Sneutrino-Antisneutrino Mixing and Neutrino Mass in Anomaly-Mediated Supersymmetry Breaking Scenario

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In supersymmetric models with nonzero Majorana neutrino mass, the sneutrino and antisneutrino mix, which may lead to same-sign dilepton signals in future collider experiments. We point out that the anomaly-mediated supersymmetry breaking scenario has a good potential to provide an observable rate of such signals for the neutrino masses suggested by the atmospheric and solar neutrino oscillations. It is noted also that the sneutrino-antisneutrino mixing can provide much stronger information on some combinations of the neutrino masses and mixing angles than the neutrino experiments.

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Current data from the atmospheric and solar neutrino experiments strongly suggest that the neutrinos have small but nonzero masses [1]. As was pointed out in Refs. [2–5], in supersymmetric models with nonzero Majorana neutrino mass, the sneutrino ($\tilde{\nu}$) and antisneutrino ($\tilde{\nu}^*$) mix to each other. The mixing rate is generically given by

$$\Delta m_{\tilde{\nu}} = C_\nu m_\nu / m_{\tilde{\nu}}, \tag{1}$$

where m_ν is the lepton number-violating ($\Delta L = 2$) Majorana neutrino mass, $m_{\tilde{\nu}}$ is the lepton-number-conserving ($\Delta L = 0$) sneutrino mass, and C_ν is determined mainly by the soft supersymmetry (SUSY) breaking in the lepton-number-violating sector of the underlying theory. The atmospheric and solar neutrino data give the neutrino square mass differences $\Delta m_\nu^2 \lesssim \mathcal{O}(10^{-3})$ eV^2, suggesting that m_ν does not exceed $\mathcal{O}(1)$ eV. Also the consideration of radiative corrections to m_ν induced by $\Delta m_{\tilde{\nu}}$ [3] leads to the bound $\Delta m_{\tilde{\nu}}/m_\nu \lesssim \mathcal{O}(4\pi/\alpha)$, implying $\Delta m_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV. Such a small mixing rate can be probed by the $\tilde{\nu}$-$\tilde{\nu}^*$ oscillation which would result in same-sign dilepton signals when the sneutrino pairs decay into charged leptons.

To have an observable rate of same-sign dilepton signals, $\tilde{\nu}$ must have enough time to mix with $\tilde{\nu}^*$ before it decays. For this, we need the sneutrino decay width $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV in view of $\Delta m_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV. Such a small decay rate would not be possible if the two-body decay channel $\tilde{\nu} \to \nu\tilde{\chi}^0$ or $\ell^-\tilde{\chi}^+$ was open for the neutralinos $\tilde{\chi}^0$ or charginos $\tilde{\chi}^+$. It was pointed out in [3] that the most plausible scenario for $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV is to have

$$m_{\tilde{\tau}_1} < m_{\tilde{\nu}} < m_{\tilde{\chi}^0}, m_{\tilde{\chi}^+}, \tag{2}$$

where $\tilde{\tau}_1$ denotes the lighter stau. Then sneutrinos decay mainly into three-body final states with a sizable branching ratio into a charged lepton: $\tilde{\nu} \to \ell^-\tilde{\tau}_1^+\nu_\tau, \nu\tilde{\tau}_1^\pm\tau^\mp$ with $\Gamma_{\tilde{\nu}} \lesssim \mathcal{O}(1)$ keV. The mass hierarchy (2) would mean that the stau is the lightest supersymmetric particle (LSP) in the minimal supersymmetric standard model (MSSM) sector, which would not be cosmologically allowed if it was stable. This difficulty can be easily avoided if one assumes a light singlet fermion ψ which has very weak couplings to the MSSM sector, e.g., a light gravitino or axino, with which $\tilde{\tau}_1$ decays into $\tau\psi$. Alternatively, one may introduce a tiny R-parity-violating coupling which would trigger $\tilde{\tau}_1 \to \ell\nu$. Still $\tilde{\tau}_1$ can live long enough inside the detector, so it is clearly distinguished from other charged sleptons.

Obviously, for a given neutrino mass, models with bigger $C_\nu/m_{\tilde{\nu}} = \Delta m_{\tilde{\nu}}/m_\nu$ have a better prospect for observable $\tilde{\nu} - \tilde{\nu}^*$ mixing. In this paper, we point out that the anomaly-mediated SUSY breaking (AMSB) scenario [6] generically predicts $C_\nu/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha) \gg 1$ if the neutrino mass is generated by SUSY-preserving dynamics at the high energy scale as in the conventional seesaw model [7]. Furthermore, the mass hierarchy (2) can be obtained on a sizable portion of the phenomenologically allowed parameter space of the AMSB model. These features are not shared by the minimal supergravity model [8] or the gauge-mediated SUSY breaking model [9], so the AMSB model has much better potential to provide observable same-sign dilepton signals induced by the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing than other SUSY breaking models. An interesting feature of the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing in the AMSB scenario is that it provides rather strong information on the neutrino mass matrix elements $(m_\nu)_{ii} = \sum_a U_{ia}^2 m_{\nu_a}$, where $i = e, \mu, \tau$ denote the flavor eigenstates while $a = 1, 2, 3$ stand for the neutrino mass eigenstates with the Maki-Nakagaa-Sakata (MNS) mixing matrix U_{ia}. Atmospheric neutrino data suggest $(m_\nu)_{\mu\mu} \simeq 3 \times 10^{-2}$ eV for hierarchical neutrino masses, while $(m_\nu)_{\mu\mu}$ can be bigger if the neutrino masses are approximately degenerate. Then the same-sign dimuon events from $\tilde{\nu}_\mu$-$\tilde{\nu}_\mu^*$ mixing can be used to distinguish $(m_\nu)_{\mu\mu} = 3 \times 10^{-2}$ eV from a bigger value of $(m_\nu)_{\mu\mu}$. Also the same-sign dielectron events from $\tilde{\nu}_e$-$\tilde{\nu}_e^*$ mixing can be used to probe $(m_\nu)_{ee}$ down to the order of 10^{-4} eV which is much smaller than the current bound on $(m_\nu)_{ee}$ from ν-less double beta decays. In the following, we will examine these points in more detail.

Anomaly mediation assumes that SUSY breaking is transmitted to the MSSM fields *mainly* by the Weyl compensator superfield Φ_0 of the supergravity multiplet [6]:

$$\Phi_0 = 1 + \theta^2 M_{\rm aux}\,, \tag{3}$$

where $M_{\rm aux} = \mathcal{O}(m_{3/2})$. The couplings of Φ_0 to the MSSM fields are determined by the super-Weyl invariance, yielding

$$\begin{aligned} S_{\rm eff} = &\int d^4x d^4\theta \Big[Z_I(Q/\sqrt{\Phi_0\Phi_0^*})\Phi_I^*\Phi_I + \frac{1}{8} g_a^{-2}(Q/\sqrt{\Phi_0\Phi_0^*}) DV^a \bar{D}^2 DV^a \Big] \\ &+ \Big[\int d^4x d^2\theta \Big(y_{IJK}\Phi_I\Phi_J\Phi_K + \frac{\Phi_0}{M}\gamma_{IJKL}\Phi_I\Phi_J\Phi_K\Phi_L \Big) + \text{H.c.} \Big], \end{aligned} \tag{4}$$

where Q denotes the renormalization scale, D and $\bar{D}$ are the supercovariant derivatives on the real gauge superfields V_a, and Φ_I are the chiral matter superfields.

The above $S_{\rm eff}$ leads to a tachyonic slepton, so one needs some additional source of SUSY breaking. The simplest possibility is an additional universal soft scalar square mass m_0^2 introduced at the grand unified theory (GUT) scale $M_{\rm GUT} \simeq 2 \times 10^{16}$ GeV. One then finds the following forms of soft supersymmetry breaking terms [6,10]:

$$\begin{aligned} \mathcal{L}_{\rm soft} = &\, m_I^2|\phi_I|^2 \\ &+ \Big(\frac{1}{2} M_a \lambda^a\lambda^a + A_{IJK} y_{IJK}\phi_I\phi_J\phi_K \\ &+ \frac{C_{IJKL}\gamma_{IJKL}}{M}\phi_I\phi_J\phi_K\phi_L + \text{H.c.} \Big), \end{aligned} \tag{5}$$

where the gaugino masses M_a, the soft scalar masses m_I, and the soft coefficients A_{IJK}, C_{IJKL} are given by

$$\begin{aligned} M_a &= -\frac{b_a\alpha_a}{4\pi} M_{\rm aux}, \\ m_I^2 &= -\frac{1}{4}\frac{d\gamma_I}{d\ln Q}|M_{\rm aux}|^2 + m_0^2, \\ A_{IJK} &= \frac{1}{2}(\gamma_I + \gamma_J + \gamma_K) M_{\rm aux}\,, \\ C_{IJKL} &= \frac{1}{2}(2 + \gamma_I + \gamma_J + \gamma_K + \gamma_L) M_{\rm aux}\,. \end{aligned} \tag{6}$$

Here $b_a = (3, -1, -33/5)$ are the one-loop beta function coefficients for $SU(3)_c \times SU(2)_L \times U(1)_Y$ in the GUT normalization and $\gamma_I = d\ln Z_I/d\ln Q$ are the anomalous dimensions of Φ_I.

Applying the above results to the sneutrino-antisneutrino mixing is rather straightforward. To be specific, we will assume that the neutrino masses are generated (mainly) by supersymmetry preserving dynamics at an energy scale M far above the weak scale. This high energy dynamics may be the exchange of heavy singlet neutrino with mass M [7], or the exchange of heavy triplet Higgs boson [11]. At the weak scale, the theory can be described by an effective superpotential including the super-Weyl invariant dimension-5 operators for neutrino masses and also the associated soft SUSY breaking terms,

$$\begin{aligned} \Delta W_{\rm eff} &= \frac{\Phi_0}{M}\gamma_{ij}(L_iH_2)(L_jH_2), \\ \Delta\mathcal{L}_{\rm soft} &= \frac{C_{ij}\gamma_{ij}}{M}(\tilde{\ell}_i h_2)(\tilde{\ell}_j h_2)\,, \end{aligned} \tag{7}$$

where $L_i(i = e, \mu, \tau)$ and $H_\alpha(\alpha = 1, 2)$ denote the lepton and Higgs doublet superfields with the scalar components $\tilde{\ell}_i$ and h_α, respectively, and $C_{ij} \simeq M_{\rm aux}$.

After the electroweak symmetry breaking, $\Delta W_{\rm eff}$ gives a neutrino mass matrix $(m_\nu)_{ij} = 2\langle h_2\rangle^2\gamma_{ij}/M$. Including the contribution from $\Delta\mathcal{L}_{\rm soft}$, the sneutrino masses are given by

$$(m_{\tilde{\nu}}^2)_{ij}\tilde{\nu}_i^*\tilde{\nu}_j + \{\tfrac{1}{2}(\Delta m_{\tilde{\nu}}^2)_{ij}\tilde{\nu}_i\tilde{\nu}_j + \text{H.c.}\}, \tag{8}$$

where the sneutrino square mass matrix $(m_{\tilde{\nu}}^2)_{ij} \simeq m_{\tilde{\nu}}^2\delta_{ij}$, with $m_{\tilde{\nu}}^2 = \frac{1}{2}M_Z^2\cos 2\beta + m_{\tilde{\ell}}^2$ for the slepton doublet square mass matrix $(m_{\tilde{\ell}}^2)_{ij} \simeq m_{\tilde{\ell}}^2\delta_{ij}$, and the $\Delta L = 2$ sneutrino square mass matrix is given by $(\Delta m_{\tilde{\nu}}^2)_{ij} = (C_{ij} + 2\mu\cot\beta)(m_\nu)_{ij}$. The $\tilde{\nu} - \tilde{\nu}^*$ mixing rate is determined by the sneutrino mass splitting $\Delta m_{\tilde{\nu}} = \Delta m_{\tilde{\nu}}^2/m_{\tilde{\nu}}$. In this regard, a *distinctive feature* of the AMSB model is that $C_{ij} \simeq M_{\rm aux}$ is induced at tree level while $m_{\tilde{\nu}}$ is loop suppressed, so that $\Delta m_{\tilde{\nu}}$ is enhanced (relative to m_ν) by the factor $M_{\rm aux}/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha)$:

$$(\Delta m_{\tilde{\nu}})_{ij} \simeq (m_\nu)_{ij} M_{\rm aux}/m_{\tilde{\nu}} = \mathcal{O}(4\pi m_\nu/\alpha)\,. \tag{9}$$

Consequently, for a given neutrino mass, the AMSB model has better potential to give a sizable $\tilde{\nu}$-$\tilde{\nu}^*$ mixing than other models with $C_{ij}/m_{\tilde{\nu}} = \mathcal{O}(1)$. Furthermore, as can be inferred from Fig. 1, a significant portion of the phenomenologically viable parameter space of the minimal AMSB model leads to the mass hierarchy (2).

Taking an analogy to the B-meson mixing, it is straightforward to compute the probability for a $\tilde{\nu}$-$\tilde{\nu}^*$ pair produced in an e^+e^- collider to yield same-sign dilepton signal [3,12]. The amplitude for $e^+(p_1) + e^-(p_2) \to \tilde{\nu}_i(q_1) + \tilde{\nu}_i^*(q_2)$ is easily computed to be

$$A_i = \tfrac{1}{2} g^2 \bar{v}(p_2)(\not{q}_1 - \not{q}_2)(X_iP_L + Y_iP_R)u(p_1)\,, \tag{10}$$

where $P_{L,R} = (1 \pm \gamma_5)/2$, $X_i = K_Z(s_W^2 - \frac{1}{2})/c_W^2 + \delta_{ie}\sum_n |V_{n1}|^2 K_n$, and $Y_i = K_Z s_W^2/c_W^2$ for the chargino $(\tilde{\chi}_n^\pm)$ mixing matrix V_{nm}, $K_Z = 1/(s - M_Z^2)$, $K_n = 1/[(p_1 - q_2)^2 - m_{\tilde{\chi}_n}^2]$, $c_W = \cos\theta_W$, $s_W = \sin\theta_W$, and the c.m. energy $\sqrt{s}$. Here K_Z represents the contributions from the Z boson-mediated s-channel diagrams, while K_n

is from the chargino-mediated t-channel diagrams. With this amplitude, the initial $\tilde{\nu}$-$\tilde{\nu}^*$ state is given by

$$|\tilde{\nu}\tilde{\nu}^*;0\rangle = \sum_i \alpha_i |\tilde{\nu}_i(\vec{q})\rangle |\tilde{\nu}_i^*(-\vec{q})\rangle + \beta_i |\tilde{\nu}_i^*(\vec{q})\rangle |\tilde{\nu}_i(-\vec{q})\rangle, \tag{11}$$

where $\alpha_i = A_i(\vec{q}_1 = \vec{q})$, $\beta_i = A_i(\vec{q}_2 = \vec{q})$ (up to overall normalization), and the momentum vector $\vec{q}$ spans only the upper hemisphere, i.e., $\cos\theta \geq 0$ for the angle θ between e^- and $\tilde{\nu}$ flight directions.

With (2), the sneutrinos decay as $\tilde{\nu} \to \ell^- \tilde{\tau}_1^+ \nu_\tau, \nu \tilde{\tau}_1^\pm \tau^\mp$. It turns out that, in most of the parameter space yielding an observable rate of same-sign dilepton signals, these decays are induced dominantly by the chargino or neutralino exchange, so the decay widths are (approximately) flavor independent. Then the effective Hamiltonian determining the evolution of (11) can be written as

$$H_{\rm eff} = \begin{pmatrix} (m_{\tilde{\nu}} - \frac{i}{2}\Gamma_{\tilde{\nu}})\delta_{ij} + (\delta m_{\tilde{\nu}})_{ij} & \frac{1}{2}(\Delta m_{\tilde{\nu}})_{ij} \\ \frac{1}{2}(\Delta m_{\tilde{\nu}})^*_{ij} & (m_{\tilde{\nu}} - \frac{i}{2}\Gamma_{\tilde{\nu}})\delta_{ij} + (\delta m_{\tilde{\nu}})^*_{ij} \end{pmatrix}, \tag{12}$$

where $\delta m_{\tilde{\nu}}$ represents the deviation from the exact degeneracy of the $\Delta L = 0$ sneutrino masses. Since $\delta m_{\tilde{\nu}} \gg \Delta m_{\tilde{\nu}}$ in our case, it is most convenient to describe the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing in the field basis in which $\delta m_{\tilde{\nu}}$ is diagonal. In the AMSB scenario, the charged lepton mass matrix can be diagonalized simultaneously with $\delta m_{\tilde{\nu}}$. In such a field basis, we find the probability P_i for the initial state (11) to produce a same-sign dilepton $\ell_i^- \ell_i^-$ or $\ell_i^+ \ell_i^+$:

$$P_i = \frac{1}{\sum_i \sigma_i} \int d\Phi_2 \frac{1}{8s} \frac{B_i^2}{(1 + x_i^2)^2} \times \left\{ \frac{1}{2} (|\alpha_i|^2 + |\beta_i|^2)(2 + x_i^2)x_i^2 + \mathrm{Re}(\alpha_i^* \beta_i) x_i^2 \right\}, \tag{13}$$

where $B_i = \mathcal{B}(\tilde{\nu}_i \to \ell_i X)$, $x_i = |(\Delta m_{\tilde{\nu}})_{ii}|/\Gamma_{\tilde{\nu}} = |(m_\nu)_{ii}| M_{\rm aux}/m_{\tilde{\nu}}\Gamma_{\tilde{\nu}}$, and σ_i denotes the total cross section for $e^+e^- \to \tilde{\nu}_i \tilde{\nu}_i^*$. Here the two-body phase space integration ($d\Phi_2$) for the initial $\tilde{\nu}\tilde{\nu}^*$ is performed for $\cos\theta \geq 0$, and we have ignored the effects suppressed by $\Delta m_{\tilde{\nu}}/\delta m_{\tilde{\nu}}$.

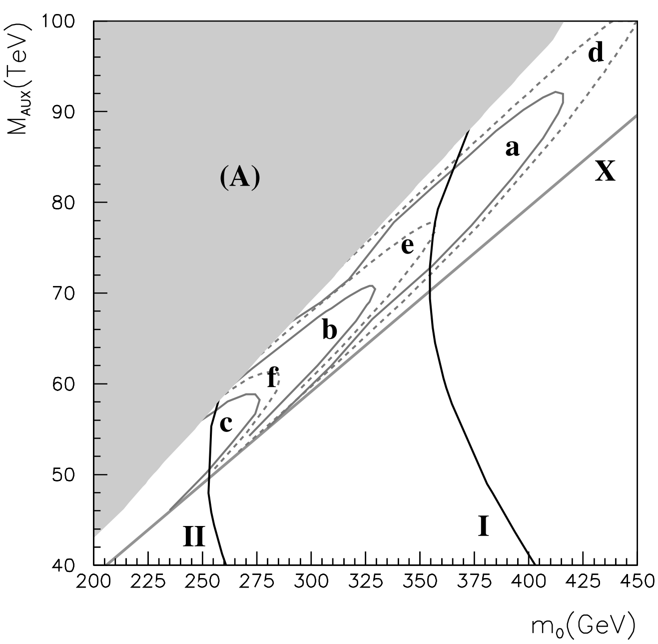

FIG. 1. Parameter regions with $\tan\beta = 5$ yielding $N_\mu \geq 20$ (inside the contour a), 10^2 (b), and 5×10^2 (c) for $(m_\nu)_{\mu\mu} = 3 \times 10^{-2}$ eV; $N_\mu \geq 10^2$ (d), 5×10^2 (e), and 2×10^3 (f) for $(m_\nu)_{\mu\mu} = 0.3$ eV, while satisfying $N_\mu \geq 5\sqrt{N_{\rm back}}$. "(A)" represents the parameter region forbidden by the stau mass bound. The upper side of line X denotes the region of LSP stau. Left sides of the lines (I) and (II) correspond to the region with $a_\mu^{\rm SUSY} \geq 5 \times 10^{-10}$ and 10^{-9}.

The same-sign dilepton probability (13) shows that in the AMSB scenario the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing provides information on the neutrino matrix elements $(m_\nu)_{\alpha\alpha} = \sum_a U_{\alpha a}^2 m_{\nu_a}$ ($\alpha = e, \mu$) where m_{ν_a} and U_{ia} denote the neutrino mass eigenvalues and the MNS mixing matrix, respectively. For instance, the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing allows us to probe $(m_\nu)_{ee}$ down to the order of 10^{-4} eV in the AMSB scenario. The information on $(m_\nu)_{\alpha\alpha}$ from the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing is different from the information on neutrino masses and mixing angles from neutrino oscillations. *So the $\tilde{\nu}$-$\tilde{\nu}^*$ can provide information on neutrino masses and mixing angles which are complementary to those from the neutrino experiments.*

Same-sign dilepton events also come from the pair-produced neutralinos which would decay $\tilde{\chi}^0\tilde{\chi}^0 \to \ell_i \ell_j \tilde{\ell}_i^* \tilde{\ell}_j^*$ ($i, j = e, \mu, \tau$) producing eventually $\ell_\alpha \ell_\beta \tilde{\tau}_1^* \tilde{\tau}_1^*$ ($\alpha, \beta = e, \mu$) through $\tilde{\ell}_i^* \to \tilde{\tau}_1^* +$ neutrinos or $\tau \to \ell_\alpha +$ neutrinos. While taking into account these backgrounds, we performed a numerical analysis to find the parameter region of the minimal AMSB model in which the same-sign dilepton events $\ell_\alpha \ell_\alpha \tilde{\tau}_1^* \tilde{\tau}_1^*$ induced by the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing are sizable enough to be distinguished from backgrounds. Specifically, we computed the number of events per year, N_α, for a future e^+e^- linear collider with the integrated luminosity 500 fb^{-1} at $\sqrt{s} = 500$ GeV under the requirement $N_\alpha \geq 5\sqrt{N_{\rm back}}$. It turns out that, in most of such parameter regions, N_α is actually bigger than $N_{\rm back}$, so the events from the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing can be clearly distinguished from the pure backgrounds.

Using the standard renormalization group (RG) analysis, the superparticle mass spectrums are obtained to compute $\Gamma_{\tilde{\nu}}$ on the parameter region of $(m_0, M_{\rm aux}, \tan\beta)$, leading to the mass hierarchy (2). When $\tan\beta$ increases for a given $M_{\rm aux}$, the mass hierarchy (2) requires a larger m_0, leading to a larger $m_{\tilde{\nu}}$. On the other hand, the enhanced left-right mixing gives an effect to lower $m_{\tilde{\tau}_1}$, so the net result is to increase $m_{\tilde{\nu}}/m_{\tilde{\tau}_1}$. As a result, for a given neutrino mass, $\Gamma_{\tilde{\nu}}$ is a sharply increasing function of $\tan\beta$, so small $\tan\beta$ is favored for sizable N_α. From a detailed numerical analysis, we find that $\tan\beta \lesssim 10$ is required to have a sizable N_α for $m_\nu \lesssim \mathcal{O}(1)$ eV.

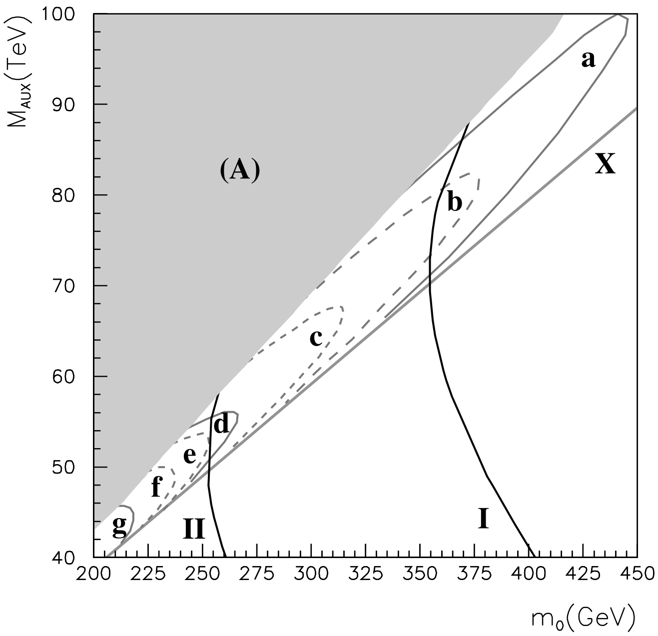

FIG. 2. Parameter regions with $\tan\beta = 5$ yielding $N_e \geq 20$ (a), 10^2 (b), and 5×10^2 (c) for $(m_\nu)_{ee} = 10^{-2}$ eV; $N_e \geq 20$ (d), 10^2 (e), and 5×10^2 (f) for $(m_\nu)_{ee} = 10^{-3}$ eV; $N_e \geq 10^2$ (g) for $(m_\nu)_{ee} = 10^{-4}$ eV, while satisfying $N_e \geq 5\sqrt{N_{\rm back}}$.

About the values of $(m_\nu)_{\mu\mu}$, we considered two cases. In the first, neutrino masses are assumed to be hierarchical, which would give $(m_\nu)_{\mu\mu} \simeq U^2_{\mu 3}\sqrt{\Delta m^2_{\rm atm}} \simeq 3 \times 10^{-2}$ eV, while, in the second case, neutrino masses are assumed to be approximately degenerate with $(m_\nu)_{\mu\mu} = 0.3$ eV. We find that the number of same-sign dimuon events per year (N_μ) for the first case is bigger than the value for the second case by about a factor of 5, so hierarchical and (approximately) degenerate neutrino masses are clearly distinguished from each other. In Fig. 1, we depict the parameter regions with $\tan\beta = 5$ yielding $N_\mu \geq 20$, 10^2, 5×10^2 for $(m_\nu)_{\mu\mu} = 3 \times 10^{-2}$ eV and $N_\mu \geq 10^2$, 5×10^2, 2×10^3 for $(m_\nu)_{\mu\mu} = 0.3$ eV. We also searched for the parameter regions with $\tan\beta = 5$ yielding $N_e \geq 20$, 10^2, 5×10^2 for $(m_\nu)_{ee} = 10^{-2}, 10^{-3}$, 10^{-4} eV and depict the results in Fig. 2. Note that the t-channel contribution to $e^+e^- \to \tilde{\nu}_e\tilde{\nu}_e^*$ enhances N_e relative to N_μ, so that we can have a sizable N_e even for $(m_\nu)_{ee} = 10^{-4}$ eV.

As available constraints on the model, we impose the Higgs, stau, chargino mass bounds, $m_h > 113.5$ GeV, $m_{\tilde{\tau}} > 89$ GeV, $m_{\tilde{\chi}_1^\pm} > 103$ GeV, respectively, and also the 2σ constraint on the $b \to s\gamma$ branching ratio, $\mathcal{B}(B \to X_s\gamma) = (2.2\text{–}4.1) \times 10^{-4}$. It has been noted that the AMSB model is severely constrained by the recent measurement of the muon anomalous magnetic moment a_μ once we require that the conventional one-loop SUSY contribution $a_\mu^{\rm SUSY} \gtrsim 10^{-9}$ which was taken as the 2σ lower bound [13]. However, the new result on the light-by-light hadronic contribution [14] essentially eliminates this difficulty. For instance, the new 1.3σ bound $a_\mu^{\rm SUSY} \gtrsim 5 \times 10^{-10}$ can be easily satisfied in the AMSB model, as shown in Figs. 1 and 2.

To conclude, we have examined the possibility of an observable same-sign dilepton signal induced by the $\tilde{\nu}$-$\tilde{\nu}^*$ mixing in the AMSB model. It is pointed out that the AMSB model has good potential to provide an observable rate of signals since the mixing rate is naturally enhanced by $m_{3/2}/m_{\tilde{\nu}} = \mathcal{O}(4\pi/\alpha)$ while the sneutrino decay rate is small enough on the sizable portion of the phenomenologically allowed parameter space of the model. Our results depicted in Figs. 1 and 2 show that this is indeed the case for the neutrino masses suggested by the atmospheric and solar neutrino data. It is noted also that the same-sign dilepton signals can be used to determine $(m_\nu)_{ee}$ and $(m_\nu)_{\mu\mu}$, providing useful information on the neutrino masses and mixing angles which is complementary to those from neutrino experiments.

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