Resonant leptogenesis with nonholomorphic R-parity violation and the LHC phenomenology

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In R-parity violating supersymmetric models both leptogenesis and the correct neutrino masses are hard to achieve together. The presence of certain soft nonholomorphic R-parity violating terms helps to resolve this problem. We consider a scenario where the lightest and the second-lightest neutralino are nearly degenerate in mass and enough CP-asymmetry can be produced through resonant leptogenesis. In this model, the lighter chargino and the lightest neutralino are highly degenerate. We have relatively lighter gauginos which can be produced at the LHC (with $\sqrt{s} = 14$ TeV) leading to heavily ionizing charged tracks. At the same time this model can also generate the correct neutrino mass scale. Thus our scenario is phenomenologically rich and testable at colliders.

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I. INTRODUCTION

Baryon Asymmetry of the Universe (BAU) is one of the major challenges of cosmoparticle physics whose theory is not yet convincing. The observed baryon asymmetry is [1]

$$X_B \equiv \frac{n_B - n_{\bar{B}}}{\hat{S}} = \frac{n_B}{\hat{S}} \approx (7.2\text{–}9.2) \times 10^{-11} \quad (95\% \text{ C.L.}), \tag{1}$$

where n_B is the number density of baryons, $n_{\bar{B}}$ is that of antibaryons, and $\hat{S}$ is the entropy density. Leptogenesis [2], leading to lepton asymmetry which partly gets converted into the baryon asymmetry through sphaleron processes, is thought of as a good candidate to describe the matter-antimatter discrepancy of the Universe, i.e., BAU. Within the supersymmetric standard model (SSM), the final baryon asymmetry is related to the initial lepton asymmetry by $B = -\frac{32}{60}L$, where B is the net baryon number and L is the net lepton number of the Universe.

Leptogenesis has drawn a significant attention as it demands lepton number violation with the hope to have possible connection with other lepton number violating processes. The canonical seesaw mechanism [3], one of the most promising ways of explaining the origin of nonzero neutrino mass, also asks for lepton number violation via heavy neutral singlet fermion exchange, i.e., the right-handed neutrino. The decays of these heavy particles can generate enough CP-asymmetry for successful leptogenesis. Thus neutrino mass generation and leptogenesis can be compatible with each other in this case. Another model of neutrino mass generation where leptogenesis occurs naturally is the Higgs triplet model. This is, of course, another realization of seesaw mechanism where a heavy Higgs scalar triplet is exchanged [4]. One should notice that in both these cases lepton-number violation (by two units) occurs at a much higher scale than the scale of electroweak symmetry breaking ($\sim 10^2$ GeV). Note also that in the supersymmetric version of seesaw mechanism R-parity is conserved.

On the other hand, R-parity violating models provide a source of neutrino masses and mixing, which is intrinsically supersymmetric in nature [5]. In R-parity nonconserving SUSY, induced by lepton number violation by odd units, realistic neutrino mass patterns and mixing angles can be generated compatible with the neutrino oscillation and reactor data. However, in the presence of these lepton number violating interactions at the scale of 10^3 GeV or so, any preexisting lepton asymmetry of the Universe would certainly be erased [6].

It was shown in [7,8] that successful leptogenesis is possible in SUSY models with R-parity violation, provided certain nonholomorphic lepton number violating interactions are taken into account. In addition, the familiar R-parity nonconserving interactions must be suppressed at the same time. In such a scenario enough CP-asymmetry can be produced in the decay of the lightest neutralino ($\tilde{\chi}_1^0$) into a charged Higgs boson and a lepton and this suppressed decay can also satisfy the out-of-equilibrium condition leading to successful leptogenesis. On the other hand, the heavier second-lightest neutralino ($\tilde{\chi}_2^0$) must not satisfy the out-of-equilibrium condition and decays very fast. It has also been shown that for a leptogenesis mechanism to be successful in the MSSM with R-parity violation, one must use only those lepton number violating terms, which are not constrained by neutrino masses. In this case the smallness of neutrino masses are explained by a radiative two-loop mechanism involving sneutrino-antisneutrino mass splitting [8].

It has been shown in Ref. [8] that one needs a very heavy spectrum of SUSY particles to get correct values of the lepton asymmetry. In fact, a hierarchical scenario has been

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considered with the assumption that the bino-dominated $\tilde{\chi}_2^0$ is heavier than the wino-dominated $\tilde{\chi}_1^0$, i.e., $M_1 > M_2$. Here, M_1 and M_2 are the $U(1)$ and $SU(2)$ gaugino mass parameters, respectively. Successful leptogenesis required that the gaugino masses must be in the range of 2–6 TeV and hence this scenario might have a very remote possibility to be tested at the LHC or the future ILC. We note in passing that the scales involved in the canonical leptogenesis models are very high and impossible to be tested directly at the present or upcoming high-energy colliders.

On the other hand, it was shown in [9] that in models where leptogenesis is driven by the decays of the right-handed neutrino (M_{N_i}), the *CP*-asymmetry can be enhanced for two nearly degenerate right-handed neutrinos. This is because in the limit $M_{N_i} - M_{N_j} \ll M_{N_i}$ the self-energy diagram dominates and due to this *resonance* effect the mass scale of the right-handed neutrinos can be lowered significantly. It was noted that sufficient lepton asymmetry can be generated even with $M_N \sim 1$ TeV [10]. The presence of TeV scale right-handed neutrinos makes this scenario phenomenologically interesting compared to other canonical leptogenesis scenarios. With this motivation, we revisited the model considered in Refs. [7,8] and focused on a scenario where we have nearly degenerate $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$. We found that enough lepton asymmetry can be generated in this case through *Resonant Leptogenesis*, which can be converted into baryon asymmetry through the sphaleron processes. In our case, the masses of $\tilde{\chi}_2^0$, $\tilde{\chi}_1^0$ and the lighter chargino $\tilde{\chi}_1^{\pm}$ are found to be around 1 TeV and thus they have a possibility to be produced at the LHC. In addition, the lighter chargino is nearly degenerate with the lighter neutralinos and can have a very slow decay leading to heavily ionizing charged tracks in the collider detector. This could be a crucial test of the present model trying to explain leptogenesis and hence the baryon asymmetry of the Universe.

II. SOFT SUPERSYMMETRY BREAKING AND R-PARITY VIOLATION WITH NONHOLOMORPHIC TERMS

In a supersymmetric theory neither gauge invariance nor supersymmetry requires the conservation of lepton and baryon number. However, the lepton and baryon number violating operators can induce a fast rate of proton decay and violate its present experimental bound. To avoid this calamity, a discrete symmetry called R-parity was introduced, which is defined as

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

where B is the baryon number, L the lepton number, and S the spin angular momentum. It is easy to check that the standard model particles have $\mathrm{R} = +1$ and their supersymmetric partners have $\mathrm{R} = -1$. An immediate consequence of R-parity conservation is that the lightest supersymmetric particle (LSP) is stable. On the other hand, one notices that proton decay is still forbidden if either baryon number or lepton number is conserved in nature and R-parity is violated. This has led to considerable theoretical and phenomenological interest in studying models in which R-parity is violated.

In an R-parity violating model, the superpotential can be written as

$$W = W_{\mathrm{MSSM}} + W_{\mathrm{RPV}}. \tag{3}$$

Here, W_{MSSM} is the superpotential of R-parity conserving minimal supersymmetric standard model (MSSM) and is given by

$$W_{\mathrm{MSSM}} = \mu H_1 H_2 + f_{ij}^e H_1 L_i e_j^c + f_{ij}^d H_1 Q_i d_j^c + f_{ij}^u H_2 Q_i u_j^c, \tag{4}$$

whereas the R-parity violating part of the superpotential is given by

$$W_{\mathrm{RPV}} = \mu_i L_i H_2 + \lambda_{ijk} L_i L_j e_k^c + \lambda'_{ijk} L_i Q_j d_k^c + \lambda''_{ijk} u_i^c d_j^c d_k^c. \tag{5}$$

Here, i, j, k are generational indices, f_{ij}^u, f_{ij}^d, and f_{ij}^e are (3×3) Yukawa matrices, Q, u^c, and d^c are the quark doublet and singlet superfields and L and e^c are the lepton doublet and singlet superfields. The two Higgs doublet superfields are H_1 and H_2 giving rise to masses for the down-type quarks (and charged leptons) and the up-type quarks, respectively. The parameters μ and μ_i have dimensions of mass and the terms $\mu_i L_i H_2$ are called the bilinear R-parity violating interactions whereas the terms involving λ, λ' and λ'' are called trilinear R-parity violating interactions.

Once supersymmetry is broken the soft supersymmetry breaking terms conserving R-parity and allowed by the standard model gauge group can be written as

$$\begin{aligned}
\mathcal{L}_{\mathrm{soft}} = & -\tilde{L}_i^{a*}(M_L^2)_{ij}\tilde{L}_j^a - \tilde{e}_i^{c*}(M_e^2)_{ij}\tilde{e}_j^c - \tilde{Q}_i^{a*}(M_Q^2)_{ij}\tilde{Q}_j^a \\
& - \tilde{u}_i^{c*}(M_u^2)_{ij}\tilde{u}_j^c - \tilde{d}_i^{c*}(M_d^2)_{ij}\tilde{d}_j^c - M_{H_1}^2 H_1^{a*} H_1^a \\
& - M_{H_2}^2 H_2^{a*} H_2^a - \varepsilon_{ab}(B H_1^a H_2^b + \mathrm{H.c.}) \\
& - \varepsilon_{ab}[(A_e f_e)_{ij} H_1^a \tilde{L}_i^b \tilde{e}_j^c + (A_u f_u)_{ij} H_2^b \tilde{Q}_i^a \tilde{u}_j^c \\
& + (A_d f_d)_{ij} H_1^a \tilde{Q}_i^b \tilde{d}_j^c + \mathrm{H.c.}] \\
& - \tfrac{1}{2}(M_3 \tilde{g}\,\tilde{g} + M_2 \tilde{W}\,\tilde{W} + M_1 \tilde{B}\,\tilde{B} + \mathrm{H.c.}).
\end{aligned} \tag{6}$$

Here, a are $SU(2)$ indices. M_3, M_2, and M_1 are the $SU(3)$, $SU(2)$, and $U(1)$ gaugino mass parameters, respectively. A_e, A_d, and A_u are the trilinear scalar couplings and B is the Higgs bilinear parameter. The Higgs doublets giving mass to the standard model fermions are

$$H_1 = \begin{pmatrix} h_1^0 \\ h_1^- \end{pmatrix}, \qquad H_2 = \begin{pmatrix} h_2^+ \\ h_2^0 \end{pmatrix}. \tag{7}$$

In an R-parity violating theory, additional soft terms may be present, and can be written as

$$\mathcal{L}_{\rm soft}^{\not{R}} = -\varepsilon_{ab}(B_i^{\prime}\tilde{L}_i^a H_2^b + A_{ijk}^{\prime e}\tilde{L}_i^a\tilde{L}_j^b\tilde{e}_k^c + A_{ijk}^{\prime d}\tilde{Q}_i^a\tilde{L}_j^b\tilde{d}_k^c) - A_{ijk}^{\prime S}\tilde{u}_i^c\tilde{d}_j^c\tilde{d}_k^c + \text{H.c.} \quad (8)$$

Following the convention of Ref. [8], 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.

In principle, there could be nonholomorphic terms in the Lagrangian. These nonholomorphic terms appear in the Lagrangian of the visible sector as an artifact of SUSY breaking in the hidden sector [11–14].

The most general set of nonholomorphic soft terms conserving R-parity is

$$\mathcal{L}_{\rm soft}^{\rm NH} = -N_{ij}^e H_2^{a*}\tilde{L}_i^a\tilde{e}_j^c - N_{ij}^d H_2^{a*}\tilde{Q}_i^a\tilde{d}_j^c - N_{ij}^u H_1^{a*}\tilde{Q}_i^a\tilde{u}_j^c + \text{H.c.} \quad (9)$$

Similarly, nonholomorphic soft terms breaking R-parity are

$$\mathcal{L}_{\rm soft}^{\rm NH\not{R}} = -N_i^{\prime B}H_1^{a*}\tilde{L}_i^a - N_i^{\prime e}H_2^{a*}H_1^a\tilde{e}_i^c - N_{ijk}^{\prime u}\tilde{L}_i^{a*}\tilde{Q}_j^a\tilde{u}_k^c - N_{ijk}^{\prime S}\tilde{u}_i^c\tilde{e}_j^c\tilde{d}_k^{c*} - N_{ijk}^{\prime d}\varepsilon_{ab}\tilde{Q}_i^a\tilde{Q}_j^b\tilde{d}_k^{c*} + \text{H.c.} \quad (10)$$

In this paper, we shall assume lepton-number violation but baryon-number conservation in the Lagrangian. This implies that $\lambda_{ijk}^{\prime\prime} = A_{ijk}^{\prime S} = N_{ijk}^{\prime d} = 0$.

If lepton number is violated by the bilinear R-parity breaking interactions $\mu_i L_i H_2$, then this induces mixing between the neutrinos with the MSSM neutralinos. In addition, the sneutrinos ($\tilde{\nu}_i$) may all acquire nonzero vacuum expectation values (VEVs). In this case, the neutralino mass matrix gets enhanced to a (7×7) mass matrix and in the basis $[\tilde{B}, \tilde{W}_3, \tilde{h}_1^0, \tilde{h}_2^0, \nu_1, \nu_2, \nu_3]$ is given by

$$\mathcal{M} = \begin{bmatrix} M_1 & 0 & -sr_Z v_1 & sr_Z v_2 & -sr_Z v_{\nu_1} & -sr_Z v_{\nu_2} & -sr_Z v_{\nu_3} \\ 0 & M_2 & cr_Z v_1 & -cr_Z v_2 & cr_Z v_{\nu_1} & cr_Z v_{\nu_2} & cr_Z v_{\nu_3} \\ -sr_Z v_1 & cr_Z v_1 & 0 & -\mu & 0 & 0 & 0 \\ sr_Z v_2 & -cr_Z v_2 & -\mu & 0 & -\mu_1 & -\mu_2 & -\mu_3 \\ -sr_Z v_{\nu_1} & cr_Z v_{\nu_1} & 0 & -\mu_1 & 0 & 0 & 0 \\ -sr_Z v_{\nu_2} & cr_Z v_{\nu_2} & 0 & -\mu_2 & 0 & 0 & 0 \\ -sr_Z v_{\nu_3} & cr_Z v_{\nu_3} & 0 & -\mu_3 & 0 & 0 & 0 \end{bmatrix}, \quad (11)$$

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 $v_1^2 + v_2^2 + v_\nu^2 = v^2 \simeq (246\ \text{GeV})^2$ and $v_\nu^2 = v_{\nu_1}^2 + v_{\nu_2}^2 + v_{\nu_3}^2$. We also define the parameter $\tan\beta = v_2/(v_1^2 + v_\nu^2)^{1/2}$.

In order to understand how a nonzero neutrino mass arises at the tree level from the above (7×7) mass matrix, let us assume that μ is much larger compared to the other entries. This means that $\tilde{h}_{1,2}^0$ form a heavy Dirac particle of mass μ which mixes very little with the other physical fields. Integrating out these heavy fields one can write down the reduced (5×5) matrix using seesaw formula in the basis $[\tilde{B}, \tilde{W}_3, \nu_1, \nu_2, \nu_3]$ as

$$\mathcal{M} = \begin{bmatrix} M_1 - s^2\delta r & sc\delta r & -s\epsilon_1 & -s\epsilon_2 & -s\epsilon_3 \\ sc\delta r & M_2 - c^2\delta r & c\epsilon_1 & c\epsilon_2 & c\epsilon_3 \\ -s\epsilon_1 & c\epsilon_1 & 0 & 0 & 0 \\ -s\epsilon_2 & c\epsilon_2 & 0 & 0 & 0 \\ -s\epsilon_3 & c\epsilon_3 & 0 & 0 & 0 \end{bmatrix}, \quad (12)$$

where

$$\delta = 2M_Z^2\frac{v_1 v_2}{v^2}\frac{1}{\mu} = \frac{M_Z^2\sin 2\beta}{\mu}\sqrt{1 - \frac{v_\nu^2}{v^2\cos^2\beta}}, \quad (13)$$

$$\epsilon_i = \frac{M_Z}{v}\left(v_{\nu_i} - \frac{\mu_i}{\mu}v_1\right), \quad (14)$$

$$r = (1 + M_2/\mu\,\sin 2\beta)/(1 - M_2^2/\mu^2). \quad (15)$$

Here, the quantity r has been introduced as a correction factor for finite values of M_2/μ.

Looking at Eq. (12), one can see that 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. The other two orthogonal combinations decouple from the neutralino mass matrix. In this case, only the eigenstate

$$\nu'_l = \nu_l + \frac{s\epsilon}{M_1}\tilde{B} - \frac{c\epsilon}{M_2}\tilde{W}_3, \tag{16}$$

gets a seesaw mass given by

$$m_{\nu'_l} = -\epsilon^2\left(\frac{s^2}{M_1} + \frac{c^2}{M_2}\right), \tag{17}$$

whereas the other two neutrinos remain massless. These massless neutrino states may get a nonzero contribution to their masses through one-loop radiative corrections [15].

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

$$\tilde{\chi}^0_2 = \tilde{B} + \frac{sc\delta r}{M_1 - M_2}\tilde{W}_3 - \frac{s\epsilon}{M_1}\nu_l, \tag{18}$$

$$\tilde{\chi}^0_1 = \tilde{W}_3 - \frac{sc\delta r}{M_1 - M_2}\tilde{B} + \frac{c\epsilon}{M_2}\nu_l. \tag{19}$$

Because of this nonzero neutrino component of the physical states $\tilde{\chi}^0_2$ and $\tilde{\chi}^0_1$, they can decay to lepton number violating two-body final states such as $\tilde{\chi}^0_1 \to l^{\mp}W^{\pm}$ [16,17]. In general, the gaugino masses can be complex and this can induce CP violation in the neutralino sector. Hence, a lepton asymmetry can be generated from the lepton number violating decays of the neutralinos. However, the asymmetry generated this way is much smaller than the required value of $\sim 10^{-10}$ [8]. This is because the asymmetry has to be much less than $(\epsilon/M_{1,2})^2$ [see, Eq. (19)]. The quantity $(\epsilon/M_{1,2})^2$ is of order $m_{\nu'_l}/M_{1,2}$ [see, Eq. (17)], and hence the asymmetry is $<5 \times 10^{-13}$ 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 results in an upper bound on $(\epsilon/M_{1,2})^2$, which is independent of $m_{\nu'_l}$. This effect also makes the asymmetry to be very much less than 10^{-10}. Even if one considers R-parity violating trilinear couplings λ and λ', it is possible to show [8] that they are also not compatible with the successful generation of a lepton or baryon asymmetry of the Universe.

III. *CP*-ASYMMETRY AND RESONANT LEPTOGENESIS FROM NEUTRALINO DECAYS

For a leptogenesis mechanism to be successful in the MSSM with R-parity violation, one needs to satisfy two requirements. First, the lepton-number violating terms must not be 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. In the line of [8], let us first assume that the bino $\tilde{B}$ is heavier than the wino $\tilde{W}_3$, although we will look at the scenario where the mass difference is very small. Because of R-parity violation, left- and right-chiral charged sleptons mix with the charged Higgs boson. Now if one assumes that the left-chiral charged slepton has a negligible mixing with the charged Higgs boson, then the $\tilde{\chi}^0_1$ decay into $l^{\mp}h^{\pm}$ is suppressed as long as the wino-bino mixing is small. This can be achieved if $\mu \gg M_1, M_2$. Hence, the heavier neutralino $\tilde{\chi}^0_2$ decays quickly and the lighter neutralino $\tilde{\chi}^0_1$ has a much slower decay. At temperatures well above $T = M_{\rm SUSY}$, 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 (bino mass), when $\tilde{\chi}^0_1$ interactions violate L_i as well as $(B - 3L_i)$ for $i = e, \mu, \tau$. Let us consider here that all other supersymmetric particles are heavier than the neutralinos, so that at temperatures below M_1 we need to consider only the interactions of $\tilde{\chi}^0_2$ and $\tilde{\chi}^0_1$.

We start with the well-known interaction of $\tilde{B}$ with l and $\tilde{l}_R$ given by [18]

$$-\frac{e\sqrt{2}}{\cos\theta_W}\left[\bar{l}\left(\frac{1-\gamma_5}{2}\right)\tilde{B}\tilde{l}_R + \text{H.c.}\right]. \tag{20}$$

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 $\tilde{\chi}^0_1$ of Eq. (19) with l and $h^{\pm}$ is given by

$$\left(\frac{sc\xi\delta r}{M_1 - M_2}\right)\left(\frac{e\sqrt{2}}{\cos\theta_W}\right)\left[\bar{l}\left(\frac{1-\gamma_5}{2}\right)\tilde{\chi}^0_1 h^- + \text{H.c.}\right], \tag{21}$$

where ξ represents the $\tilde{l}_R - h^-$ mixing because of nonholomorphic R-parity violation and is assumed to be real. In the absence of nonholomorphic terms it is very hard to generate a large right-handed slepton and charged Higgs mixing without generating a large left-handed slepton and charged Higgs mixing as well. In order to achieve this we assume that B'_i and $N_i^{\prime B}$ in Eqs. (8) and (10) are negligible, thus the left-handed slepton and charged Higgs do not mix heavily. The term $N_i^{\prime e}$ produces the mixing between $\tilde{l}_R$ and charged Higgs. ξ is proportional to $N_i^{\prime e}$ and measures the strength of the nonholomorphic coupling. However, the parameter δ of Eq. (13) is complex. The nontrivial CP phase in the above contributes negligibly to the neutron electric dipole moment because the magnitude of δ is very small [8].

In Fig. 1 we show the lepton number violating decay processes (a) $\tilde{\chi}^0_2 \leftrightarrow l^{\pm}_R h^{\mp}$ and (b) $\tilde{\chi}^0_1 \leftrightarrow l^{\pm}_R h^{\mp}$, at the tree level as well as the one-loop (c) self-energy and (d) vertex corrections of $\tilde{\chi}^0_1$ decay. The decay width of $\tilde{\chi}^0_2$ is given by

$$\begin{aligned}\Gamma_{\tilde{\chi}^0_2} &= \Gamma(\tilde{\chi}^0_2 \to l^+h^-) + \Gamma(\tilde{\chi}^0_2 \to l^-h^+)\\ &= \frac{1}{4\pi}\xi^2\frac{e^2}{c^2}\frac{(M^2_{\tilde{\chi}^0_2} - m^2_h)^2}{M^3_{\tilde{\chi}^0_2}},\end{aligned} \tag{22}$$

while that of the $\tilde{\chi}^0_1$ is

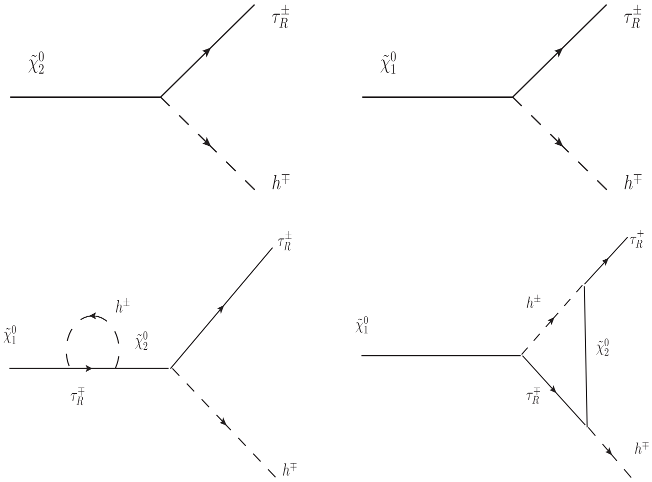

FIG. 1. Tree-level diagrams for (a) $\tilde{\chi}_2^0$ decay and (b) $\tilde{\chi}_1^0$ decay (through their bino content), and the one-loop (c) self-energy and (d) vertex correction diagrams for $\tilde{\chi}_1^0$ decay.

$$\Gamma_{\tilde{\chi}_1^0} = \Gamma(\tilde{\chi}_1^0 \to l^+ h^-) + \Gamma(\tilde{\chi}_1^0 \to l^- h^+)$$
$$= \frac{1}{4\pi}\xi^2\left(\frac{es|\delta|r}{M_1 - M_2}\right)^2 \frac{(M_{\tilde{\chi}_1^0}^2 - m_h^2)^2}{M_{\tilde{\chi}_1^0}^3}. \qquad (23)$$

Here, m_h is the mass of the charged Higgs boson (h^-). Let us also mention that for our choice of parameter (shown later) the radiative decay $\tilde{\chi}_2^0 \to \tilde{\chi}_1^0\gamma$ and the 3-body decay $\tilde{\chi}_2^0 \to \tilde{\chi}_1^0 f\bar{f}$ are very much suppressed and do not contribute to the decay width.

Calculating the interference between the tree-level and self-energy + vertex correction diagrams of Fig. 1 one obtains the following CP-asymmetry from the decay of $\tilde{\chi}_1^0$ [8]:

$$\in = \frac{\Gamma(\tilde{\chi}_1^0 \to l^+ h^-) - \Gamma(\tilde{\chi}_1^0 \to l^- h^+)}{\Gamma_{\tilde{\chi}_1^0}} \qquad (24)$$

$$= \frac{\alpha\xi^2}{2\cos^2\theta_W}\frac{\mathrm{Im}\delta^2}{|\delta|^2}\left(1 - \frac{m_h^2}{M_{\tilde{\chi}_1^0}^2}\right)^2 \frac{x^{1/2} f(x)}{(1-x)}, \qquad (25)$$

where $x = M_{\tilde{\chi}_1^0}^2 / M_{\tilde{\chi}_2^0}^2$ and

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

If the $\tilde{\chi}_1^0$ decay rate satisfies the out-of-equilibrium condition, then a lepton asymmetry may be generated from the above decay asymmetry.

If the sphaleron interactions [19] are still in equilibrium during the generation of this lepton asymmetry, they will convert it into a baryon asymmetry of the Universe [20]. For a strongly first-order electroweak phase transition, the sphaleron interactions freeze out at the temperature of the phase transition (the critical temperature). For a second-order or weakly first-order phase transition, the sphaleron interactions freeze out at a temperature lower than the critical temperature [21]. In order to calculate the final baryon asymmetry, the precise values of the critical temperature and the freeze out temperature would have to be determined. Here we shall assume that the sphaleron interactions freeze out at a temperature $\sim$100 GeV. So, as long as the lepton asymmetry is generated at a temperature above, say $\sim$100 GeV, it will be converted to a baryon asymmetry of the Universe using the relation $B = -\frac{32}{60}L$.

If the $\tilde{\chi}_1^0$ decay rate is much less than the expansion rate of the Universe, the generated lepton asymmetry is of the order of the decay asymmetry given in Eq. (25). In other words, the out-of-equilibrium condition reads as:

$$K_{\tilde{\chi}_1^0} = \frac{\Gamma_{\tilde{\chi}_1^0}}{H(M_{\tilde{\chi}_1^0})} \ll 1, \qquad (27)$$

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

$$H(T) = \sqrt{\frac{4\pi^3 g_*}{45}}\frac{T^2}{M_{\mathrm{Planck}}}, \qquad (28)$$

with g_* the number of massless degrees of freedom which is 106.75 in this case corresponding to the standard model (SM) degrees of freedom and $M_{\mathrm{Pl}} \sim 10^{19}$ GeV is the Planck scale.[1]

However, in order to present a realistic and reliable estimation of the lepton asymmetry generated from neutralino decay, we solve the full Boltzmann equations [22]. In our scenario, we have the Boltzmann equations (including the washout effects) the same as in [8],

$$\frac{dX_{\tilde{\chi}_1^0}}{dz} = -zK_{\tilde{\chi}_1^0}\frac{K_1(z)}{K_2(z)}(X_{\tilde{\chi}_1^0} - X_{\tilde{\chi}_1^0}^{\mathrm{eq}}),$$
$$\frac{dX_L}{dz} = zK_{\tilde{\chi}_1^0}\frac{K_1(z)}{K_2(z)}\left[\varepsilon(X_{\tilde{\chi}_1^0} - X_{\tilde{\chi}_1^0}^{\mathrm{eq}}) - \frac{1}{2}\frac{X_{\tilde{\chi}_1^0}}{X_\gamma}X_L\right]$$
$$- z\left(\frac{M_{\tilde{\chi}_2^0}}{M_{\tilde{\chi}_1^0}}\right)^2 K_{\tilde{\chi}_2^0}\left[\frac{1}{2}\frac{K_1(zM_{\tilde{\chi}_2^0}/M_{\tilde{\chi}_1^0})}{K_2(zM_{\tilde{\chi}_2^0}/M_{\tilde{\chi}_1^0})}\frac{X_{\tilde{\chi}_2^0}}{X_\gamma}X_L\right.$$
$$\left. + 2\frac{X_L}{X_\gamma}\frac{\gamma_{\mathrm{scatt.}}^{\mathrm{eq}}}{\hat{S}\Gamma_{\tilde{\chi}_2^0}}\right], \qquad (29)$$

where K_1 and K_2 are the modified Bessel's functions, $z \equiv M_{\tilde{\chi}_1^0}/T$, $K_{\tilde{\chi}_2^0} = \Gamma_{\tilde{\chi}_2^0}/H(M_{\tilde{\chi}_2^0})$, and $\hat{S} = g_* \frac{2\pi^2}{45}T^3$ is the entropy density. The number densities per comoving volume have been defined as $X_i = n_i/\hat{S}$ in terms of the number densities of particles 'i'.

Once the washout effects are included [as in Eq. (29)], it is very hard to generate the lepton asymmetry of correct order [$O(10^{-10})$] keeping nearly degenerate neutralinos

[1]In our numerical analysis we have used $M_{\mathrm{Pl}} = 0.9 \times 10^{19}$ GeV.

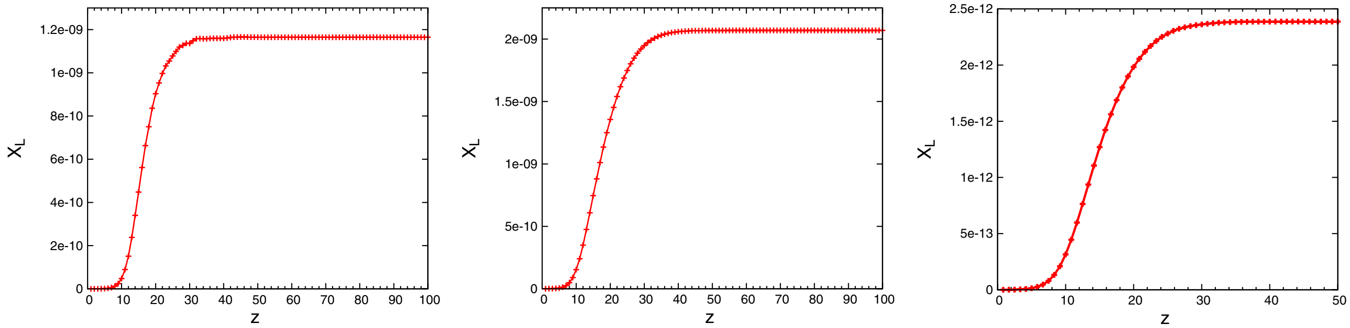

FIG. 2 (color online). Lepton asymmetry (X_L) vs z with $g_* = 106.75$ for Set I (left), Set II (middle), and Set III (right).

$\sim$1 TeV. In order to find a reliable and stable solution we need very large values of $\mu \sim 40$–75 TeV and $\tan\beta \sim 60$–70. We consider three sets of parameters to estimate the lepton asymmetry:

$$\begin{aligned} \text{Set I: } M_{\tilde{\chi}_2^0} &= 1520.998 \text{ GeV}, \\ M_{\tilde{\chi}_1^0} &= 1520.997 \text{ GeV}, \qquad \tan\beta = 64, \\ \xi &= 1.85 \times 10^{-5}, \qquad \mu = 75 \text{ TeV}, \\ M_1 &= 1521 \text{ GeV}, \qquad M_2 = 1520.3 \text{ GeV}, \end{aligned} \tag{30}$$

$$\begin{aligned} \text{Set II: } M_{\tilde{\chi}_2^0} &= 1380.9998 \text{ GeV}, \\ M_{\tilde{\chi}_1^0} &= 1380.9997 \text{ GeV}, \qquad \tan\beta = 72, \\ \xi &= 0.64 \times 10^{-5}, \qquad \mu = 43 \text{ TeV}, \\ M_1 &= 1381 \text{ GeV}, \qquad M_2 = 1380.4 \text{ GeV}, \end{aligned} \tag{31}$$

and

$$\begin{aligned} \text{Set III: } M_{\tilde{\chi}_2^0} &= 1680.9 \text{ GeV}, \qquad M_{\tilde{\chi}_1^0} = 1680.8 \text{ GeV}, \\ \tan\beta &= 65, \qquad \xi = 0.67 \times 10^{-5}, \qquad \mu = 43 \text{ TeV}, \\ M_1 &= 1681.0 \text{ GeV}, \qquad M_2 = 1680.4 \text{ GeV}, \end{aligned} \tag{32}$$

with $m_h = 180$ GeV, $M_Z = 91.19$ GeV.

The resulting evolution of the lepton asymmetry (X_L) is shown in Fig. 2 (left and middle) for the parameter choices of Set I and Set II, respectively. We see from these two figures that a large asymmetry of order 10^{-10} is produced at $T \sim 100$ GeV provided we have nearly degenerate $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$, very large μ and $\tan\beta$ and values of ξ around 1×10^{-5}. In Fig. 2 (right), the lepton asymmetry is shown for the parameters given in Set III. Note that in order to have a large lepton asymmetry the splitting between $M_{\tilde{\chi}_2^0}$ and $M_{\tilde{\chi}_1^0}$ is required to be much smaller compared to the splitting between the gaugino parameters M_1 and M_2. But when the splitting between $M_{\tilde{\chi}_2^0}$ and $M_{\tilde{\chi}_1^0}$ is increased a little bit, lepton asymmetry falls very sharply. This shows the importance of the requirement of very highly degenerate neutralinos. On the other hand, this highly degenerate neutralino scenario (with larger splitting between M_1 and M_2) is difficult to achieve in practice and might need additional fine-tuning in this model or could be taken as a hint in favor of nonminimal SUSY models. It is also important to note that the above discussion is based on the tree-level neutralino mass matrix. One should also include radiative corrections at one-loop order to check the stability of the results presented here. However, that is a subject of a separate study and we will not take it up in the present paper. Our main objective here is to present the idea of resonant leptogenesis in the MSSM with nonholomorphic R-parity violating soft SUSY breaking interactions and its testability at the LHC.

IV. LARGE VALUES OF $\tan\beta$ AND μ

In the general MSSM scenario, the maximum value of $\tan\beta \lesssim 50$ is restricted by the perturbative limit of the supersymmetric Yukawa couplings. Nevertheless, the possibility of large values of $\tan\beta (\gg 50)$ has been considered in various context within the MSSM. For example, in the context of up-down Yukawa unification and Higgs mediated FCNC, $\tan\beta \gg 50$ has been pointed out [23,24]. Some other aspects of Higgs phenomenology with $\tan\beta$ as large as 130 have been studied in [25]. Very recently the authors of [26] have pointed out that the parameter space of MSSM includes a region where the down-type fermion masses are generated by the loop induced couplings to the up-type Higgs doublet. In this region of MSSM, a large value of $\tan\beta \gtrsim 100$ is consistent with the perturbativity of the SUSY Yukawa couplings of the down-type fermions to H_1. In Fig. 3, we have shown the variation of lepton asymmetry with $\tan\beta$, keeping $\tan\beta$ large. From the above plot it is clear that as we increase $\tan\beta$ the lepton asymmetry increases keeping other parameters fixed. This justifies our choice of large value for $\tan\beta$.

The bilinear R-parity conserving term $\mu H_1 H_2$ in Eq. (4) introduces μ as a free parameter of the theory. There is no known symmetry that protects μ from having a value $\sim M_{\rm Pl}$. However, from phenomenological point of view one would expect that the value of μ should be around 100 GeV or 1 TeV scale to avoid unnatural fine-tuning in the theory. This is evident from the electroweak symmetry breaking condition that connects μ with the mass of

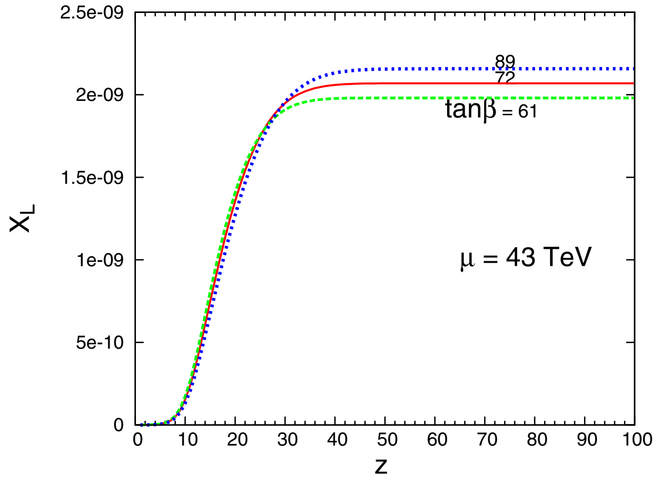

FIG. 3 (color online). Lepton asymmetry vs $\tan\beta$ with $g_* = 106.75$ for Set II.

the Z-boson by the following relation (in the limit of large $\tan\beta$) [27]

$$m_Z^2 = -2(|\mu|^2 + m_{H_2}^2) + \frac{2}{\tan^2\beta}(m_{H_1}^2 - m_{H_2}^2) + \mathcal{O}(1/\tan^4\beta). \tag{33}$$

In order to have the correct value for m_Z, the input parameters $m_{H_2}^2$, $m_{H_1}^2$, and μ on the right-hand side of Eq. (33) should be within an order or two of magnitude of m_Z^2 in the absence of any fine cancellation between various terms. However, if one admits some amount of fine-tuning then the value of μ in the range of 40–75 TeV is possible. It must be noted here that typical viable solutions for the MSSM still requires significant cancellation [27]. We have seen in this paper that to achieve a reliable and stable solution for the lepton asymmetry, μ needs to be very large, i.e., in the range of 40–75 TeV. Thus this situation might be more fine-tuned than the general MSSM scenario. We note in passing that in some studies [28] in the context of Lepton Flavor Violation (LFV) in the Higgs boson decay large values of μ ($\sim$ 25 TeV) has been suggested.

The results from the studies cited above can also be applied in the scenario under consideration. For example, Ref. [25] pointed out that μ has to be large and positive and $\tan\beta > 50$. Thus in our case we find solutions for the lepton asymmetry for low bino and wino masses ($\sim$ 1 TeV) while μ is very large $\sim$40 TeV and $\tan\beta > 50$. Having rather low $M_{\tilde{\chi}_1^0}$, $M_{\tilde{\chi}_1^\pm}$, and $M_{\tilde{\chi}_2^0}$ can give us possible distinct and interesting signatures at the LHC.

Because of large μ, the higgsino sector is decoupled from the wino and bino sectors. The third and the fourth neutralino as well as the heavier chargino are very massive (as their masses are controlled mostly by μ) and are out of reach of the LHC. On the other hand, the lightest and the next-to-lightest neutralino and the lighter chargino are within the reach of the LHC. In addition, since the lighter chargino is also nearly degenerate with the lighter neutralinos, it is expected that the lighter chargino will be long-lived and produce heavily ionizing charged tracks in the detectors at the LHC. It must be noted though that the pair production cross section of a 1.3 TeV $\tilde{\chi}_1^\pm$ or the associated production cross section of $\tilde{\chi}_1^0\tilde{\chi}_1^\pm$ is of the order of 0.01–0.1 fb at the LHC with $\sqrt{s} = 14$ TeV. This means one can see signals of such a scenario only with a large integrated luminosity (~ 300 fb^{-1}). Nevertheless, from the above discussion it is clear that this model of leptogenesis is phenomenologically rich and possibly testable at the LHC. In [29], a class of high-scale nonuniversal scenario is suggested, where at the electroweak scale nearly degenerate neutralinos can be achieved.

It is noted in [8] that the nonholomorphic terms $N_1^{\prime e}$ can generate neutrino masses. The light neutrino mass is given as:

$$m_\nu = \frac{1}{256\pi^4}\frac{e^2}{\sin^2\theta_W}\mu^2\frac{m_\tau^2}{v^2} \times \xi^2 M_{\tilde{\chi}_1^0}\frac{M_{\tilde{\nu}}^2 - M_{\tilde{\chi}_1^0}^2 - M_{\tilde{\chi}_1^0}^2\ln(M_{\tilde{\nu}}^2/M_{\tilde{\chi}_1^0}^2)}{(M_{\tilde{\nu}}^2 - M_{\tilde{\chi}_1^0}^2)^2}, \tag{34}$$

when the slepton $\tilde{l}^+$ that mixes with h^+ is mainly $\tilde{\tau}^+$. In our scenario, we find that correct order of neutrino mass (0.049 eV) is achieved for the parameters in Set II [see Eq. (31)] with $M_{\tilde{\nu}} = 1381.2$ GeV. For the other set of parameters, the correct order of neutrino masses can be generated with a suitable choice of $M_{\tilde{\nu}}$. Thus a realistic scheme of radiative neutrino mass generation which originates from the same nonholomorphic terms can be accommodated along with other phenomenological aspects, discussed in earlier sections, in our scenario.

V. CONCLUSIONS

We discuss the possibility of resonant leptogenesis in an R-parity violating supersymmetric standard model with nonholomorphic supersymmetry breaking terms in the scalar potential. We work within a parameter space where the lighter neutralinos, $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$ are nearly degenerate in mass. In this framework we find out a consistent scenario where neutrino masses, low-scale leptogenesis, and interesting collider signatures exist simultaneously. As the masses of the $\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ are around the TeV scale, they can be produced at the LHC with small but finite cross sections. This makes our model phenomenologically rich and accessible at the LHC with signatures involving heavily ionizing charged tracks. Thus this low-scale resonant leptogenesis model is testable at the LHC which might give an indication of the presence of nonholomorphic couplings.

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