

Sneutrino mixing

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Abstract. In supersymmetric models with nonvanishing Majorana neutrino masses, the sneutrino and antisneutrino mix. The conditions under which this mixing is experimentally observable are studied, and mass-splitting of the sneutrino mass eigenstates and sneutrino oscillation phenomena are analyzed.

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1. Introduction

In the minimal Standard Model, as well as in the minimal supersymmetric extension of it (MSSM) neutrinos are exactly massless [1]. The present direct laboratory upper bounds on their masses are [2]

$$m_{\nu_e} \lesssim 10 \text{ eV}, \quad m_{\nu_\mu} \leq 0.17 \text{ MeV}, \quad m_{\nu_\tau} \leq 18 \text{ MeV}, \quad (1)$$

and cosmological constraints require stable neutrinos to be lighter than about 100 eV [3]. However, the solar [4] and atmospheric [5] neutrino puzzles, the LSND results [6] and models of mixed (hot and cold) dark matter [7] suggest that neutrinos may have a small mass, $m_\nu \sim 10^{-5} - 10 \text{ eV}$. The most attractive way to get such masses is to introduce Majorana neutrino mass terms that violate lepton number (L) by two units.

In this talk, based upon Ref. [8, 9] (see also [10] for an independent study), we consider a supersymmetric Standard Model with Majorana neutrino masses. In such models, lepton number violation can generate interesting phenomena in the slepton sector. In addition to generating small neutrino masses, the $\Delta L = 2$ operators introduce a mass splitting and mixing in the sneutrino-antisneutrino system. The sneutrino and antisneutrino will then no longer be mass eigenstates.

This is analogous to the neutral meson systems[11]. For example, in the B^0 system the effect of a small $\Delta B = 2$ perturbation to the leading $\Delta B = 0$ mass term results in a mass splitting between the heavy and light B^0 , which are no longer pure B^0 and \bar{B}^0 states. The very small mass splitting, $\Delta m_B/m_B = 6 \times 10^{-14}$ [2], can be measured by observing flavor oscillations. The flavor is tagged in B -decays by the final state lepton charge. Since $x_d \equiv \Delta m_B/\Gamma_B \approx 0.7$ [2], there is time for the flavor to oscillate before the meson decays. When B mesons are produced in pairs (for example in e^+e^- collider operating at the $\Upsilon(4S)$ resonance) the same sign dilepton signal indicates that only one of the B oscillated. This time-integrated same sign dilepton sample is used to determine the tiny mass splitting.

The sneutrino system can exhibit similar behavior. The lepton number is tagged in sneutrino decay using the charge of the outgoing lepton. If the sneutrino has time to mix before it decays, namely if

$$x_{\tilde{\nu}} \equiv \frac{\Delta m_{\tilde{\nu}}}{\Gamma_{\tilde{\nu}}} \gtrsim 1, \quad (2)$$

and if the branching ratio of the sneutrino decay into a charged lepton is significant, then we can directly measure a non-zero sneutrino mass splitting via the same sign dilepton signal. When the sneutrinos are pair produced, *e.g.* in e^+e^- collisions, the two leptons from the sneutrino decays are used. When the sneutrino is produced together with a charged lepton, *e.g.* in

hadron collider via cascade decays, the lepton from the sneutrino decay and the associated produced lepton are used. In both cases a measurable same sign dilepton signal is expected.

The neutrino mass and the sneutrino mass splitting are both consequences of the small breaking of lepton number. Therefore, they are expected to be related. Thus, we can use upper bounds (or indications) of neutrino masses to set bounds on the sneutrino mass splitting. We will consider the consequences of two cases: (i) ν_τ with a mass near its present laboratory upper limit, $m_\nu \sim 10$ MeV; and (ii) light neutrinos of mass less than 100 eV.

In order to derive specific results, one must specify a model for the lepton number violation. In the following we concentrate on two models of neutrino masses: the see-saw mechanism and R-parity violation. We compute the sneutrino mass splitting in each model and its relation to the neutrino mass. We then briefly discuss the consequences for sneutrino phenomenology in colliders

2. The Supersymmetric See-Saw Model

Consider an extension of the MSSM where one adds a right-handed neutrino superfield, \hat{N} , with a bare mass $M \gg m_Z$. We consider a one generation model (*i.e.*, we ignore lepton flavor mixing) and assume CP conservation. We employ the most general R-parity conserving renormalizable superpotential and attendant soft-supersymmetry breaking terms. For this work, the relevant terms in the superpotential are (following the notation of Ref. [12])

$$W = \epsilon_{ij} \left[\lambda \hat{H}_2^i \hat{L}^j \hat{N} - \mu \hat{H}_1^i \hat{H}_2^j \right] + \frac{1}{2} M \hat{N} \hat{N}. \quad (3)$$

The D -terms are the same as in the MSSM. The relevant terms in the soft-supersymmetry-breaking scalar potential are:

$$V_{\text{soft}} = m_{\tilde{L}}^2 \tilde{\nu}^* \tilde{\nu} + m_{\tilde{N}}^2 \tilde{N}^* \tilde{N} + (\lambda A_\nu H_2^2 \tilde{\nu} \tilde{N}^* + M B_N \tilde{N} \tilde{N} + \text{h.c.}). \quad (4)$$

When the neutral Higgs field vacuum expectation values are generated [$\langle H_i^0 \rangle = v_i/\sqrt{2}$, with $\tan \beta \equiv v_2/v_1$ and $v_1^2 + v_2^2 = v^2 = (246 \text{ GeV})^2$], one finds that the light neutrino mass is given by the usual one generation see-saw result

$$m_\nu = \frac{m_D^2}{M}, \quad (5)$$

where $m_D \equiv \lambda v_2$ and we drop terms higher order in m_D/M .

The sneutrino masses are obtained by diagonalizing a 4×4 squared-mass matrix. Here, it is convenient to define: $\tilde{\nu} = (\tilde{\nu}_1 + i\tilde{\nu}_2)/\sqrt{2}$ and

$\tilde{N} = (\tilde{N}_1 + i\tilde{N}_2)/\sqrt{2}$. Then, the sneutrino squared-mass matrix separates into CP-even and CP-odd blocks:

$$-\mathcal{L}_{\text{mass}} = \frac{1}{2} (\phi_1 \quad \phi_2) \begin{pmatrix} \mathcal{M}_+^2 & 0 \\ 0 & \mathcal{M}_-^2 \end{pmatrix} \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix}, \quad (6)$$

where $\phi_i \equiv (\tilde{\nu}_i \quad \tilde{N}_i)$ and

$$\mathcal{M}_\pm^2 = \begin{pmatrix} m_{\tilde{L}}^2 + \frac{1}{2}m_Z^2 \cos 2\beta + m_D^2 & m_D[A_\nu - \mu \cot \beta \pm M] \\ m_D[A_\nu - \mu \cot \beta \pm M] & M^2 + m_D^2 + m_{\tilde{N}}^2 \pm 2B_N M \end{pmatrix}. \quad (7)$$

In the following derivation we assume that M is the largest mass parameter. Then, to first order in $1/M$, the two light sneutrino eigenstates are $\tilde{\nu}_1$ and $\tilde{\nu}_2$, with corresponding squared masses:

$$m_{\tilde{\nu}_{1,2}}^2 = m_{\tilde{L}}^2 + \frac{1}{2}m_Z^2 \cos 2\beta \mp \frac{1}{2}\Delta m_{\tilde{\nu}}^2, \quad (8)$$

where the squared mass difference $\Delta m_{\tilde{\nu}}^2 \equiv m_{\tilde{\nu}_2}^2 - m_{\tilde{\nu}_1}^2$ is of order $1/M$. Thus, in the large M limit, we recover the two degenerate sneutrino states of the MSSM, usually chosen to be $\tilde{\nu}$ and $\bar{\tilde{\nu}}$. For finite M , these two states mix with a 45° mixing angle, since the two light sneutrino mass eigenstates must also be eigenstates of CP. The sneutrino mass splitting is easily computed using $\Delta m_{\tilde{\nu}}^2 = 2m_{\tilde{\nu}}\Delta m_{\tilde{\nu}}$, where $m_{\tilde{\nu}} \equiv \frac{1}{2}(m_{\tilde{\nu}_1} + m_{\tilde{\nu}_2})$ is the average of the light sneutrino masses. We find that the ratio of the light sneutrino mass difference relative to the light *neutrino* mass [eq. (5)] is given by (to leading order in $1/M$)

$$r_\nu \equiv \frac{\Delta m_{\tilde{\nu}}}{m_\nu} \simeq \frac{2(A_\nu - \mu \cot \beta - B_N)}{m_{\tilde{\nu}}}. \quad (9)$$

The magnitude of r_ν depends on various supersymmetric parameters. Naturalness constrains supersymmetric mass parameters associated with particles with non-trivial electroweak quantum numbers to be roughly of order m_Z [13]. Thus, we assume that μ , A_ν , and $m_{\tilde{L}}$ are all of order the electroweak scale. The parameters M , $m_{\tilde{N}}$, and B_N are fundamentally different since they are associated with the $SU(2) \times U(1)$ singlet superfield \tilde{N} . In particular, $M \gg m_Z$, since this drives the see-saw mechanism. Since M is a supersymmetry-conserving parameter, the see-saw hierarchy is technically natural. The parameters $m_{\tilde{N}}$ and B_N are soft-supersymmetry-breaking parameters; their order of magnitude is less clear. Since \tilde{N} is an electroweak gauge group singlet superfield, supersymmetry-breaking terms associated with it need not be directly tied to the scale of electroweak symmetry breaking. Thus, it is possible that $m_{\tilde{N}}$ and B_N are much larger than m_Z . Since B_N enters directly into the formula for the light sneutrino mass splitting [eq. (9)], its value is critical for sneutrino phenomenology. If $B_N \sim \mathcal{O}(m_Z)$, then $r_\nu \sim \mathcal{O}(1)$, which implies that the sneutrino mass splitting is of order the *neutrino* mass. However, if $B_N \gg m_Z$, then the sneutrino mass splitting is significantly enhanced.

3. R-parity violation

Consider an extension of the MSSM where R-parity is not imposed (in this case, we do not add right handed neutrinos). Again, we ignore lepton flavor mixing and assume CP conservation. In this model one neutrino mass arises at tree level from neutrino mixing with the neutralinos via sneutrino VEVs or quadratic terms (“ μ -terms”) in the superpotential [14]. The sneutrino splitting arises from sneutrino mixing with the Higgs fields [9]. The other neutrino masses and sneutrino splittings arise at one loop [15] and are not considered here.

In models without R-parity there is a priori nothing to distinguish the lepton-doublet supermultiplets \hat{L}_i from the down-Higgs supermultiplet \hat{H}_d , as both transform as $(2)_{-1/2}$ under $SU(2)_L \times U(1)_Y$. With one generation there are two $Y = -1/2$ doublets which we denote by \hat{H}_1 and \hat{H}_3 . Then, in the superpotential the single μ -term of the MSSM is now extended to a vector

$$W = -\epsilon_{ij}\mu_\alpha \hat{H}_\alpha^i \hat{H}_2^j, \quad (10)$$

where $\alpha = 1, 3$ here and it what follows. The trilinear terms in the superpotential are irrelevant here. The D -terms are the same as in the MSSM. The single SUSY breaking B term of the MSSM is also extended to a vector, and SUSY breaking scalar masses are extended into a matrix

$$V_{\text{soft}} = m_{\alpha\beta}^2 H_\alpha H_\beta^* + m_2^2 |H_2|^2 + (B_\alpha H_\alpha H_2 + \text{h.c.}). \quad (11)$$

where H_1, H_2 and H_3 are the neutral components of the scalar fields. Finally, the single down type Higgs vev, v_1 , is also extended to a vector, v_α .

To get the neutrino mass we consider the neutralino mass matrix [14]. For simplicity we consider only one generation. The full 5×5 tree-level neutralino mass matrix with rows and columns corresponding to $\{\tilde{B}, \tilde{W}_3, \tilde{H}_2^0, \tilde{H}_1^0, \tilde{H}_3^0\}$ is

$$M^{(n)} = \begin{pmatrix} M_1 & 0 & m_Z s \frac{v_2}{v} & -m_Z s \frac{v_1}{v} & -m_Z s \frac{v_3}{v} \\ 0 & M_2 & -m_Z c \frac{v_2}{v} & m_Z c \frac{v_1}{v} & m_Z c \frac{v_3}{v} \\ m_Z s \frac{v_2}{v} & -m_Z c \frac{v_2}{v} & 0 & \mu_1 & \mu_3 \\ -m_Z s \frac{v_1}{v} & m_Z c \frac{v_1}{v} & \mu_1 & 0 & 0 \\ -m_Z s \frac{v_3}{v} & m_Z c \frac{v_3}{v} & \mu_3 & 0 & 0 \end{pmatrix}, \quad (12)$$

where $c = \cos \theta_W$, $s = \sin \theta_W$ and $v^2 = v_1^2 + v_2^2 + v_3^2$. We define

$$\mu \equiv \left(\sum_\alpha \mu_\alpha^2 \right)^{1/2}, \quad v_d \equiv \left(\sum_\alpha v_\alpha^2 \right)^{1/2}, \quad \cos \xi \equiv \frac{\sum_\alpha v_\alpha \mu_\alpha}{v_d \mu}. \quad (13)$$

Note that ξ measures the alignment of v_α and μ_α . The product of the masses is then

$$\det M^n = (m_Z^2 \mu^2 m_{\tilde{\gamma}}) \cos^2 \beta \sin^2 \xi, \quad (14)$$

where $\tan \beta = v_2/v_d$ and $m_{\tilde{\gamma}} = \cos^2 \theta_W M_1 + \sin^2 \theta_W M_2$. In the MSSM with R-parity [12], where the neutrino would be massless, the product of the four non-vanishing masses is

$$\det M_0^{(n)} = \mu (-M_1 M_2 \mu + \sin 2\beta m_Z^2 m_{\tilde{\gamma}}). \quad (15)$$

To first order in the neutrino mass, the neutralino masses are unchanged by the R-parity violating terms. Thus, we get [16]

$$m_\nu = \frac{\det M^{(n)}}{\det M_0^{(n)}} = \rho_\nu m_Z \cos^2 \beta \sin^2 \xi, \quad (16)$$

with

$$\rho_\nu = \frac{m_Z \mu m_{\tilde{\gamma}}}{(-M_1 M_2 \mu + \sin 2\beta m_Z^2 m_{\tilde{\gamma}})}. \quad (17)$$

We find $\rho_\nu \sim 1$ for $\mu \sim M_1 \sim M_2 \sim m_Z$.

The sneutrino mass splitting is a result of the difference in sneutrino mixing with the CP even and CP odd Higgs fields. The mass-squared matrices are given by

$$M_{\text{odd}}^2 = \frac{1}{v^2} \begin{pmatrix} a_{12}^2 v_2^2 + a_{13}^2 v_3^2 & a_{12}^2 v_1 v_2 & -a_{13}^2 v_1 v_3 \\ a_{12}^2 v_1 v_2 & a_{12}^2 v_1^2 + a_{23}^2 v_3^2 & a_{23}^2 v_2 v_3 \\ -a_{13}^2 v_1 v_3 & a_{23}^2 v_2 v_3 & a_{13}^2 v_1^2 + a_{23}^2 v_2^2 \end{pmatrix}, \quad (18)$$

$$M_{\text{even}}^2 = \frac{1}{v^2} \times \begin{pmatrix} m_Z^2 v_1^2 + a_{12}^2 v_2^2 + a_{13}^2 v_3^2 & -(m_Z^2 + a_{12}^2) v_1 v_2 & -(m_Z^2 + a_{13}^2) v_1 v_3 \\ -(m_Z^2 + a_{12}^2) v_1 v_2 & a_{12}^2 v_1^2 + m_Z^2 v_2^2 + a_{23}^2 v_3^2 & -(m_Z^2 + a_{23}^2) v_2 v_3 \\ -(m_Z^2 + a_{13}^2) v_1 v_3 & -(m_Z^2 + a_{23}^2) v_2 v_3 & a_{13}^2 v_1^2 + a_{23}^2 v_2^2 + m_Z^2 v_3^2 \end{pmatrix}$$

where

$$a_{12}^2 \equiv \frac{B_1 v^2}{v_1 v_2}, \quad a_{23}^2 \equiv \frac{B_3 v^2}{v_2 v_3}, \quad a_{13}^2 \equiv \frac{m_{13}^2 v^2}{v_1 v_3}. \quad (19)$$

Note that M_{odd}^2 includes the massless Goldstone boson and two massive CP-odd scalars. We work in the basis where $\mu_3 = 0$. Then, $v_3 \rightarrow 0$ only when both $m_{13}^2 \rightarrow 0$ and $B_3 \rightarrow 0$ in such a way that a_{ij} is finite [14]. Thus, the only small parameter is v_3 . We use $v_3 = v \cos \beta \sin \xi$ and find that to lowest order in $\sin \xi$

$$\Delta m_{\tilde{\nu}} = \rho_{\tilde{\nu}} m_Z \cos^2 \beta \sin^2 \xi, \quad (20)$$

where $\rho_{\tilde{\nu}} \sim \rho_\nu$ is given in [9]. In particular, we find

$$r_\nu \equiv \frac{\Delta m_{\tilde{\nu}}}{m_\nu} = \frac{\rho_{\tilde{\nu}}}{\rho_\nu} \sim 1. \quad (21)$$

Thus, we conclude that in models where R-parity violation is the only source of lepton number violation, $r_\nu \simeq \mathcal{O}(1)$, and no enhancement of the sneutrino mass splitting is possible.

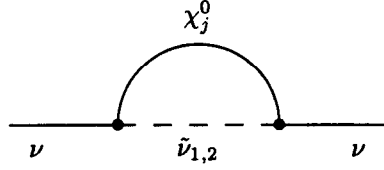


Figure 1. One-loop contribution to the neutrino mass due to sneutrino mass splitting.

4. Loop Effects

In the previous sections, we took into account only tree level contributions to the neutrino and sneutrino mass matrices. However, in some cases, one-loop effects can substantially modify r_ν . In general, the existence of a sneutrino mass splitting generates a one-loop contribution to the neutrino mass. Note that this effect is generic, and is independent of the mechanism that generates the sneutrino mass splitting. Similarly, the existence of a Majorana neutrino mass generates a one-loop contribution to the sneutrino mass splitting. In models discussed in this paper we found that $r_\nu \gtrsim 1$ at tree level, and therefore, the latter effect can be safely neglected. In contrast, the one-loop correction to the neutrino mass is potentially significant, and may dominate the tree-level mass. We have computed exactly the one-loop contribution to the neutrino mass $[m_\nu^{(1)}]$ from neutralino/sneutrino loops shown in Fig. 1. In the limit of $m_\nu, \Delta m_{\tilde{\nu}} \ll m_{\tilde{\nu}}$, the formulae simplify, and we find

$$m_\nu^{(1)} = \frac{g^2 \Delta m_{\tilde{\nu}}}{32\pi^2 \cos^2 \theta_W} \sum_j f(y_j) |Z_{jZ}|^2, \quad (22)$$

where $f(y_j) = \sqrt{y_j} [y_j - 1 - \ln(y_j)] / (1 - y_j)^2$, with $y_j \equiv m_{\tilde{\nu}}^2 / m_{\tilde{\chi}_j^0}^2$, and $Z_{jZ} \equiv Z_{j2} \cos \theta_W - Z_{j1} \sin \theta_W$ is the neutralino mixing matrix element that projects out the \tilde{Z} eigenstate from the j th neutralino. One can check that $f(y_j) < 0.566$, and for typical values of y_j between 0.1 and 10, $f(y_j) > 0.25$. Since Z is a unitary matrix, we find $m_\nu^{(1)} \approx 10^{-3} m_\nu^{(0)} r_\nu^{(0)}$, where $r_\nu^{(0)}$ is the tree-level ratio. If $r_\nu^{(0)} \gtrsim 10^3$, then the one-loop contribution to the neutrino mass cannot be neglected. Moreover, r_ν cannot be arbitrarily large without unnatural fine-tuning. Writing the neutrino mass as $m_\nu = m_\nu^{(0)} + m_\nu^{(1)}$, and assuming no unnatural cancellation between the two terms, we conclude that

$$r_\nu \equiv \frac{\Delta m_{\tilde{\nu}}}{m_\nu} \lesssim 2 \times 10^3. \quad (23)$$

5. Phenomenological Consequences

Based on the analysis presented above, we take $1 \lesssim r_\nu \lesssim 10^3$. If r_ν is near its maximum, and if there exists a neutrino mass in the MeV range, then the corresponding sneutrino mass difference is in the GeV range. Such a large mass splitting can be observed directly in the laboratory. For example, in e^+e^- annihilation, third generation sneutrinos are produced via Z -exchange. Since the two sneutrino mass eigenstates are CP-even and CP-odd respectively, sneutrino pair production occurs only via $e^+e^- \rightarrow \tilde{\nu}_1\tilde{\nu}_2$. In particular, the pair production processes $e^+e^- \rightarrow \tilde{\nu}_i\tilde{\nu}_i$ (for $i = 1, 2$) are forbidden. If the low-energy supersymmetric model incorporates some R-parity violation, then sneutrinos can be produced as an s-channel resonance in e^+e^- collisions [17, 18]. Then, for a sneutrino mass difference in the GeV range, two sneutrino resonant peaks could be distinguished.

A smaller sneutrino mass splitting can be probed using the same sign dilepton signal if $x_{\tilde{\nu}} \gtrsim 1$. Here we must rely on sneutrino oscillations. Assume that the sneutrino decays with significant branching ratio via chargino exchange: $\tilde{\nu} \rightarrow \ell^\pm + X$. Since this decay conserves lepton number, the lepton number of the decaying sneutrino is tagged by the lepton charge. Then in $e^+e^- \rightarrow \tilde{\nu}_1\tilde{\nu}_2$, the probability of a same sign dilepton signal is

$$P(\ell^+\ell^+) + P(\ell^-\ell^-) = \chi_{\tilde{\nu}} [\text{BR}(\tilde{\nu} \rightarrow \ell^\pm + X)]^2, \quad (24)$$

where

$$\chi_{\tilde{\nu}} \equiv x_{\tilde{\nu}}^2 / [2(1 + x_{\tilde{\nu}}^2)], \quad (25)$$

is the integrated oscillation probability, which arises in the same way as the corresponding quantity that appears in the analysis of B meson oscillations [11]. At hadron collider, where the sneutrino are produced mainly via $\chi_2^+ \rightarrow \tilde{\nu}\ell^+$ the probability of a same sign dilepton signal is

$$P(\ell^+\ell^+) + P(\ell^-\ell^-) = \chi_{\tilde{\nu}} [\text{BR}(\tilde{\nu} \rightarrow \ell^\pm + X)]. \quad (26)$$

We have considered the constraints on the supersymmetric model imposed by the requirements that $x_{\tilde{\nu}} \sim \mathcal{O}(1)$ and $\text{BR}(\tilde{\nu} \rightarrow \ell^\pm + X) \sim 0.5$. We examined two cases depending on whether the dominant $\tilde{\nu}$ decays involve two-body or three-body final states.

If the dominant sneutrino decay involves two-body final states, then we must assume that $m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}^\pm} < m_{\tilde{\nu}}$. Then, the widths of the two leading sneutrino decay channels, with the latter summed over both final state charges, are given by [18, 19]

$$\begin{aligned} \Gamma(\tilde{\nu} \rightarrow \tilde{\chi}_j^0 \nu) &= \frac{g^2 |Z_{jZ}|^2 m_{\tilde{\nu}}}{32\pi \cos^2 \theta_W} B(m_{\tilde{\chi}_j^0}^2 / m_{\tilde{\nu}}^2), \\ \Gamma(\tilde{\nu} \rightarrow \tilde{\chi}^\pm \ell^\mp) &= \frac{g^2 |V_{11}|^2 m_{\tilde{\nu}}}{8\pi} B(m_{\tilde{\chi}^\pm}^2 / m_{\tilde{\nu}}^2), \end{aligned} \quad (27)$$

where $B(x) \equiv (1-x)^2$, V_{11} is one of the mixing matrix elements in the chargino sector, and Z_{jZ} is the neutralino mixing matrix element defined below eq. (22), and we take $m_\ell = 0$. For example, for $m_{\tilde{\nu}} \sim \mathcal{O}(m_Z)$ we find

$$\begin{aligned}\Gamma(\tilde{\nu} \rightarrow \chi_j^0 \nu) &\approx \mathcal{O}\left(|Z_{jZ}|^2 B(m_{\tilde{\chi}_j^0}^2/m_{\tilde{\nu}}^2) \times 1 \text{ GeV}\right) \\ \Gamma(\tilde{\nu} \rightarrow \chi^+ \ell) &\approx \mathcal{O}\left(|V_{11}|^2 B(m_{\tilde{\chi}^+}^2/m_{\tilde{\nu}}^2) \times 1 \text{ GeV}\right).\end{aligned}\quad (28)$$

Typically, $B \gtrsim 10^{-2}$ in eq. (27). Thus, for the third generation sneutrino, a significant same-sign dilepton signal can be generated with $m_{\nu_\tau} = 10 \text{ MeV}$, even if $\tau_\nu \sim 1$ and the light chargino/neutralino mixing angles are of $\mathcal{O}(1)$. If the lightest chargino and two lightest neutralinos are Higgsino-like, then the mixing angle factors in eq. (27) are suppressed. For $|\mu| \sim m_Z$ and gaugino mass parameters not larger than 1 TeV, the square of the light chargino/neutralino mixing angles must be of $\mathcal{O}(10^{-2})$ or larger. Thus, if τ_ν is near its maximum value ($\tau_\nu \sim 10^3$), then one can achieve $x_{\tilde{\nu}} \sim 1$ for *neutrino* masses as low as about 100 eV.

If no open two-body decay channel exists, then we must consider the possible sneutrino decays into three-body final states. In this case we require that $m_{\tilde{\nu}} < m_{\tilde{\chi}^0}, m_{\tilde{\chi}^+}$. Again, we assume that there exists a significant chargino-mediated decay rate with charged leptons in the final state. The latter occurs in models in which the $\tilde{\tau}_R$ is lighter than the sneutrino. In this case, the rate for chargino-mediated three-body decay $\tilde{\nu}_\ell \rightarrow \tilde{\tau}_R \nu_\tau \ell$ can be significant. The $\tilde{\tau}_R$ with $m_{\tilde{\tau}_R} < m_{\tilde{\nu}}$ can occur in radiative electroweak breaking models of low-energy supersymmetry if $\tan\beta$ is large. However, in the context of the MSSM, such a scenario would require that $\tilde{\tau}_R$ is the lightest supersymmetric particle (LSP), a possibility strongly disfavored by astrophysical bounds on the abundance of stable heavy charged particles. Thus, we go beyond the usual MSSM assumptions and assume that the $\tilde{\tau}_R$ decays. This can occur in gauge-mediated supersymmetry breaking models [20] where $\tilde{\tau}_R \rightarrow \tau \tilde{g}_{3/2}$, or in R-parity violating models where $\tilde{\tau}_R \rightarrow \tau \nu$. Here, we have assumed that intergenerational lepton mixing is small; otherwise the $\Delta L = 2$ sneutrino mixing effect is diluted.

We have computed the chargino and neutralino-mediated three-body decays of $\tilde{\nu}_\ell$. In the analysis presented here, we have not considered the case of $\ell = \tau$, which involves a more complex final state decay chain containing two τ -leptons. For simplicity, we present analytic formulae in the limit where the mediating chargino and neutralinos are much heavier than the $\tilde{\tau}_R$. In addition, we assume that the lightest neutralino is dominated by its bino component. We have checked that our conclusions do not depend strongly on these approximations. Then, the rates for the chargino and neutralino-mediated sneutrino decays (the latter summed over both final

state charges) are

$$\begin{aligned}\Gamma(\tilde{\nu}_\ell \rightarrow \ell^- \tilde{\tau}^+ \nu_\tau) &= \frac{g^4 m_\nu^3 m_\tau^2 \tan^2 \beta f_{\tilde{\chi}^+}(m_\tau^2/m_\nu^2)}{1536\pi^3(m_W^2 \sin 2\beta - M_2\mu)^2}, \\ \Gamma(\tilde{\nu}_\ell \rightarrow \tau^\pm \tilde{\tau}^\mp \nu_\ell) &= \frac{g'^4 m_\nu^5 f_{\tilde{\chi}^0}(m_\tau^2/m_\nu^2)}{3072\pi^3 M_1^4},\end{aligned}\quad (29)$$

for $\ell = \mu, e$, where the M_i are gaugino mass parameters and

$$\begin{aligned}f_{\tilde{\chi}^+}(x) &= (1-x)(1+10x+x^2) + 6x(1+x)\ln x, \\ f_{\tilde{\chi}^0}(x) &= 1 - 8x + 8x^3 - x^4 - 12x^2 \ln x.\end{aligned}\quad (30)$$

As an example, for $\tan \beta = 20$ (consistent with a light $\tilde{\tau}_R$ as noted above) and $m_\tau^2/m_\nu^2 = 0.64$, reasonable values for the other supersymmetric parameters can be found such that $\Gamma(\tilde{\nu}_\ell \rightarrow \ell^\pm \tilde{\tau}^\mp \nu_\tau) \sim \Gamma(\tilde{\nu}_\ell \rightarrow \tau^\pm \tilde{\tau}^\mp \nu_\ell) \sim \mathcal{O}(1 \text{ eV})$. In this case, for $r_\nu \sim 1 [10^3]$, a significant like-sign dilepton signal could be observed for light neutrino masses as low as $1 \text{ eV} [10^{-3} \text{ eV}]$.

6. Conclusions

Non-zero Majorana neutrino masses imply the existence of $\Delta L = 2$ phenomena. In particular, in supersymmetric models, we expect sneutrino-antisneutrino mixing. The resulting sneutrino mass splitting is generally of the same order as the light neutrino mass, although an enhancement of up to three orders of magnitude is conceivable. If the mass of the tau neutrino is near its present experimental bound, $m_\nu \sim 10 \text{ MeV}$, then it may be possible to directly observe the sneutrino mass splitting in the laboratory. Even if neutrino masses are small (of order 1 eV), some supersymmetric models yield an observable sneutrino oscillation signal. Remarkably, model parameters exist where sneutrino mixing phenomena are detectable for *neutrino* masses as low as $m_\nu \sim 10^{-3} \text{ eV}$ (a mass suggested by the solar neutrino anomaly). Thus, sneutrino mixing and oscillations could provide a novel opportunity to probe lepton-number violating phenomena in the laboratory.

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