Sneutrino Mixing Phenomena

Yuval Grossman[1] and Howard E. Haber[2]
[1]Stanford Linear Accelerator Center, Stanford University, Stanford, California 94309
[2]Santa Cruz Institute for Particle Physics, University of California, Santa Cruz, California 95064
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In any model with nonzero Majorana neutrino masses, the sneutrino and antisneutrino of the supersymmetric extended theory mix. We outline the conditions under which sneutrino-antisneutrino mixing is experimentally observable. The mass splitting of the sneutrino mass eigenstates and sneutrino oscillation phenomena are considered. [S0031-9007(97)03063-9]

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In the standard model, neutrinos are exactly massless [1]. However, a number of experimental hints suggest that neutrinos may have a small mass. The solar neutrino puzzle can be solved by invoking the Mikheyev-Smirnov-Wolfenstein (MSW) mechanism, with the neutrino squared-mass difference of $\Delta m^2 \simeq 6 \times 10^{-6}$ eV^2 [2]. The atmospheric neutrino puzzle could be explained if $\Delta m^2 \simeq 10^{-2}$ eV^2 [3]. The Liquid Scintillation Neutrino Detector (LSND) experiment has reported a signal that, if interpreted as neutrino oscillations, implies $\Delta m^2 \sim \mathcal{O}(1 \text{ eV}^2)$ [4]. To accommodate this data, the standard model must be extended; the simplest models simply add Majorana neutrino mass terms that violate lepton number (L) by two units.

One must also extend the standard model in order to accommodate light Higgs bosons in a more fundamental unified theory that incorporates gravity. Models of low-energy supersymmetry [5] are attractive candidates for the theory of TeV scale physics. However, in the minimal supersymmetric extension of the standard model (MSSM), neutrinos are also exactly massless. In this paper, we wish to consider a supersymmetric extension of an extended standard model that contains Majorana neutrino masses. In such models, the lepton number violation can generate interesting phenomena in the sector of supersymmetric leptons. The effect of $\Delta L = 2$ operators is to introduce a mass splitting and mixing into the sneutrino–antisneutrino system (this observation was also made recently in Ref. [6]). The sneutrino and antisneutrino will then no longer be mass eigenstates.

This phenomena is analogous to the effect of a small $\Delta B = 2$ perturbation to the leading $\Delta B = 0$ mass term in the B system [7]. This 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, $\Delta m_B / m_B = 6 \times 10^{-14}$ [8], can be measured by observing flavor oscillations. The flavor is tagged in B decays by the final state lepton charge. Since $x_d \equiv \Delta m_B / \Gamma_B \approx 0.7$ [8], there is time for the flavor to oscillate before the meson decays. Then the time-integrated same sign dilepton signal is used to determine Δm_B.

The sneutrino system can exhibit similar behavior. The lepton number is tagged in sneutrino decay using the charge of the outgoing lepton. The relevant scale is the sneutrino width (as emphasized in the context of lepton flavor oscillation in Ref. [9]). If the sneutrino mass splitting is large, namely,

$$x_{\tilde{\nu}} \equiv \frac{\Delta m_{\tilde{\nu}}}{\Gamma_{\tilde{\nu}}} \gtrsim 1 \,, \tag{1}$$

and the sneutrino branching ratio into a charged lepton is significant, then a measurable same-sign dilepton signal is expected.

The neutrino mass and the sneutrino mass splitting are related as a consequence of the lepton number violating interactions and supersymmetry breaking. Thus, we can use upper bounds (or indications) of neutrino masses to set bounds on the sneutrino mass splitting. At present, neutrino mass bounds obtained from direct laboratory measurements imply [8]: $m_{\nu_e} \lesssim 10$ eV, $m_{\nu_\mu} \le 0.17$ MeV, and $m_{\nu_\tau} \le 24$ MeV. Cosmological constraints require stable neutrinos to be lighter than about 100 eV. For example, models of mixed dark matter require a neutrino mass of order 10 eV [10]. For unstable neutrinos, the mass limits are more complex and model dependent [11]. In this paper we will consider the consequences of two cases: (i) ν_τ with a mass near its present laboratory upper limit, and (ii) light neutrinos of mass less than 100 eV.

Some model-independent relations among the neutrino and sneutrino $\Delta L = 2$ masses (and other $\Delta L = 2$ phenomena) have been derived in Ref. [6]. However, in order to derive specific results, it is useful to exhibit an explicit model of lepton number violation. In the following, we concentrate on the seesaw model for neutrino masses [1], as it exhibits all the interesting features. We compute the sneutrino mass splitting in this model and discuss its implications for sneutrino phenomenology at e^+e^- colliders. (We also briefly mention some consequences of lepton number violation arising from R-parity nonconservation.) A more complete presentation will be given in Ref. [12].

The supersymmetric seesaw model.—Consider an extension of the MSSM where one adds a right-handed neutrino superfield, $\hat{N}$, with a bare mass $M \gg m_Z$. For simplicity 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. [13])

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

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

$$V_{\text{soft}} = m_{\tilde{L}}^2 \tilde{\nu}^* \tilde{\nu} + m_{\tilde{N}}^2 \tilde{N}^* \tilde{N} + (\lambda A_\nu H_2^2 \tilde{\nu} \tilde{N}^* + M B_N \tilde{N} \tilde{N} + \text{H.c.}) . \qquad (3)$$

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

$$\mathcal{L}_{\text{mass}} = \tfrac{1}{2} (\phi_1 \quad \phi_2) \begin{pmatrix} \mathcal{M}_+^2 & 0 \\ 0 & \mathcal{M}_-^2 \end{pmatrix} \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix}, \qquad (4)$$

where $\phi_i \equiv (\tilde{\nu}_i \quad \tilde{N}_i)$ and $\mathcal{M}_\pm^2$ consist of the following 2×2 blocks:

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

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,

$$m_{\tilde{\nu}_{1,2}}^2 = m_{\tilde{L}}^2 + \tfrac{1}{2} m_Z^2 \cos 2\beta \mp \tfrac{1}{2} \Delta m_{\tilde{\nu}}^2 , \qquad (5)$$

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 $\Delta m_{\tilde{\nu}}^2 = 2 m_{\tilde{\nu}} \Delta m_{\tilde{\nu}}$, 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 is given by (to leading order in $1/M$)

$$r_\nu \equiv \frac{\Delta m_{\tilde{\nu}}}{m_\nu} \simeq \frac{2(A_\nu - \mu \cot\beta - B_N)}{m_{\tilde{\nu}}} . \qquad (6)$$

The magnitude of r_ν depends on various supersymmetric parameters. Naturalness constrains supersymmetric mass parameters associated with particles with nontrivial electroweak quantum numbers to be roughly of order m_Z [14]. 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 $\hat{N}$. In particular, $M \gg m_Z$, since this drives the seesaw mechanism. Since M is a supersymmetry-conserving parameter, the seesaw 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 $\hat{N}$ is an electroweak gauge group singlet superfield, supersymmetry-breaking terms associated with it need not be 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. (6)], 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.

We have also considered other possible models of lepton number violation [12]. For example, in models of R-parity violation (but with no right-handed neutrino), a sneutrino mass splitting is also generated whose magnitude is of order the corresponding neutrino mass. Thus, 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.

Loop effects.—In the previous section, all formulas given involved tree-level parameters. However, in some cases, one-loop effects can substantially modify Eq. (6). 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. However, the latter effect can be safely neglected. Any one-loop contribution to the sneutrino mass splitting must be roughly $\Delta m_{\tilde{\nu}}^{(1)} \sim (g^2/16\pi^2) m_\nu$; thus the tree-level result $r_\nu \gtrsim \mathcal{O}(1)$ cannot be significantly modified. In contrast, the one-loop correction to the neutrino mass is potentially significant, and may dominate the tree-level mass $[m_\nu^{(0)} \simeq m_D^2/M]$. 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

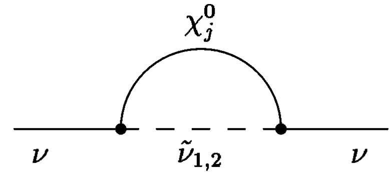

FIG. 1. One-loop contribution to the neutrino mass due to sneutrino mass splitting.

of $m_\nu, \Delta m_{\tilde{\nu}} \ll m_{\tilde{\nu}}$, the formulas simplify, and we find

$$m_\nu^{(1)} = \frac{g^2 \Delta m_{\tilde{\nu}}}{32\pi^2 \cos^2\theta_W} \sum_j f(y_j) |Z_{jZ}|^2, \qquad (7)$$

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 jth 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 computed in Eq. (6). 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

$$r_\nu \equiv \frac{\Delta m_{\tilde{\nu}}}{m_\nu} \lesssim 2 \times 10^3. \qquad (8)$$

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 $e^+e^- \to \tilde{\nu}_1\tilde{\nu}_2$. In particular, the pair production processes $e^+e^- \to \tilde{\nu}_i\tilde{\nu}_i$ (for $i = 1, 2$) are forbidden. If the low-energy supersymmetric model incorporates some R-parity violation, then sneutrinos can be produced in e^+e^- via an s-channel resonance [15,16]. 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 in e^+e^- annihilation 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: $\tilde{\nu} \to \ell^\pm + X$. Since this decay conserves lepton number, the lepton number of the decaying sneutrino is tagged by the lepton charge. Then in $e^+e^- \to \tilde{\nu}_1\tilde{\nu}_2$, the probability of a same-sign dilepton signal is

$$P(\ell^+\ell^+) + P(\ell^-\ell^-) = \chi_{\tilde{\nu}}[BR(\tilde{\nu} \to \ell^\pm + X)]^2, \qquad (9)$$

where $\chi_{\tilde{\nu}} \equiv x_{\tilde{\nu}}^2/[2(1 + x_{\tilde{\nu}}^2)]$ is the integrated oscillation probability, which arises in the same way as the corresponding quantity that appears in the analysis of B meson oscillations [7]. We have considered the constraints on the supersymmetric model imposed by the requirements that $x_{\tilde{\nu}} \sim \mathcal{O}(1)$ and B$(\tilde{\nu} \to \ell^\pm + X) \sim 0.5$. 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 $m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}^+} < m_{\tilde{\nu}}$. Then, the widths of the two leading sneutrino decay channels are [16,17]

$$\Gamma(\tilde{\nu} \to \tilde{\chi}_j^0 \nu) = \frac{g^2 |Z_{iZ}|^2 m_{\tilde{\nu}}}{32\pi \cos^2\theta_W} B(m_{\tilde{\chi}_j^0}^2/m_{\tilde{\nu}}^2),$$
$$\Gamma(\tilde{\nu} \to \tilde{\chi}^+ \ell^-) = \frac{g^2 |V_{11}|^2 m_{\tilde{\nu}}}{16\pi} B(m_{\tilde{\chi}^+}^2/m_{\tilde{\nu}}^2), \qquad (10)$$

(assuming $m_\ell = 0$), 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. (7). For example, for $m_{\tilde{\nu}} \sim \mathcal{O}(m_Z)$ we find $\Gamma(\tilde{\nu} \to \chi_j^0 \nu) \approx \mathcal{O}[|Z_{jZ}|^2 B(m_{\tilde{\chi}_j^0}^2/m_{\tilde{\nu}}^2) \times 1 \text{ GeV}]$ and $\Gamma(\tilde{\nu} \to \chi^+ \ell) \approx \mathcal{O}[|V_{11}|^2 B(m_{\tilde{\chi}^+}^2/m_{\tilde{\nu}}^2) \times 1 \text{ GeV}]$. We require that the sneutrino and chargino are sufficiently separated in mass, so that the emitted charged lepton will not be too soft and can be identified experimentally. This implies that $B \gtrsim 10^{-2}$ in Eq. (10). Thus, for the third generation sneutrino, a significant same-sign dilepton signal can be generated with $m_{\nu_\tau} = 10$ MeV, even if $r_\nu \sim 1$ 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. (10) 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 r_ν is near its maximum value ($r_\nu \sim 10^3$), 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 $m_{\tilde{\nu}} < m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}^+}$. 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 $\tilde{\nu}_\ell \to \ell^- \tilde{\tau}_R^+ \nu_\tau$ 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 [18]. 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 [19] where $\tilde{\tau}_R \to \tau + \tilde{g}_{3/2}$, or in R-parity violating models where $\tilde{\tau}_R \to \tau\nu$. 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 formulas in the limit where the mediating chargino and neutralinos are much heavier than the $\tilde{\tau}_R^{\pm}$. 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 are

$$\Gamma(\tilde{\nu}_\ell \to \ell^- \tilde{\tau}_R^+ \nu_\tau) = \frac{g^4 m_{\tilde{\nu}}^3 m_\tau^2 \tan^2\beta \, f_{\tilde{\chi}^+}(m_{\tilde{\tau}}^2/m_{\tilde{\nu}}^2)}{3 \times 2^9 \pi^3 (m_W^2 \sin 2\beta - M_2 \mu)^2},$$

$$\Gamma(\tilde{\nu}_\ell \to \nu_\ell \tilde{\tau}_R^{\pm} \tau^{\mp}) = \frac{g'^4 m_{\tilde{\nu}}^5 f_{\tilde{\chi}^0}(m_{\tilde{\tau}}^2/m_{\tilde{\nu}}^2)}{3 \times 2^{10} \pi^3 M_1^4}, \tag{11}$$

for $\ell = \mu, e$ where the M_i are gaugino mass parameters, $f_{\tilde{\chi}^+}(x) = (1 - x)(1 + 10x + x^2) + 6x(1 + x)\ln x$ and $f_{\tilde{\chi}^0}(x) = 1 - 8x + 8x^3 - x^4 - 12x^2 \ln x$. As an example, for $\tan\beta = 20$ (consistent with a light $\tilde{\tau}_R$ as noted above) and $m_{\tilde{\tau}}^2/m_{\tilde{\nu}}^2 = 0.64$, reasonable values for the other supersymmetric parameters can be found such that $\Gamma(\tilde{\nu}_\ell \to \ell^- \tilde{\tau}_R^+ \nu_\tau) \sim \Gamma(\tilde{\nu}_\ell \to \nu_\ell \tilde{\tau}_R^{\pm} \tau^{\mp}) \sim O(1 \text{ eV})$. 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 eV $[10^{-3}$ eV$]$.

In conclusion, nonzero Majorana neutrino masses imply the existence of $\Delta L = 2$ phenomena. In low-energy supersymmetric models, such phenomena also lead to sneutrino-antisneutrino mixing with the corresponding mass eigenstates split in mass. The sneutrino mass splitting is generally of the same order as the light neutrino mass, although an enhancement of up to 3 orders of magnitude is conceivable. If the mass of the ν_τ is near its present experimental bound, 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 at e^+e^- colliders. Remarkably, model parameters exist where sneutrino mixing phenomena are detectable for *neutrino* masses as low as $m_\nu \sim 10^{-3}$ 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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