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Vacuum oscillations in the supersymmetric standard model

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We analyze the spectrum and mixing among neutrinos in the minimal supersymmetric standard model with explicit breaking of R parity. It is shown that (i) the mixing among neutrinos could be large and (ii) the nonzero neutrino mass is constrained to be $\leq 10^{-5}$ eV from arguments based on baryogenesis. Thus vacuum oscillations of neutrinos in this model may offer a solution of the solar neutrino problem. The allowed space of the supersymmetric parameters consistent with this solution is determined.

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I. INTRODUCTION

The neutrino masses [1] are known to solve some outstanding problems, notably, the solar neutrino, the atmospheric neutrino, and the dark matter problems. Theoretically, the presence of a nonzero neutrino mass provides a window into physics beyond the standard electroweak model. The generation of neutrino masses is possible either in the presence of neutral Higgs bosons transforming as an $SU(2)_L$ triplet and/or if there exist additional neutral fermions with which the conventional neutrinos could mix. The most popular example of the latter kind is provided by the seesaw mechanism [1] in which left-handed neutrinos obtain their masses through mixing with right-handed neutrinos. Another example is provided in supersymmetric theory [2] which automatically contains additional neutral fermions, namely, gauginos and Higgsinos. However, in this theory neutrinos cannot mix with the latter if the Lagrangian possesses a symmetry, called R parity (R) [3], which distinguishes between matter and supermatter. But breakdown of R parity can lead to mixing of neutrinos with gauginos and Higgsinos [4, 5] and hence to the mass of neutrinos. In fact, as long as the terms associated with R parity breaking are small, the natural seesaw mechanism is operative, with gauginos and Higgsinos playing the role of the "right-handed neutrino" of the conventional seesaw mechanism. The possibility of the generation of a nonzero neutrino mass in supersymmetry (SUSY) models with broken R parity is extensively discussed in the literature, both in the supersymmetric standard model (SSM) [5, 6] and in some of its extensions [7]. The neutrino masses have in fact been used to put constraints on the amount of admissible violation of R [5–8]. There exist independent constraints on the strength of R breaking. They arise by requiring [9] that the $(B-L)$ violation associated with R breaking should not erase the baryon asymmetry in the presence of the sphelaron-induced $(B+L)$ violation [10]. This can impose severe restrictions on the neutrino masses. In this paper we wish to systematically analyze the structure of neutrino masses and mixing in the SSM with explicitly broken R parity utilizing the restrictions imposed on R-parity-violating interactions.

The neutrino mass can offer a solution [1] to the solar neutrino problem if the relevant (mass)^2 difference Δ lies either around 10^{-5}–10^{-6} eV^2 or around 10^{-10} eV^2. The matter-induced resonant oscillations [11] deplete the neutrino flux in the former while vacuum or "just so" oscillations [12] are responsible for the depletion in the latter case. Moreover, the vacuum solution is feasible only if the mixing angle θ between oscillating neutrinos is large, typically $\sin^2(2\theta) \sim 0.75$–1 [1]. Such solutions therefore require extremely tiny neutrino masses which can arise, for example, from Planck scale physics [13]. We shall show that the constraints coming from baryogenesis in the SSM in fact restrict the neutrino mass to be as small as $\sim 10^{-5}$ eV, a value just required in order to solve the solar neutrino problem. The large mixing required for this purpose also follows naturally in the SSM with broken R parity, as we will see. While restrictions on R parity breaking coming from the baryogenesis are well known, their implications for obtaining the vacuum solution to the solar neutrino problem have not been stressed before. The purpose of the present note is to point out this possibility and at the same time determine the allowed region of parameters which realize it.

We discuss the SSM with R parity violation and summarize the constraints coming from baryogenesis in the next section. In Sec. III we discuss the structure of neutralino masses in the presence of R parity violation. An effective seesaw mechanism allows one to reduce the 7×7 mass matrix to a 3×3 effective neutrino mass matrix and

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makes it possible to discuss the mixing among neutrinos analytically. We determine in this section restrictions imposed on the conventional parameters of the SSM if one wants to solve the solar neutrino problem. The R violation causes the neutralinos to decay into Z^* and a neutrino. We determine this coupling and discuss its implications in Sec. IV. Conclusions are presented in the last section.

II. R-PARITY VIOLATING SSM

We shall confine ourselves to the SSM. This is characterized by the following superpotential in standard notation:

$$W_0 = \varepsilon_{ab}\left[h_{ij}\hat{L}_i^a\hat{H}_1^b\hat{E}_j^C + h'_{ij}\hat{Q}_i^a\hat{H}_1^b\hat{D}_j^C + h''_{ij}\hat{Q}_i^a\hat{H}_2^b\hat{U}_j^C + \mu\hat{H}_1^a\hat{H}_2^b\right]. \quad (1)$$

In addition to $SU(2)_L\otimes U(1)_Y$, this potential is also invariant under the discrete R parity under which quarks, leptons, Higgs bosons, and gauge bosons are even while their superpartners are odd. R can be broken in SSM explicitly by the terms

$$W_{R\!\!\!/} = \varepsilon_{ab}\left[\lambda_{ijk}\hat{L}_i^a\hat{L}_j^b\hat{E}_k^C + \lambda'_{ijk}\hat{L}_i^a\hat{Q}_j^b\hat{D}_k^C + \epsilon_i\hat{L}_i^a\hat{H}_2^b\right] + \lambda''_{ijk}\hat{U}_i^C\hat{D}_j^C\hat{D}_k^C\ . \quad (2)$$

The presence of both lepton- and baryon-number-violating terms in Eq. (2) leads to difficulties with the proton lifetime. From now on we therefore set the coupling λ''_{ijk} to zero. The full superpotential W is now the sum of W_0 and $W_{R\!\!\!/}$.

Even in the absence of R-parity-violating terms $W_{R\!\!\!/}$, R can be broken spontaneously if the sneutrino acquires a nonzero vacuum expectation value (VEV) $\tilde{\nu}_L$ [14]. This possibility is allowed in the SSM [15], but the resulting VEV of the sneutrino field is large, typically around the weak scale for a natural range of parameters. Moreover, one generates a Majoron which is strongly coupled to Z. Such a Majoron is in conflict with the invisible width of the Z. Hence one must abandon the idea of spontaneous R violation in the the SSM. Restrictions on parameters of SSM implied by this requirement were worked out in [16]. We shall assume these parameters to lie in the range determined in [16] and hence $\langle\tilde{\nu}_L\rangle$ would be assumed zero in the absence of R-breaking terms. But now if one introduces a small explicit R-breaking term $\varepsilon_{ab}\epsilon_i\hat{L}_i^a\hat{H}_2^b$ then the sneutrino VEV automatically gets generated [5, 18]. In this case, the VEV is related to the explicit rather than to the spontaneous violation of R parity and hence is not accompanied by a massless Majoron. Moreover, as long as the R-breaking parameters are small, the sneutrino VEV's also remain small and one avoids conflict with phenomenology which requires the sneutrino VEV to be small independent of the existence of the Majoron [17]. On the other hand, the R breaking induced by ϵ_i and the consequent sneutrino VEV can lead to interesting predictions for neutrino masses.

The scalar potential following from W and general soft supersymmetric breaking terms has the form [18]

$$\begin{aligned} V_{\text{Higgs}} = {} & \mu_1^2|\phi_1|^2 + \mu_2^2|\phi_2|^2 + \mu_{L_i}^2(\varphi_i^\dagger\varphi_i) + \frac{1}{2}\lambda_1\left[|\phi_1|^4 + |\phi_2|^4 + (\varphi_i^\dagger\varphi_i)^2 + 2|\phi_1|^2(\varphi_i^\dagger\varphi_i) - 2|\phi_2|^2(\varphi_i^\dagger\varphi_i)\right] \\ & + \lambda_2|\phi_1|^2|\phi_2|^2 - (\lambda_1+\lambda_2)|\phi_1^\dagger\phi_2|^2 + \left[\lambda_3(\phi_1^\dagger\phi_2) + \text{H.c.}\right] \\ & + \left[i\kappa_i(\phi_1^T\tau_2\varphi_i) + \text{H.c.}\right] + \left[i\kappa'_i(\phi_2^T\tau_2\varphi_i) + \text{H.c.}\right] + \cdots\ , \end{aligned} \quad (3)$$

where $\varphi_i \equiv L_i$, $\phi_2 \equiv H_2$, and $\phi_1 \equiv -i\tau_2 H_1^*$, and the ellipsis indicates terms of the potential not relevant for minimization. In deriving Eq. (3) we have assumed for simplicity that all ϵ_i are equal. Then the parameters of the Higgs potential are (we will neglect all possible CP-violating phases)

$$\begin{aligned} &\mu_1^2 = m_1^2 + |\mu|^2, \quad \mu_2^2 = m_2^2 + |\mu|^2 + \epsilon_i\epsilon_i, \\ &\mu_{L_i}^2 = m_{L_i}^2 + |\epsilon_i|^2 \\ &\lambda_1 = \frac{1}{4}(g^2 + g'^2), \quad \lambda_2 = \frac{1}{2}g^2 - \lambda_1, \\ &\lambda_3 = -m_{12}^2, \quad \kappa_i = \mu\epsilon_i. \end{aligned} \quad (4)$$

The parameters m_i^2, m_{12}^2, $m_{L_i}^2$, and κ'_i are soft breaking parameters. κ'_i gets related to ϵ_i at the Planck scale in the SSM. The g and g' are the gauge couplings.

Because of the presence of the κ_i and κ'_i terms in (3) the minimization invariably leads to nonzero VEV's $\omega_i \equiv \langle\tilde{\nu}_{iL}\rangle$ of the sneutrino. These are given by [18]

$$\omega_i = \frac{\kappa_i v_1 + \kappa'_i v_2}{\mu_{L_i}^2 + \frac{1}{2}\lambda_1\left(|v_1|^2 - |v_2|^2 + \sum_k|\omega_k|^2\right)}\ . \quad (5)$$

We are working here in the unconventional basis in which the ϵ term in W is not rotated away. Even if one chooses to utilize this freedom of rotation, the essential ingredients remain unchanged. In particular, the VEV for the sneutrino gets generated independent [5] of the basis one chooses. Note from Eq. (5) that ω_i vanish when the R-breaking terms κ_i, κ'_i, and ϵ_i are taken as zero. In this limit the model reduces to the minimal standard model and thus has two scalars and a massive pseudoscalar. The spectrum does not contain a Majoron since R is not spontaneously broken. The majoron of the SSM with spontaneously broken R symmetry [15, 16] in fact now has a mass [16] $\sim \frac{\kappa v_1 + \kappa' v_2}{\omega}$ which is of the order of the weak scale from Eq. (5).

As we will see in the next section, the parameters ϵ_i and ω_i determine the tree level neutrino masses and mixing. We therefore summarize the restrictions [9] on these parameters which follow from the baryogenesis [10]. Lepton number violation induced by ϵ_i, κ_i, or κ_i' could erase the existing baryon (or $B-L$) asymmetry if the $(B+L)$-violating sphaleron interactions are simultaneously in equilibrium with the lepton-number-violating interactions. The constraints on ϵ_i, κ_i, and κ_i' follow by demanding that the corresponding interaction be out of thermal equilibrium when the sphaleron interactions are in equilibrium, i.e., for $T \geq 100$ GeV. The rates for the L-violating interactions characterized by ϵ_i are typically given by $\Gamma_2 \sim \epsilon^2/T$. These interactions are out of thermal equilibrium for $T > 100$ GeV if $\Gamma_2 < 20HT^2/M_p$ for $T \sim T_C \sim 100$ GeV (H is the Hubble constant). This immediately implies [9]

$$\epsilon_i \leq 10^{-6} \text{ GeV}. \tag{6}$$

Likewise, requiring the rates for dimension-3 interactions characterized by κ_i, κ_i' to be less than the expansion rate H at $T \sim T_C$ one obtains

$$\kappa_i, \kappa_i' \leq 10^{-4} \text{ GeV}^2. \tag{7}$$

The constraints on κ_i, κ_i' can be translated into constraints on the VEV ω_i through Eq. (5). If one takes $v_1 \sim v_2 \sim \mu_{L_i} \sim 100$ GeV then Eq. (7) implies

$$\omega_i \leq 10^{-6} \text{ GeV}. \tag{8}$$

The exact limits on ω_i depend upon the model parameters. But we shall regard the limit on ω_i as given in Eq. (8) as indicative of the typical limit and work out the consequences of Eqs. (6) and (8) in the next section.

The restrictions displayed in Eqs. (6) and (8) are generic constraints which hold in a general situation. If some of the ϵ_i and κ_i' are zero then the theory automatically possesses a global lepton number symmetry corresponding to the ith lepton number. The presence of such a global symmetry could prevent [19] the erasure of the baryon asymmetry. The other nonzero ϵ_i would not be constrained in this case. We shall disregard this possibility and assume no global symmetry L_i to be exact.

III. NEUTRINO MASSES IN THE SSM

Neutrino masses arise in the SSM of the last section through three different sources. First, the nonzero ϵ_i directly induce mixing of the neutrino with Higgsinos. Secondly, the VEV ω_i induced by the presence of ϵ_i gives rise to mixing between neutrinos and gauginos. These two sources contribute at the tree level. But since the lepton number is violated, one could radiatively generate the direct Majorana mass term among neutrinos [5, 6]. Their strength is also controlled by the basic parameters ϵ_i and other R-breaking parameters in Eq. (2). We shall assume that the tree level contribution dominates over the radiatively generated masses. Since the baryogenesis constrains the tree level mass very significantly $\leq 10^{-5}$ eV, it is reasonable to neglect the radiative contributions and concentrate on the tree level masses.

In addition to the three neutrinos, the SSM contains two gauginos ($\tilde{B}$,$\tilde{W}_3$) and two Higgsinos ($\tilde{H}_1$,$\tilde{H}_2$). The neutralino mass matrix has the following form [2, 4, 16] in the basis $\chi'^T = (\nu, -i\tilde{B}, -i\tilde{W}_3, \tilde{H}_1, \tilde{H}_2)$:

$$\mathcal{M}_0 = \begin{pmatrix} 0 & m \\ m^T & M_4 \end{pmatrix} , \tag{9}$$

where m is a 3×4 matrix given by

$$m = \begin{pmatrix} -\frac{g'}{2}\omega_1 & \frac{g}{2}\omega_1 & 0 & -\epsilon_1 \\ -\frac{g'}{2}\omega_2 & \frac{g}{2}\omega_2 & 0 & -\epsilon_2 \\ -\frac{g'}{2}\omega_3 & \frac{g}{2}\omega_3 & 0 & -\epsilon_3 \end{pmatrix} \tag{10}$$

and M_4 is the usual neutralino mass matrix [2] describing neutralino mixing in the absence of R parity breaking:

$$M_4 = \begin{pmatrix} cM & 0 & -\frac{1}{2}g'v_1 & \frac{1}{2}g'v_2 \\ 0 & M & \frac{1}{2}gv_1 & -\frac{1}{2}gv_2 \\ -\frac{1}{2}g'v_1 & \frac{1}{2}gv_1 & 0 & -\mu \\ \frac{1}{2}g'v_2 & -\frac{1}{2}gv_2 & -\mu & 0 \end{pmatrix} . \tag{11}$$

M is the common gaugino mass parameter and $c = 5\alpha_1/\alpha_2 = 0.5$ [16]. Since the parameters (ϵ_i, ω_i) entering m are expected to be much smaller than the ones appearing in M_4, the neutralino mass matrix $\mathcal{M}_0$ has a seesaw structure. Hence the the neutrino masses and mixing are derived from an effective mass matrix of the form

$$\begin{aligned} m_{\text{eff}} &= -m\, M_4^{-1}\, m^T \\ &= \frac{(cg^2+g'^2)}{D} \begin{pmatrix} A_1^2 & A_1A_2 & A_1A_3 \\ A_1A_2 & A_2^2 & A_2A_3 \\ A_1A_3 & A_2A_3 & A_3^2 \end{pmatrix} , \end{aligned} \tag{12}$$

where we have defined

$$\mathbf{A} \equiv \mu\boldsymbol{\omega} - v_1\boldsymbol{\epsilon} \tag{13}$$

and

$$D \equiv 4\frac{\det M_4}{M} = 2\mu\left[-2cM\mu + v_1v_2\left(cg^2+g'^2\right)\right] . \tag{14}$$

Although the 7×7 neutralino mass matrix is quite complex, the neutrino masses can be approximately described by a simple structure displayed in Eq. (12). m_{eff} can be diagonalized by an orthogonal matrix O:

$$O\, m_{\text{eff}}\, O^T = \text{diag}(0,\ 0,\ m_\nu) , \tag{15}$$

with the only nonzero neutrino mass given by

$$m_\nu = \text{tr}(m_{\text{eff}}) = \frac{(cg^2+g'^2)}{D}|\mathbf{A}|^2 . \tag{16}$$

The matrix O^T can be parametrized by

$$O^T = \begin{pmatrix} \cos\theta_{13} & 0 & -\sin\theta_{13} \\ \sin\theta_{23}\sin\theta_{13} & \cos\theta_{23} & \sin\theta_{23}\cos\theta_{13} \\ \sin\theta_{13} & \sin\theta_{23} & \cos\theta_{13}\cos\theta_{23} \end{pmatrix} \tag{17}$$

with the mixing angles given by

$$\tan\theta_{13} = -\frac{A_1}{\sqrt{A_2^2+A_3^2}}, \qquad \tan\theta_{23} = \frac{A_2}{A_3}\,. \tag{18}$$

We note the following.

Two of the eigenvalues of $m_{\rm eff}$ are zero. This is not an artifact of the seesaw approximation, but follows in a more general situation with the full 7×7 matrix $\mathcal{M}_0$. It is easy to see that the following ψ_0 represents two eigenvectors of $\mathcal{M}_0$ with zero eigenvalues:

$$\begin{aligned}\psi_0^T = (&(\boldsymbol{\epsilon}\wedge\boldsymbol{\omega})_1 x_3 - A_2 x_6,\\ &(\boldsymbol{\epsilon}\wedge\boldsymbol{\omega})_2 x_3 + A_1 x_6,\ (\boldsymbol{\epsilon}\wedge\boldsymbol{\omega})_3 x_3,\ 0,\ 0,\ x_6,\ 0)\end{aligned} \tag{19}$$

for arbitrary x_3 and x_6. This feature of the neutralino masses is a direct consequence of the restricted structure of $\mathcal{M}_0$ implied by the particle and charge assignments in the SSM.

The nonzero eigenvalue is typically given by

$$m_\nu \sim \frac{\epsilon^2}{M}\,. \tag{20}$$

For $\epsilon \sim 10^{-6}$ GeV and $M \sim 100$ GeV one has $m_\nu \sim 10^{-5}$ eV. This is in the right range for a solution of the solar neutrino problem through vacuum oscillations [20].

If ϵ_i do not display any hierarchy then both mixing angles are automatically large. In fact for $\epsilon_1 \sim \epsilon_2 \sim \epsilon_3$ and $\omega_1 \sim \omega_2 \sim \omega_3$ we have

$$\tan\theta_{23} \sim 1, \qquad \tan\theta_{13} \sim -\frac{1}{\sqrt{2}}\,. \tag{21}$$

Thus if all ϵ_i and ω_i are flavor independent and near their limit coming from baryogenesis then one naturally generates the (mass)^2 difference and the mixing angles required for the vacuum solution to the solar neutrino problem. The details depend upon other parameters as well and we will now present the quantitative analysis.

Given the mixing matrix O and the mass m_ν, the survival probability for the solar ν_e after time t is given by

$$P_{\nu_e\nu_e} = 1 - \sin^2\theta_{13}\sin^2\frac{\Delta t}{2p}, \tag{22}$$

where $\Delta \equiv m_\nu^2$. This displays a simple two-generation-like structure due to the fact that there is only one nontrivial (mass)^2 difference. The restrictions on the parameters θ_{13} and Δ have been worked out in detail [12] combining observations of all the four solar neutrino detectors. The allowed ranges of these parameters are given by

$$\begin{aligned}&\Delta \simeq (0.5\text{–}1.0)\times 10^{-10}\ \mathrm{eV}^2,\\ &\sin^2(2\theta_{13}) \simeq 0.75\text{–}1.0\,.\end{aligned} \tag{23}$$

Note that the expected value of $\tan\theta_{13} \sim -1/\sqrt{2}$ when ϵ_i and ω_i are flavor independent falls within the allowed range. The exact value of Δ depends upon the parameters of the SSM in addition to ϵ_i and ω_i. These SSM parameters are tightly constrained by various observations at the CERN e^+e^- collider LEP and the $p\bar{p}$ collider. We shall use these constraints and show that values of Δ in the above range are possible for ϵ_i and ω_i near the limit from baryogenesis.

Various observables in the SSM can be expressed in terms of the three basic parameters μ, M, and $\tan\beta$. Observations at LEP and $p\bar{p}$ colliders have been used [19] to restrict these parameters. In the following we fix the ω_i, ϵ_i near their limit coming from baryogenesis and then show that it is possible to obtain the Δ in the band required for the vacuum oscillation solution to the solar neutrino problem, for a range of values (μ, M) allowed by the other observables. Specifically we choose $\omega_i = \epsilon_i \equiv \frac{\omega}{\sqrt{3}}$ independent of the flavor. As already remarked, for these values, the mixing angle $\tan\theta_{13} = -1/\sqrt{2}$ and is in the allowed band.

The constraints on μ, M, and $\tan\beta$ coming from the nonobservation of the decay $Z \to \chi^+\chi^-$ ($\chi^\pm$ being the chargino) is found [21] to be very restrictive among various restrictions that are possible on the SSM parameters from the LEP and $p\bar{p}$ collider experiments [22]. As an illustration, we reproduce in Fig. 1 (solid line) the allowed values of μ and M for $\tan\beta = 4$, obtained by requiring that the lighter of the charginos be heavier than 45 GeV. To be specific, we have taken $\omega^2 = 2\times 10^{-12}$ GeV^2 and plotted the curves in the μ, M plane corresponding to $\Delta = 10^{-10}$ eV^2 and $\Delta = 0.5\times 10^{-10}$ eV^2 (broken line). It is seen that there exists a sizable region in the μ, M plane which is allowed by various observations and which offers a solution to the solar neutrino problem. The allowed region is dependent on the chosen values of ϵ and ω. This dependence is displayed in Fig. 2, which shows the variation of Δ with the basic parameters ϵ_i and ω_i. We once again assume $\epsilon_i \equiv \frac{\epsilon}{\sqrt{3}}$ and $\omega_i \equiv \frac{\omega}{\sqrt{3}}$ independent of i and plot the region in the ϵ and $\delta \equiv \epsilon/\omega$ plane, for which Δ lies between 10^{-10} eV^2 and 5×10^{-11} eV^2. $\tan\beta$ is chosen to be 4 and the curves are shown for two typical values of the pair (μ, M). This figure highlights the fact

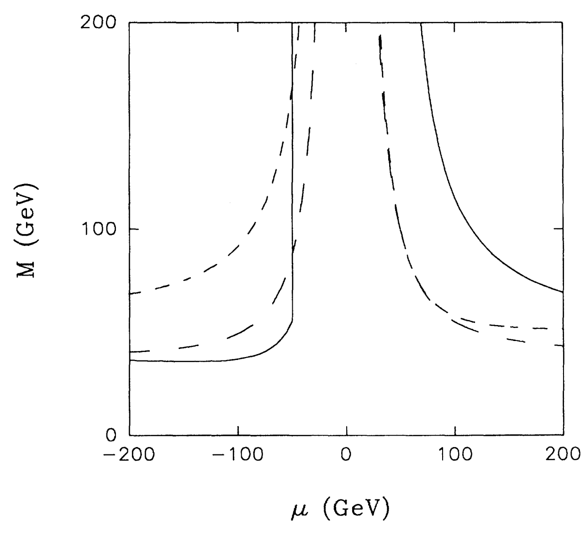

FIG. 1. The allowed region in the μ-M plane corresponding to the chargino mass > 45 GeV (solid line) and $5\times 10^{-11} < \Delta < 10^{-10}$ eV^2 (broken line). $\tan\beta$ is chosen to be 4 and $\omega^2 = 2\times 10^{-12}$ eV^2. The region above each curve is allowed.

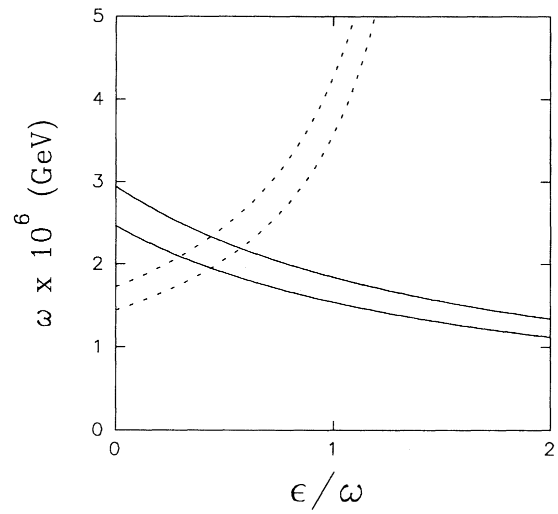

FIG. 2. Band of values of ω and $\frac{\epsilon}{\omega}$ allowed by $5\times10^{-11} < \Delta < 10^{-10}$ eV^2. The solid and the dotted curves are for $(\mu, M) = (100, 100)$ GeV, and (−100,100) GeV, respectively. $\tan\beta$ is chosen to be 4. The region above each curve is allowed.

that for $\epsilon \sim \omega$ and ϵ in the range allowed by the baryogenesis constraint, one could get Δ in the range required for solving the solar neutrino problem.

The bound on the neutrino mass ($m_\nu \leq 10^{-5}$ eV) following here from the baryogenesis constraint is to be contrasted with an analogous bound [23] on the (Majorana) masses of neutrinos in the generic seesaw type model. The lepton number violation appears in this model through a large Majorana mass for the neutrino. It generates the left-handed neutrino mass through a dimension-5 operator. Baryogenesis constraints on this operator lead to a typical mass $m_\nu < 50$ eV. In contrast, dimension-2 and -3 terms are responsible for the neutrino masses and the baryogenesis constraint on them translates into a much stronger limit, $m_\nu < 10^{-5}$ eV, in the context of the SSM considered here.

IV. NEUTRALINO DECAY

The lightest R-odd particle is not allowed to decay in the SSM with R symmetry. Because of its stability and neutrality, the lightest supersymmetric particle (LSP) is considered an ideal candidate for the dark matter of the universe [24]. Usually, one expects a combination of neutralinos to be the LSP [2]. The presence of even a tiny amount of R violation such as the one considered here can make the LSP unstable on a cosmological scale. The couplings which make the LSP unstable have been considered in the literature [4, 8, 16]. The treatment of the neutralino mass matrix based on the seesaw approximation makes it possible to write down the couplings of neutralino to neutrino analytically without neglecting the intergenerational mixing. We give these couplings below for completeness and discuss their consequences.

R parity violation causes mixing between two neutral particles (neutrinos and gauginos) transforming differently under the SU(2) × U(1) group. As is well known, the coupling of Z to fermions no longer remains flavor diagonal in this case. Neutralinos couple to Z through the following equation in the weak basis χ':

$$-\mathcal{L}_Z = \frac{g}{2\cos\theta_W}\,[\overline{\chi}'\,\overline{\sigma}_\mu T_3\,\chi']\;Z^\mu\;, \tag{24}$$

where $T_3 \equiv (1,1,1,0,0,-1,1)^T$. Let U be the matrix which diagonalizes the $\mathcal{M}_0$ of Eq. (1):

$$U\,\mathcal{M}_0\,U^T = \mathrm{diag}(0,0,m_\nu,m_4,m_5,m_6,m_7)\;. \tag{25}$$

m_α, $\alpha = 4,5,6,7$, are neutralino masses. The form of U is well known in the seesaw approximation:

$$U = \begin{pmatrix} O_\nu(1-\frac{1}{2}\rho\rho^T) & -O_\nu\rho \\ O_\lambda\rho^T & (1-O_\lambda\frac{1}{2}\rho^T\rho) \end{pmatrix}\;. \tag{26}$$

Here $\rho \equiv m\,M_4^{-1}$. O_ν (O_λ) diagonalize m_{eff} (M_4) of Eqs. (11) and (12).

The flavor nondiagonal coupling of neutralinos to Z follows from Eqs. (24)–(26) after some algebra:

$$-\mathcal{L}_{\nu\chi} = \frac{g}{4\cos\theta_W}\,[F^*_{i\alpha}\overline{\nu}_{iL}\gamma_\mu\chi_{\alpha L} - F_{i\alpha}\overline{\nu}_{iR}\gamma_\mu\chi_{\alpha R}]\,Z^\mu\;, \tag{27}$$

where

$$\begin{aligned} F_{i\alpha} = {} & -\frac{-2\mu|\mathbf{A}|}{D}[g'(O_\lambda)_{\alpha4} - cg(O_\lambda)_{\alpha5}]\delta_{i3} \\ & -\left[\frac{v_2}{\mu}(g'^2+cg^2)(O_\nu)_{ij}A'_j - 4c\epsilon_j M\right](O_\lambda)_{\alpha6} \end{aligned} \tag{28}$$

with $A'_i = A_i + 2\epsilon_i v_1$ and $|\mathbf{A}|$ and D defined before.

The elements $O_{\alpha\beta}$ appearing in the equation above depend upon the composition of the LSP in terms of neutralinos. They are given as in the MSSM [2]. But it follows from the structure of the above equation that couplings of the LSP to $Z\nu$ are not suppressed by the mixing factor coming from the neutralino mixing [25]. Thus, irrespective of its composition, the LSP will decay into a real (or virtual) Z and the neutrino. The typical strength is given from Eq. (28) by $\frac{g\epsilon}{\mu}$. For $\epsilon \sim 10^{-6}$ GeV, this strength is great enough to make the LSP decay on the cosmological scale but is not large enough to cause its decay and hence any signature in the laboratory. The lifetime following from Eqs. (27) and (28) is typically given by

$$\begin{aligned} \tau & \sim \tau_\mu\left(\frac{g\epsilon}{\mu}\right)^{-2}\left(\frac{M_{\text{LSP}}}{m_\mu}\right)^{-5} \\ & \sim 1.6\times10^{-4}\ \text{sec}\ \left(\frac{M_{\text{LSP}}}{50\ \text{GeV}}\right)^{-5} \\ & \quad\times\left(\frac{100\ \text{GeV}}{\mu}\right)^{-2}\left(\frac{\epsilon}{10^{-6}\ \text{GeV}}\right)^{-2}\;, \end{aligned} \tag{29}$$

where τ_μ is the muon lifetime. It is seen from the above that the typical LSP with 50 GeV mass will not be able to decay inside the detector but it will be too short lived to have any cosmological signature. The LSP in this case would contribute to the invisible Z width but this contribution is suppressed by the factor $(\frac{\epsilon}{\mu})^2$ compared to other fermion contributions and thus is practically negligible.

V. CONCLUSIONS

We have discussed in this paper the possibility of obtaining a solution to the solar neutrino problem through vacuum oscillations in the SSM. This requires an extremely tiny (mass)^2 difference $\Delta \sim 10^{-10}$ eV^2. As we have discussed, the limit on the neutrino masses coming from baryogenesis implies $\Delta \leq 10^{-10}$ eV^2. Hence if R-breaking parameters are near this limit then the SSM offers a vacuum solution to the solar problem. This becomes more interesting due to the fact that the relevent mixing angle predicted in this model can be large as required for the vacuum solution if the R-parity-breaking parameters do not display any hierarchy in flavor space.

A similar analysis of the Mikheyev-Smirnov-Wolfenstein (MSW) solution to the solar neutrino problem in supersymmetric theory has been earlier carried out in an extension of the SSM involving right-handed neutrinos and other gauged singlet superfields [7]. This analysis concentrated on the spontaneous rather than the explicit violation of R. In such a case the baryogenesis may not restrict the amount of R violation. In contrast, it is not possible to break R spontaneously in the minimal case considered here, and one should then consider baryogenesis constraints. In spite of its strength, these constraints do allow interesting physical effect namely a solution to the solar neutrino problem as we we have argued here.

It is hard to explain why the R-parity-breaking parameters ϵ_i are as small as is required from the baryogenesis limit. But if they do have such values then they may be responsible for causing vacuum oscillations of the solar neutrinos.

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