeV, or a cancellation of the three terms together. In the case where there is no cancellation between these terms the charged slepton which mixs with the charged Higgs boson must have predominantly an electron or μ flavor. In the case of a cancellation between these terms, all flavors are possible. For other sets of parameters which lead to large asymmetries, these numerical bounds can be relaxed easily by a factor of 2 to 4. In the more general case of three families, similar constraints and relations are obtained.

7 Two-Loop Neutrino Mass

It is interesting to note that in addition to inducing a lepton asymmetry, the nonholomorphic terms $N_i^{\prime e}$ could also generate a neutrino mass. Since lepton number is violated at most by one unit in each term, the neutrino mass should include at least two lepton-violating vertices in a loop diagram.

There exists a one-loop diagram (Fig. 5), which contributes to the sneutrino "Majorana" mass. In general, the sneutrinos could have diagonal lepton number conserving masses, i.e. $\tilde{\nu}^*\tilde{\nu}$. But there can be also lepton-number violating mass terms, i.e. $\tilde{\nu}\tilde{\nu}$ [25, 26, 27]. In the present model, the nonholomorphic terms give rise to a lepton-number violating sneutrino-antisneutrino mixing term, i.e. $\mathcal{L} \ni -\frac{1}{2}\delta m^2_{\tilde{\nu}_{ij}}\tilde{\nu}_i\tilde{\nu}_j + h.c.$. In the case of one family, we get

$$\delta m^2_{\tilde{\nu}} \sim \frac{1}{8\pi^2}\frac{\mu^2\xi^2}{v^2}m_l^2. \tag{70}$$

This lepton-number violating sneutrino mass can then induce a Majorana neutrino mass [26].

In the present case the one-loop diagram of Fig. 6 gives a neutrino mass

$$\begin{aligned} m_\nu &\sim \frac{1}{32\pi^2}\frac{e^2}{\sin^2\theta_W}\delta m^2_{\tilde{\nu}} M_{\tilde{W}'_3}\frac{M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3} - M^2_{\tilde{W}'_3}\ln\left(M^2_{\tilde{\nu}}/M^2_{\tilde{W}'_3}\right)}{\left(M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3}\right)^2} \\ &\sim \frac{1}{256\pi^4}\frac{e^2}{\sin^2\theta_W}\mu^2\frac{m_l^2}{v^2}\xi^2 M_{\tilde{W}'_3}\frac{M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3} - M^2_{\tilde{W}'_3}\ln\left(M^2_{\tilde{\nu}}/M^2_{\tilde{W}'_3}\right)}{\left(M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3}\right)^2}. \end{aligned} \tag{71}$$

In the case where the lepton l which mixes with h^+ is essentially τ, we get

$$m_\nu = \frac{1}{256\pi^4}\frac{e^2}{\sin^2\theta_W}\mu^2\frac{m_\tau^2}{v^2}\xi^2 M_{\tilde{W}'_3}\frac{M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3} - M^2_{\tilde{W}'_3}\ln\left(M^2_{\tilde{\nu}}/M^2_{\tilde{W}'_3}\right)}{\left(M^2_{\tilde{\nu}} - M^2_{\tilde{W}'_3}\right)^2}, \tag{72}$$

which has the correct order of magnitude. For example with the parameters of Eq. (47) and taking $M_{\tilde{\nu}} \sim M_{\tilde{B}'}$ we get $m_{\nu_\tau} \sim 0.1$ eV. The value of ξ we need for having the right order of magnitude for the asymmetry is therefore also the one we need to have a neutrino mass, in agreement with the present data on atmospheric neutrinos. In Eq. (71), the factor $\delta m^2_{\tilde{\nu}}$ appears because there is GIM (Glashow-Iliopoulos-Maiani) suppression from summing over all possible neutral slepton eigenstates in the loop. In Fig. 5 the two-point functions of the form $f(m^2, m'^2, p) \equiv (i/\pi^2) \int d^4k (k^2 - m^2)^{-1}((k+p)^2 - m'^2)^{-1}$ have been (roughly) approximated by ~ 1 while in Fig. 6 the two-point functions have been calculated explicitly (as required by the fact that for these diagrams, a GIM suppression mechanism is operative). This can also be understood from another point of view. Since the diagonal terms of the sneutrino mass come from the lepton-number conserving interactions, they should not contribute to the Majorana mass of a neutrino. Only the lepton-number violating sneutrino mass, which is the mass-squared difference, should contribute to the Majorana neutrino mass. This makes the neutrino mass proportional to the mass-squared difference after GIM cancellation. Note that combining both one-loop diagrams, a two-loop diagram is obtained which is similar to the diagram proposed in Ref. [28] with a different lepton-number violating soft term.

In the case of three families, the induced neutrino mass terms involving ν_e are suppressed by the small value of the electron mass with respect to the μ or τ mass. Therefore, unless N'^e_e is much larger than N'^e_τ and N'^e_μ, ν_e essentially decouples and acquires a very small mass; we get $m_{\nu_e} \sim 10^{-8}$ eV or less. In this case, the mass matrix of the sneutrinos in the $\mu - \tau$ sector is of the form:

$$\mathcal{L} \ni -\frac{1}{2} \Phi^\dagger_{\tilde{\nu}} \mathcal{M}^2_{\tilde{\nu}} \Phi_{\tilde{\nu}}, \tag{73}$$

with $\Phi_{\tilde{\nu}} = (\tilde{\nu}_\mu, \tilde{\nu}_\tau, \tilde{\nu}^\dagger_\mu, \tilde{\nu}^\dagger_\tau)^T$ and

$$\mathcal{M}^2_{\tilde{\nu}} = \begin{pmatrix} M^2_{L_\mu} & 0 & \delta m^2_{\tilde{\nu}_{\mu\mu}} & \delta m^2_{\tilde{\nu}_{\mu\tau}} \\ 0 & M^2_{L_\tau} & \delta m^2_{\tilde{\nu}_{\mu\tau}} & \delta m^2_{\tilde{\nu}_{\tau\tau}} \\ \delta m^2_{\tilde{\nu}_{\mu\mu}} & \delta m^2_{\tilde{\nu}_{\mu\tau}} & M^2_{L_\mu} & 0 \\ \delta m^2_{\tilde{\nu}_{\mu\tau}} & \delta m^2_{\tilde{\nu}_{\tau\tau}} & 0 & M^2_{L_\tau} \end{pmatrix}, \tag{74}$$

where for simplicity we have assumed in Eq. (9) a diagonal matrix $M^2_L = diag(M^2_{L_\mu}, M^2_{L_\tau})$ and with

$$\delta m^2_{\tilde{\nu}_{ij}} \sim \frac{1}{8\pi^2} \frac{\mu^2 \xi^2}{v^2} m_{l_i} m_{l_j} \frac{N'^e_i N'^e_j}{(N'^e)^2}. \tag{75}$$

From the mass matrix of Eq. (74), the diagrams of Fig. 6 induce then the following neutrino mass term

$$\mathcal{L} \ni -\frac{1}{2} \Psi^\dagger_\nu \mathcal{M}_\nu \Psi_\nu, \tag{76}$$

with $\Psi_\nu = (\nu_\mu, \nu_\mu^\dagger, \nu_\tau, \nu_\tau^\dagger)^T$ and

$$\mathcal{M}_\nu \sim A \left(\begin{array}{cccc} 0 & \Delta_{\mu-\mu} & 0 & \Delta_{\mu-\tau} \\ \Delta_{\mu-\mu} & 0 & \Delta_{\mu-\tau} & 0 \\ 0 & \Delta_{\mu-\tau} & 0 & \Delta_{\tau-\tau} \\ \Delta_{\mu-\tau} & 0 & \Delta_{\tau-\tau} & 0 \end{array} \right) , \tag{77}$$

where

$$\begin{aligned} \Delta_{i-i} &= M_{\tilde{W}_3'} \Big(f(M_{L_i}^2 + \delta m^2_{\tilde{\nu}_{ii}}, M^2_{\tilde{W}_3'}, 0) - f(M_{L_i}^2 - \delta m^2_{\tilde{\nu}_{ii}}, M^2_{\tilde{W}_3'}, 0) \Big) \\ &\sim 2\delta m^2_{\tilde{\nu}_{ii}} M_{\tilde{W}_3'} \frac{M_{L_i}^2 - M^2_{\tilde{W}_3'} - M^2_{\tilde{W}_3'} \ln\left(M^2_{L_i}/M^2_{\tilde{W}_3'} \right)}{\left(M^2_{L_i} - M^2_{\tilde{W}_3'} \right)^2}, \end{aligned} \tag{78}$$

$$\begin{aligned} \Delta_{\mu-\tau} &= \frac{2\delta m^2_{\tilde{\nu}_{\mu\tau}}}{M^2_{L_\tau} - M^2_{L_\mu}} M_{\tilde{W}_3'} \Big(f(M^2_{L_\tau}, M^2_{\tilde{W}_3'}, 0) - f(M^2_{L_\mu}, M^2_{\tilde{W}_3'}, 0) \Big) \\ &\sim 2\delta m^2_{\tilde{\nu}_{\mu\tau}} M_{\tilde{W}_3'} (M^2_{L_\tau} - M^2_{L_\mu})^{-1} (M^2_{L_\tau} - M^2_{\tilde{W}_3'})^{-1} (M^2_{L_\mu} - M^2_{\tilde{W}_3'})^{-1} \\ &\times \left[M^2_{L_\mu} M^2_{\tilde{W}_3'} \ln\left(\frac{M^2_{L_\mu}}{M^2_{\tilde{W}_3'}} \right) + M^2_{L_\tau} M^2_{\tilde{W}_3'} \ln\left(\frac{M^2_{\tilde{W}_3'}}{M^2_{L_\tau}} \right) + M^2_{L_\tau} M^2_{L_\mu} \ln\left(\frac{M^2_{L_\tau}}{M^2_{L_\mu}} \right) \right], \end{aligned} \tag{79}$$

and

$$A = \frac{1}{64\pi^2} \frac{e^2}{\sin^2\theta_W}. \tag{80}$$

This matrix can lead easily to a maximal mixing between the μ and τ neutrinos. This will be the case in particular if $\Delta_{\mu-\mu} \sim \Delta_{\tau-\tau}$ which implies

$$\frac{m_\tau N_\tau'^e}{M^2_{L_\tau}} \sim \frac{m_\mu N_\mu'^e}{M^2_{L_\mu}}. \tag{81}$$

In addition it can be seen easily that in the limit where $M^2_{L_\mu} = M^2_{L_\tau}$, the determinant of the neutrino mass matrix vanishes, leading to a large hierarchy of masses (as required by atmospheric and solar neutrino experiments, taking into account the fact that the mass of ν_e is below 10^{-8} eV in the present scenario). For example with the parameters of Eq. (47) and taking in addition $N_\mu'^e = 14 N_\tau'^e$, $M_{L_\mu} \sim M_{\tilde{B}'} = 6$ TeV, and $M_{L_\tau} \sim 7.5$ TeV, we obtain one neutrino with a mass $\sim 10^{-3}$ eV and one with a mass $\sim 10^{-5}$ eV in addition to the electron neutrino with a mass below 10^{-8} eV. In this case the mixing between the μ and τ flavors is large ($\sin 2\alpha = 0.99$) while that of the electron flavor with the two other flavors is very much suppressed. Note that the values of the lepton-number violating mass terms $\delta m^2_{\tilde{\nu}_{ij}}$ induced by Fig. 5 are several orders of magnitude below the phenomenological bounds

$\delta m_e < 350$ MeV, $\delta m_\mu < 50$ GeV, and $\delta m_\tau < 450$ GeV obtained for $M_{SUSY} \sim 1$ TeV in Ref. [27].

In summary, from the above qualitative estimate, we observe that realistic neutrino masses could be accomodated easily in the present scenario, in agreement with atmospheric and solar neutrino experiments. A large mixing and a hierarchy of neutrino masses appear rather naturally. A more quantitative estimate would require an explicit calculation of the two-loop integrals involved, but since there are still many free parameters, it will not add much to our understanding in any case.

8 Conclusion

We have studied a model of leptogenesis in a R-parity violating supersymmetric model. The lightest neutralino $\tilde{W}_3'$ is assumed to be mostly the $SU(2)$ gaugino but its decay into $l^\pm h^\mp$ is suppressed because the required $\tilde{l}_L - h^+$ mixing is negligible. On the other hand, $\tilde{W}_3'$ has a small component of $\tilde{B}$, the $U(1)$ gaugino, which decays readily because the required $\tilde{l}_R - h^+$ mixing is of order 10^{-3} from the presence of nonholomorphic R-parity violating soft terms in the Lagrangian. The decay asymmetry of $\tilde{W}_3'$ is then evolved into a lepton asymmetry of the Universe by solving the Boltzmann equations in detail numerically. We demonstrate how each term in the equations affects the eventual outcome of the proposed scenario. The charged scalar mass matrix and the neutralino sector are discussed in detail. A realistic scenario of radiative neutrino mass generation in two loops is presented, which originates from the same lepton-number violating nonholomorphic terms.

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A Complete Charged Scalar Mass Matrix

From Eqs. (7) to (12), neglecting small terms of order $(\lambda_{ijk})^2$ and $\lambda_{ijk}\lambda'_{lmn}$, the complete charged scalar mass matrix is given by Eq. (50) with (see also Refs. [15, 29] for the holomorphic part):

$$\begin{aligned}
\mathcal{M}^2_{h_1^- - h_1^+} &= \frac{g^2}{4}(v_2^2 - v_\nu^2) - B\frac{v_2}{v_1} - \mu\mu_i\frac{v_{\nu_i}}{v_1} + \frac{1}{2}f^e_{ij}f^e_{kj}v_{\nu_i}v_{\nu_k} - N'^B_i\frac{v_{\nu_i}}{v_1}, &(82)\\
\mathcal{M}^2_{h_1^- - h_2^+} &= \frac{g^2}{4}v_1v_2 - B, &(83)\\
\mathcal{M}^2_{h_1^- - \tilde{e}^*_{i_L}} &= \frac{g^2}{4}v_1v_{\nu_i} + \mu\mu_i - \frac{f^e_{kj}f^e_{ij}}{2}v_1v_{\nu_k} + \frac{1}{2}f^e_{lj}(\lambda_{ikj} - \lambda_{kij})v_{\nu_l}v_{\nu_k} + N'^B_i, &(84)\\
\mathcal{M}^2_{h_1^- - \tilde{e}^c_i} &= -\frac{1}{\sqrt{2}}f^e_{ji}\mu_j v_2 - \frac{1}{\sqrt{2}}(A_e f_e)_{ji}v_{\nu_j} + \frac{1}{\sqrt{2}}N'^e_i v_2, &(85)\\
\mathcal{M}^2_{h_2^- - h_2^+} &= \frac{g^2}{4}(v_1^2 + v_\nu^2) - B\frac{v_1}{v_2} - B'_i\frac{v_{\nu_i}}{v_2}, &(86)\\
\mathcal{M}^2_{h_2^- - \tilde{e}^*_{i_L}} &= \frac{g^2}{4}v_2v_{\nu_i} - B'_i, &(87)\\
\mathcal{M}^2_{h_2^- - \tilde{e}^c_i} &= +\frac{1}{\sqrt{2}}f^e_{ji}(\mu v_{\nu_j} - \mu_j v_1) + \frac{1}{\sqrt{2}}\lambda_{kji}(\mu_k v_{\nu_j} - \mu_j v_{\nu_k})\\
&\quad +\frac{1}{\sqrt{2}}N^e_{ji}v_{\nu_j} + \frac{1}{\sqrt{2}}N'^e_i v_1, &(88)\\
\mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^*_{i_L}} &= (M_L^2)_{ij} - \frac{1}{8}(g^2 - g'^2)(v_1^2 - v_2^2 + v_\nu^2)\delta_{ij} + \frac{g^2}{4}v_{\nu_i}v_{\nu_j} + \frac{1}{2}f^e_{jk}f^e_{ik}v_1^2 + \mu_j\mu_i\\
&\quad +\frac{1}{2}f^e_{jl}v_1v_{\nu_k}(\lambda_{kil} - \lambda_{ikl}) + \frac{1}{2}f^e_{il}v_1v_{\nu_k}(\lambda_{kjl} - \lambda_{jkl}), &(89)\\
\mathcal{M}^2_{\tilde{e}_{j_L} - \tilde{e}^c_i} &= \frac{1}{\sqrt{2}}f^e_{ji}\mu v_2 + \frac{1}{\sqrt{2}}(A_e f_e)_{ji}v_1 + \frac{1}{\sqrt{2}}(A'^e_{kji} - A'^e_{jki})v_{\nu_k}\\
&\quad +\frac{1}{\sqrt{2}}\mu_k(\lambda_{kji} - \lambda_{jki})v_2 + \frac{1}{\sqrt{2}}N^e_{ji}v_2, &(90)\\
\mathcal{M}^2_{\tilde{e}^{c*}_j - \tilde{e}^c_i} &= (M^2_{\tilde{e}^c})_{ji} - \frac{g'^2}{4}(v_1^2 - v_2^2 + v_\nu^2)\delta_{ij} + \frac{1}{2}f^e_{ki}f^e_{lj}v_{\nu_k}v_{\nu_l} + \frac{1}{2}f^e_{ki}f^e_{kj}v_1^2\\
&\quad +\frac{1}{2}f^e_{kj}(\lambda_{lki} - \lambda_{kli})v_1v_{\nu_l} + \frac{1}{2}f^e_{ki}(\lambda_{lkj} - \lambda_{klj})v_1v_{\nu_l}. &(91)
\end{aligned}$$

In Eqs. (82) and (86), the tadpoles conditions for the h_1^0 and h_2^0 fields have been used:

$$\Big(m^2_{H_1} + \mu^2 + \frac{1}{8}(g^2 + g'^2)(v_1^2 - v_2^2 + v_\nu^2)\Big)v_1 + (\mu\mu_i + N'^B_i)v_{\nu_i} = 0, \qquad (92)$$

$$\Big(m^2_{H_2} + \mu^2 + \mu_i^2 - \frac{1}{8}(g^2 + g'^2)(v_1^2 - v_2^2 + v_\nu^2)\Big)v_2 + Bv_1 + B'_i v_{\nu_i} = 0. \qquad (93)$$

In Eqs. (82) to (91), v_{ν_i} are given by the corresponding tadpole conditions for the $\tilde{\nu}_i$ fields:

$$\Big((M_L^2)_{ji} + \mu_i\mu_j + \frac{1}{8}(g^2 + g'^2)(v_1^2 - v_2^2 + v_\nu^2)\delta_{ij}\Big)v_{\nu_j} + B'_i v_2 + N_i'^B v_1 + \mu\mu_i v_1 \;=\; 0. \quad (94)$$

Without nonholomorphic terms, it is difficult to obtain a large $\tilde{e}_R - h^+$ mixing without generating a large $\tilde{e}_L - h^+$ mixing as well (which would induce an undesirably large $\tilde{W}_3' \to h^+ e_L$ decay rate) or without requiring very fine tuning between the values of v_{ν_i} and μ_i. In the case of one family (putting all indices equal), this can be seen easily. First, the $\mathcal{M}^2_{h_2^- - \tilde{e}^c}$ matrix element is proportional to the neutrino mass and hence very small. Second, the $\mathcal{M}^2_{h_1^- - \tilde{e}^c}$ matrix element, neglecting a small term proportional to the neutrino mass, is proportional to the $\mathcal{M}^2_{\tilde{e}_L - \tilde{e}^c}$ matrix element. Hence it can be shown easily that it is not possible to have a sufficiently large $\mathcal{M}^2_{h_1^- - \tilde{e}^c}$ matrix element (inducing ξ of order 10^{-3}) together with a sufficiently small $\tilde{e}_L - h^-$ mixing . The latter mixing gets a contribution $\sim \xi(\mu/\mu_l)[\mathcal{M}^2_{\tilde{e}_L - \tilde{e}^c}/max(\mathcal{M}^2_{\tilde{e}_L - \tilde{e}_L^*}, \mathcal{M}^2_{\tilde{e}^{c*} - \tilde{e}^c})]$. Now, in the case of three families, due to the A'^e terms in $\mathcal{M}^2_{\tilde{e}_L - \tilde{e}^c}$, both matrix elements are not any more proportional and the $\tilde{e}_L^i - h^+$ mixings can be made as small as necessary independently of the value of ξ. However, for values of $\tilde{e}_L$ and $\tilde{e}^c$ masses of the order 4 TeV or more (see section 4), a value of ξ around 10^{-3} requires that the $\mathcal{M}^2_{h_1^- - \tilde{e}^c}$ matrix element is of order $10^{-3} \times (4\ \text{TeV})^2 \simeq (125\ \text{GeV})^2$ which implies very large values of μ_l and v_{ν_l}. Hence extreme fine tuning between the values of μ_i and v_{ν_i} is needed to obtain a small enough neutrino mass in Eq. (18).

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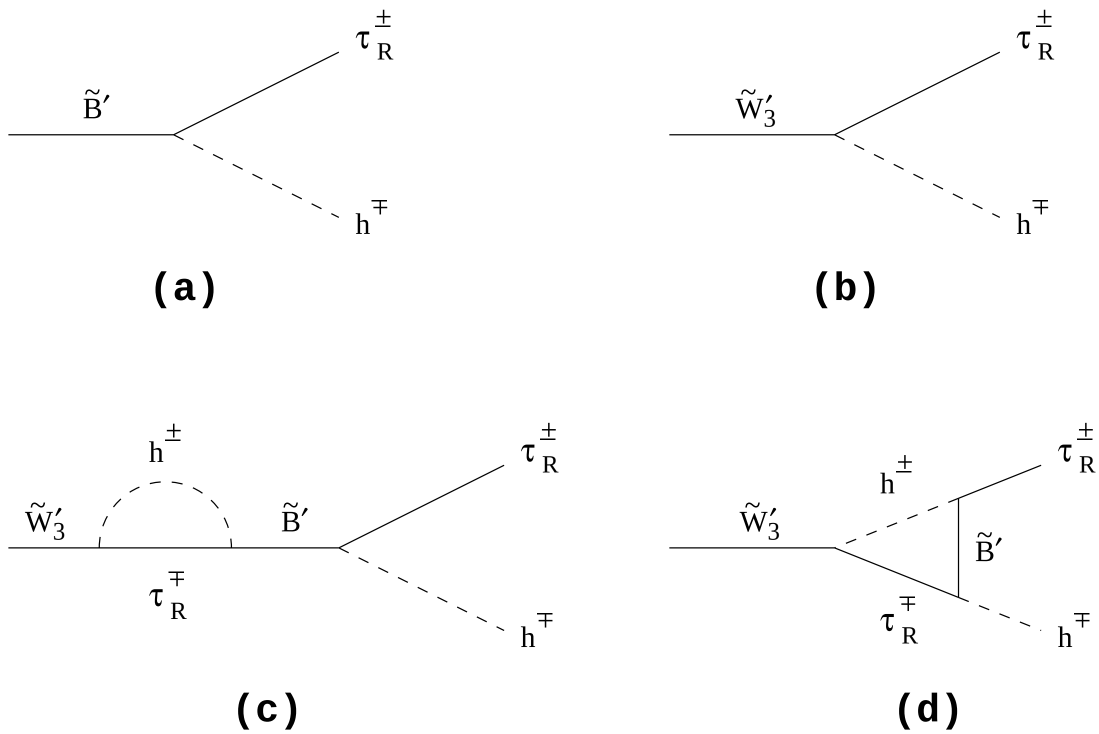

Figure 1: Tree-level diagrams for (a) $\tilde{B}'$ decay and (b) $\tilde{W}_3'$ decay (through their $\tilde{B}$ content), and the one-loop (c) self-energy and (d) vertex diagrams for $\tilde{W}_3'$ decay which have absorptive parts of opposite lepton number.

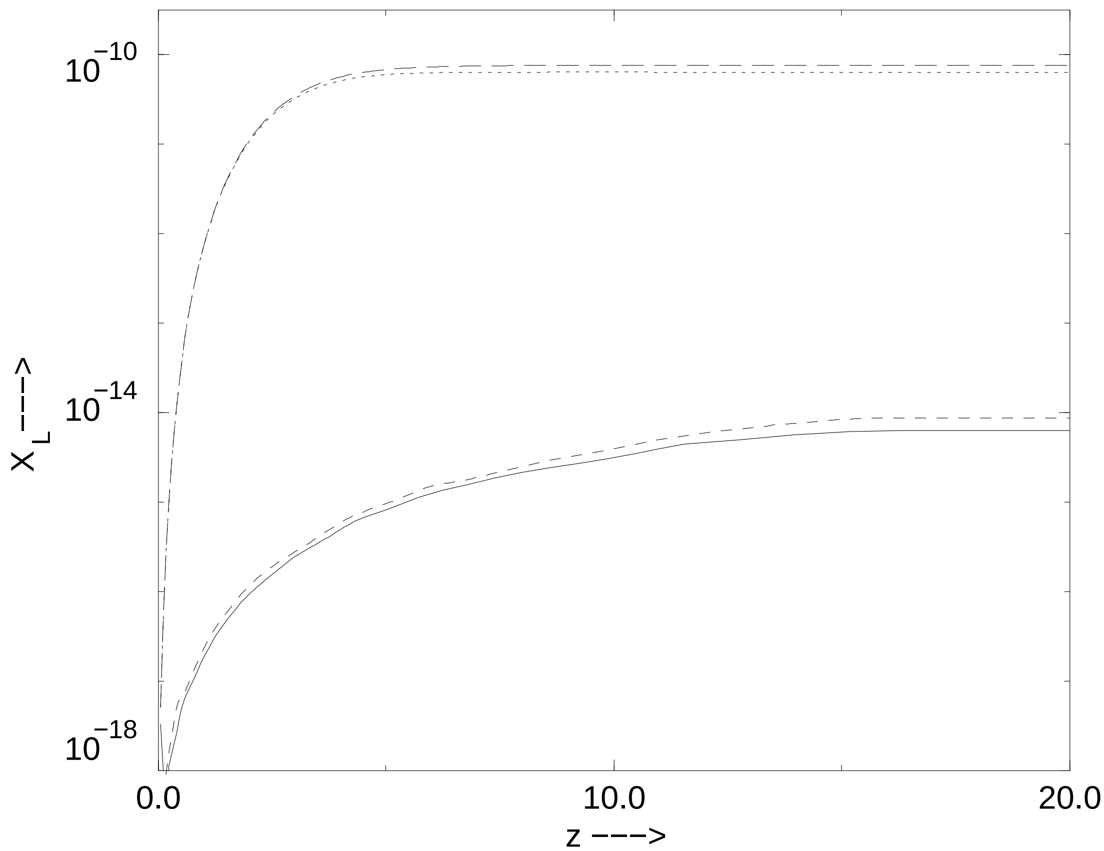

Figure 2: Leptonic asymmetry X_L as a function of $z = M_{\tilde{W}_3'}/T$ as obtained with the parameters given in the text, including all the contributions (solid); taking out the scattering term (short-dashed); not considering in addition the inverse decay of the $\tilde{B}'$ damping term (dotted); and without the inverse decay of the $\tilde{W}_3'$ damping term (long-dashed). In the last case, since all damping terms have been taken out, the asymptotic result is just $X_L = \varepsilon n_\gamma/(2s)$.

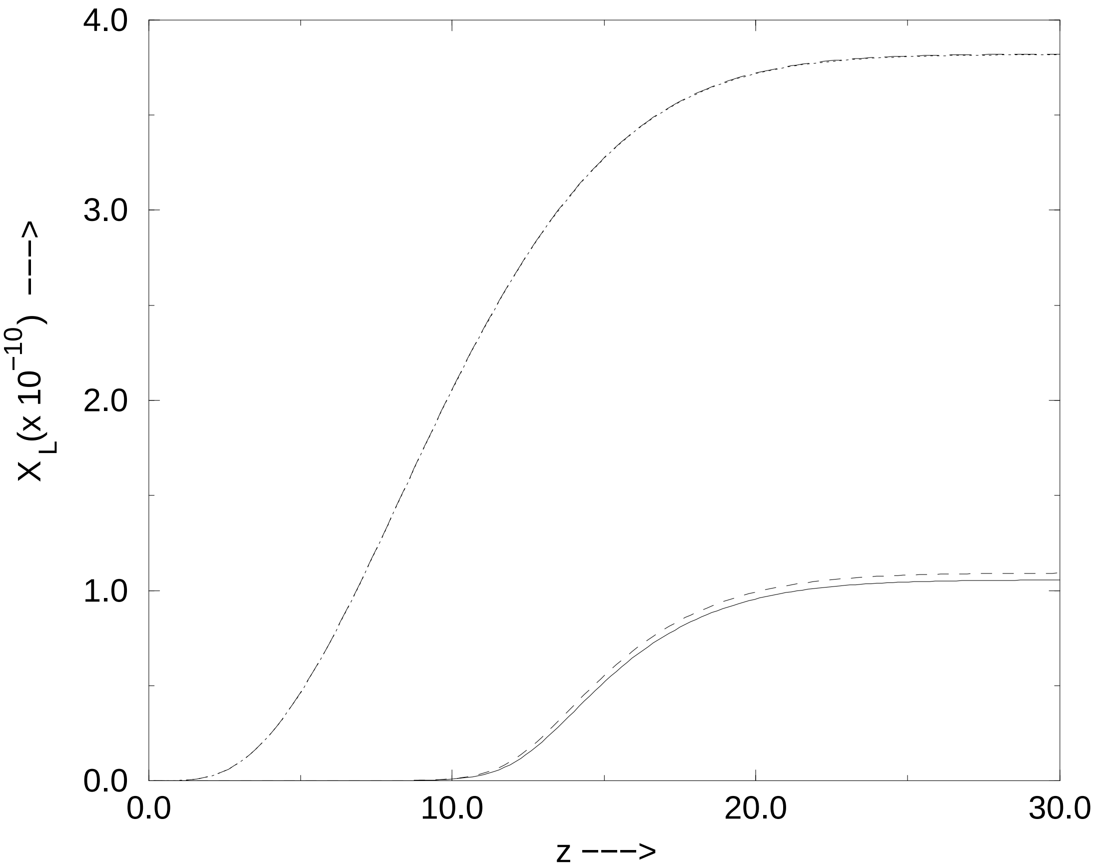

Figure 3: Leptonic asymmetry X_L as a function of $z = M_{\tilde{W}_3'}/T$ as obtained with the parameters given by the set Eq. (47), including all the contributions (solid); taking out the scattering term (short-dashed); not considering in addition the inverse decay of the $\tilde{B}'$ damping term (dotted); and without the inverse decay of the $\tilde{W}_3'$ damping term (long-dashed). In the last case, since all damping terms have been taken out, the asymptotic result is just $X_L = \varepsilon n_\gamma/(2s)$.

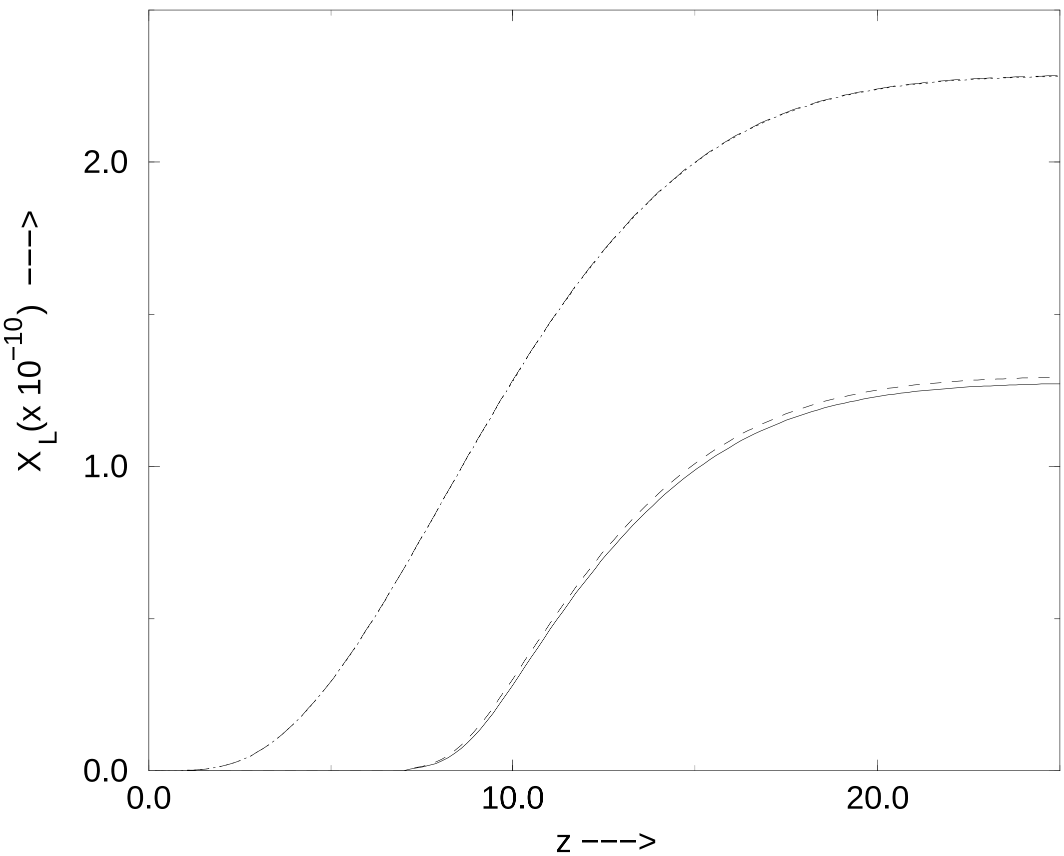

Figure 4: Leptonic asymmetry X_L as a function of $z = M_{\tilde{W}_3'}/T$ as obtained with the parameters given by the set Eq. (48), including all the contributions (solid); taking out the scattering term (short-dashed); not considering in addition the inverse decay of the $\tilde{B}'$ damping term (dotted); and without the inverse decay of the $\tilde{W}_3'$ damping term (long-dashed). In the last case, since all damping terms have been taken out, the asymptotic result is just $X_L = \varepsilon n_\gamma/(2s)$.

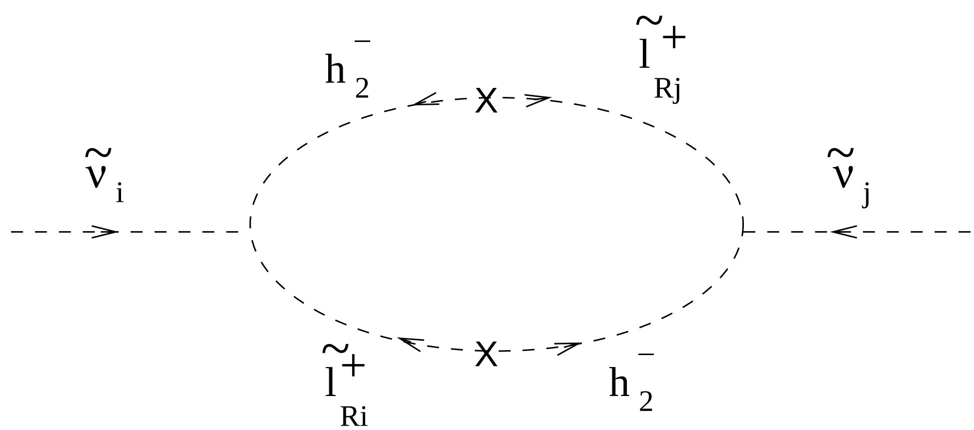

Figure 5: One-loop diagram contributing to the sneutrino “Majorana” mass.

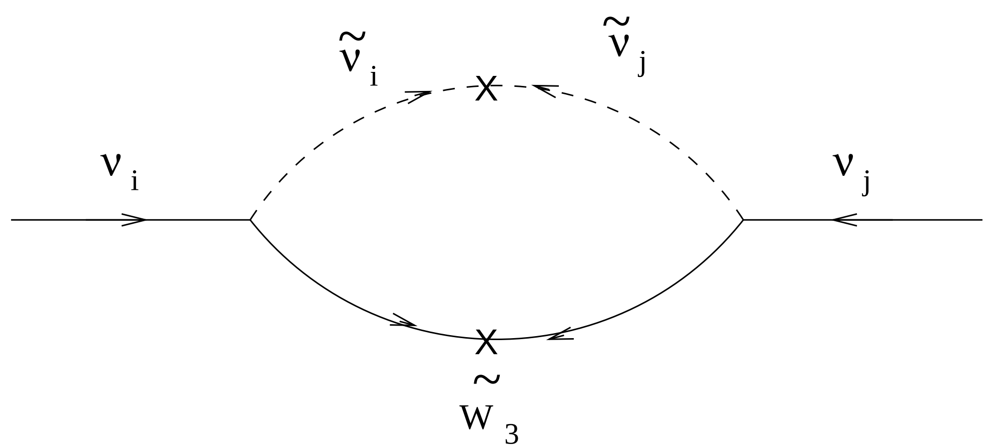

Figure 6: One-loop diagram contributing to the neutrino mass, induced by the sneutrino “Majorana” mass.