Implication of Super-Kamiokande Data on R-parity Violation

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Abstract

R-parity violating bilinear (soft) terms in the supersymmetric standard model would be the leading source for nonzero neutrino masses and mixing. We point out that the mixing between neutralinos (charginos) and neutrinos (charged leptons) driven by the bilinear terms take factorized forms, which may enable us to probe the neutrino mixing parameters in a collider. It is then shown that the Super-Kamiokande data on atmospheric neutrinos require all the baryon number violating couplings to be substantially suppressed: $\lambda''_{\text{any}} < 10^{-9}$.

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The supersymmetric standard model allows for the baryon number $B$ and the lepton number $L$ violation. Such $B$ and $L$ violations are usually discarded by imposing the R-parity under which the superpotential is divided into R-parity conserving and violating parts:

\[ h_u^i H_2 Q_i U_i^c + h_d^i H_1 L_i E_i^c + h^c_i H_1 L_i E_i^c + \mu H_1 H_2; \]

\[ \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \lambda''_{ijk} U_i^c D_j^c D_k^c + \epsilon_i \mu L_i H_2, \]

where $i, j, k$ are generation indices. There is, however, no compelling theoretical justification for this requirement. Rather, explicit R-parity violation would well be the origin of nonzero neutrino masses and mixing required by the recent observation of neutrino oscillation in the Super-Kamiokande \[1\]. One of the essential features of $L$ and $R$ violation is the generation of the mixing between neutralinos (charginos) and neutrinos (charged leptons). These mixing come from the presence of $\Delta L = 1$ bilinear soft terms which generate nonzero vacuum expectation values (VEVs) for sneutrino fields unaligned with $\epsilon_i$ \[2\]. As we will see, the neutralino-neutrino and chargino-charged lepton mixing (denoted by $\Theta^N$ and $\Theta^{L,R}$, respectively) are factorized into $R_p$ conserving and violating parts as $\Theta_{ij} = c_i \xi_j$, thereby leading to important phenomenological consequences. If the so-called tree level mass coming from the neutralino-neutrino mixing gives the dominant contribution, one can express the sizes of $\Theta$’s in terms of the atmospheric neutrino mass scale found in the Super-Kamiokande data. Then, the factorization property enables us to relate the neutrino mixing parameters to the $L$ violating sparticle decay rates to charged leptons, which are detectable in a collider. It also turns out that the mixing $\Theta$ combined with the $B$ violation can lead to a fast nucleon decay into a neutrino or charged lepton, unless all the $B$ violating couplings $\lambda''$ are suppressed very strongly.

Let us start our discussion on the mixing $\Theta$ coming from the $L$ violating bilinear terms. The relevant soft supersymmetry (SUSY) breaking terms are

\[ V_{\text{soft}} = m_{H_1}^2 |H_1|^2 + m_{L_i}^2 |L_i|^2 + (m_{L_i H_1}^2 L_i H_1^1 + B H_1 H_2 + B_i L_i H_2 + \text{h.c.}). \]

When the lepton number violating terms are small, $\epsilon_i, \lambda, \lambda' << 1$, as indicated by small neutrino masses advocated by current experiments, it is particularly convenient to use the linear approximation to rotate away the $\epsilon_i$ terms from the superpotential \[1\]:

\[ H_1 \rightarrow H'_1 = H_1 + \epsilon_i L_i, \]
\[ L_i \rightarrow L'_i = L_i - \epsilon_i H_1. \]

which is valid only up to $O(\epsilon_i)$. In this basis, the lepton number is defined as in the supersymmetric limit, the merit of which is that the mixing $\Theta$ can be expressed only in terms of the basis independent quantities $\langle \tilde{\nu}_i' \rangle = \langle \tilde{\nu}_i \rangle - \epsilon_i \langle H_1 \rangle$ as implied by Eq. \[3\]. Upon the minimization of the scalar potential, the lepton number violating soft parameters $B_i', m_{L_i H_2}^{r2}$ induce nonzero VEVs $\langle \tilde{\nu}_i' \rangle \equiv -v_1 \xi_i$ \[3\]:

\[ \xi_i \equiv \frac{B_i' \tan \beta + m_{L_i H_1}^{r2}}{m_{L_i}^2 + M_Z^2 \cos 2\beta/2}, \]
where \( \tan \beta = v_2/v_1 \), \( v_1 = \langle H_1 \rangle \) and \( v = 174 \) GeV. Due to the sneutrino VEVs, neutrinos mix with neutralinos in the \( 7 \times 7 \) mass matrix in the \( (\nu_i; B, W_3, \tilde{H}_1^0, \tilde{H}_2^0) \) basis:

\[
\begin{pmatrix}
    m_{\nu}^{\text{loop}} & m_D \\
    m_D & M_N
\end{pmatrix}
\]

(5)

where \( m_{\nu}^{\text{loop}} \) comes from 1-loop diagrams involving the trilinear couplings \( \lambda, \lambda', \) and \( M_N \) is the usual neutralino mass matrix. Since the \( 3 \times 4 \) matrix \( m_D \),

\[
m_D = \begin{pmatrix}
    M_Z s_W \xi_1 \cos \beta & -M_Z c_W \xi_1 \cos \beta & 0 & 0 \\
    M_Z s_W \xi_2 \cos \beta & -M_Z c_W \xi_2 \cos \beta & 0 & 0 \\
    M_Z s_W \xi_3 \cos \beta & -M_Z c_W \xi_3 \cos \beta & 0 & 0
\end{pmatrix},
\]

(6)

is much smaller than \( M_N \), it is enough to use see-saw formula to find the mixing matrix \( \Theta^N \) between weak eigenstate fields \( (B, W_3, \tilde{H}_1^0, \tilde{H}_2^0) \) and \( \nu_j \):

\[
\Theta^N_{ij} = -(M_N^{-1} m_D^T)_{ij} \equiv \frac{M_Z}{F_N} b_i \xi_j \cos \beta \quad \text{with}
\]

\[
b_1 = -\frac{s_W M_2}{M_1 c_W^2 + M_2 s_W^2}, \quad b_2 = \frac{c_W M_1}{M_1 c_W^2 + M_2 s_W^2},
\]

\[
b_3 = -\sin \beta \frac{M_2}{\mu}, \quad b_4 = \cos \beta \frac{M_Z}{\mu},
\]

\[
F_N = \frac{M_1 M_2}{M_1 c_W^2 + M_2 s_W^2} + \frac{M_2^2}{\mu} \sin 2\beta.
\]

(7)

Here \( s_W = \sin \theta_W, \) etc, and \( \theta_W \) is the weak mixing angle.

In a similar way, the mixing between charged leptons and charginos can be obtained from the following \( 5 \times 5 \) mass matrix for the fields, \( (l_i^\pm; \tilde{W}^\pm, \tilde{H}_{2,1}^\pm) \):

\[
\begin{pmatrix}
    m_C & m_L \\
    m_R & M_C
\end{pmatrix}
\]

with \( m_L = \begin{pmatrix}
    -\sqrt{2} M_W \xi_1 \cos \beta & 0 \\
    -\sqrt{2} M_W \xi_2 \cos \beta & 0 \\
    -\sqrt{2} M_W \xi_3 \cos \beta & 0
\end{pmatrix},
\]

(8)

\[
\begin{pmatrix}
    0 & 0 & 0 \\
    m_e \xi_1 & m_\mu \xi_2 & m_\tau \xi_3
\end{pmatrix},
\]

where \( m_C \) is the diagonal charged lepton mass matrix with elements \( m_{li} \), and \( M_C \) is the usual chargino matrix. While the charged lepton masses \( m_C \) remain untouched to a good approximation, the left-handed charginos \( (\tilde{W}^-, \tilde{H}_1^-) \) mix with the left-handed leptons \( (e, \mu, \tau) \) through the matrix,

\[
\Theta^L_{ij} = -(m_C M_C^{-1})^T_{ij} \equiv c_i^L \xi_j \cos \beta
\]

with \( c_1^L = \sqrt{2} M_W/F_C \), \( c_2^L = -2 \sin \beta \frac{M_2^2}{\mu F_C} \),

\[
F_C = M_2 + M_W^2 \sin 2\beta/\mu.
\]

(9)

The mixing between the charginos \( (\tilde{W}^+, \tilde{H}_2^+) \) and the anti-leptons \( (e^c, \mu^c, \tau^c) \) is given by the matrix,
$$\Theta^R_{ij} = \left(M_C^{-1}(\Theta^L m_C - m_R)\right)_{ij} \equiv c^R_i m_j \xi_j \cos \beta F_C,$$

with \(c^R_i = -\sqrt{2}(1 - \cot \beta \frac{\mu}{M_2}) \frac{M_W M_2}{\mu F_C}\),

\(c^R_2 = \frac{2}{\cos \beta} (1 + \cos^2 \beta \frac{M_2 W}{M_2^2}) \frac{M_2^2}{\mu F_C}.\)  \(10\)

Note that \(\Theta^R\) contains the lepton mass suppression.

The \(L\) violating processes arising from the mixing \(\Theta^R\) show also the factorization property, which may provide a striking connection to the neutrino physics, as we will see. Here, we should mention that the physical quantities involving the mixing of \(\tilde{W}\) or \(\tilde{H}_1\) are not given by \(\Theta^N\) or \(\Theta^L\) alone but by the \(\Theta^N_{3j} - \Theta^L_{2j}\). Recall then that in the basis where \(\epsilon_i\) are not rotated away \(\epsilon\), \(\Theta^N_{3j}\) and \(\Theta^L_{2j}\) contain extra \(\epsilon_j\) which are canceled out in the above combination. Therefore, in either basis, the quantities \(\xi_i\) dictate the properties of the neutralino-neutrino and chargino-charged lepton mixing, and thus the corresponding neutrino masses shown below.

The see-saw reduced mass matrix from \(\Theta^R\) gives so-called the tree level neutrino masses:

$$m_{\nu_{ij}}^{\text{tree}} = \frac{M_2^2}{F_N} \xi_i \xi_j \cos^2 \beta,$$ \(11\)

which, as is well-known, gives a mass to only one neutrino. If the tree mass \(11\) dominates over the loop mass \(m_{\nu}^{\text{loop}}\), the heaviest neutrino mass is given by \(m_{\nu_3} \equiv M_2^2 \xi^2 \cos^2 \beta / F_N\) with \(\xi = \sqrt{\xi_1^2 + \xi_2^2 + \xi_3^2}\). Now the matrix \(\Theta^N\) can be written as

$$\Theta^N_{ij} = \frac{b_i \xi_j}{\xi} \left(\frac{m_{\nu_3}}{F_N}\right)^{1/2},$$ \(12\)

which shows that neutralinos mix only with the heaviest neutrino \(\nu_3\). Here we take \(m_{\nu_3} \simeq \sqrt{\Delta m_{\text{atm}}^2} \simeq 5 \times 10^{-2} \text{ eV}\) as implied by the Super-Kamiokande data \(\|\), and then one finds

$$\xi \cos \beta \simeq 0.7 \times 10^{-6} \frac{F_N}{M_Z} \left(\frac{1}{0.05 \text{eV}}\right)^{1/2}.\) \(13\)

In addition, the mass matrix \(\|\) determines two mixing angles related to the heaviest neutrino \(\|\) and thus gives rise to the oscillation amplitudes with \(\Delta m_{\text{atm}}^2 = \Delta m_{32,31}^2\) as follows:

$$\sin^2 \theta^\text{atm}_{\nu_e} = \frac{4 \xi_2^2}{\xi_2^2 - \xi_1^2},$$

$$\sin^2 \theta^\text{atm}_{\nu_{\mu}} = \frac{4 \xi_3^2}{\xi_3^2 - \xi_2^2}.\) \(14\)

The other mixing angles related to the solar neutrino oscillation can be found only after taking the next leading masses \(m_{\nu}^{\text{loop}}\). The Super-Kamiokande data \(\|\) requires large mixing for \(\nu_{\mu} - \nu_{\tau}\) oscillation, and thus \(\xi \sim \xi_2 \sim \xi_3\). A limit on \(\xi_1^2 / \xi^2\) can come from the CHOOZ \(\bar{\nu}_e\)
disappearance experiment [8]: \( \xi_i^2/\xi_1^2 \lesssim 0.05 \) which is applicable for \( \Delta m^2_{\text{atm}} \gtrsim 2 \times 10^{-3} \text{eV}^2 \).

It is now important to notice that \( \xi_i^2 \)'s appear also in the mixing \( \Theta^{L,R} \), which can lead to L violating sparticle decays. Of particular interest is the lightest sparticle (LSP) decay with same-sign dilepton signal \((l_j l_j (+4 \text{ jets}))\) inside a collider [9]. Typical example is the two body decay, \( \chi^0_1 \rightarrow l_j W \) [9]. Having the decay rate,

\[
\Gamma(l_j) \sim G_F M_{\chi^0_1}^3 |\xi_j|^2 \cos^2 \beta/2\pi, \tag{15}
\]

its decay length is roughly \( \lesssim 1 \text{ mm} \). Then, one will be able to probe the neutrino oscillation amplitudes (14) by counting the same-sign lepton events. For example, in the \( e^+e^- \) collider, taking \( \tau \) in one hemisphere in the direction of incident \( e^- \) and counting the same-sign leptons in the other hemisphere gives the relation

\[
(e\tau) : (\mu\tau) : (\tau\tau) = \xi_1^2 : \xi_2^2 : \xi_3^2. \tag{16}
\]

The CHOOZ result implies that \( \frac{(e\tau)}{(\mu\tau) + (\tau\tau)} \lesssim 0.05 \) independently of \( \Delta m^2_{\text{atm}} \). Otherwise, one can conclude that \( \Delta m^2_{\text{atm}} < 2 \times 10^{-3} \text{eV}^2 \). For \( M_{\chi^0_1} < M_W \), the LSP has only three body decay modes like \( \chi^0_1 \rightarrow l_j + 2\text{ jets} \) through \( W/Z \) boson exchanges which can barely occur inside a collider. To get a qualitative estimation of the decay length, let us take bino \( \tilde{B} \) as the LSP. Its decay rate normalized by the muon decay rate is then given by \( \Gamma_{\chi^0_1}/\Gamma_{\mu} \approx 13.4 b_1^2 (m_{\nu_3}/F_N) (M_{\chi^0_1}/m_{\mu})^5 \) where the numerical factor comes from the final state summation.

If one assumes the unification relation \( M_1 = \frac{5}{3} M_W M_2 \), one gets \( b_1 = 3/8 s_W, F_N = 3 M_1/8 s_W^2 \), and thus

\[
\tau_{\chi^0_1} \sim 0.7 m \left( \frac{0.05 \text{ eV}}{m_{\nu_3}} \right) \left( \frac{M_W}{M_{\chi^0_1}} \right)^4. \tag{17}
\]

Of course, details will depend on the mass parameters in the neutralino sector [10].

Dominance of the tree mass over the loop mass is generally true if there is no alignment in soft masses [11]. But, the loop-to-tree mass ratio is then too small for the loop mass to account for the solar neutrino problem. The situation can be different if supersymmetry breaking process is flavor blind as in the framework of minimal supergravity, or gauge mediated supersymmetry breaking, which is usually required for suppressing flavor changing processes in supersymmetric theories. In this framework, the parameters \( B'_i \) and \( m^2_{L_i H_1} \) are zero at the mediation scale of supersymmetry breaking \( M_m \), but their nonzero values (which are linearly dependent on \( \lambda, \lambda' \) or \( \epsilon_i \) given at \( M_m \)) are generated at the weak scale through renormalization group (RG) evolution [12], and thus both the tree and loop mass are of the radiative origin. In order to get a rough idea how nonzero values of \( \xi_i \) are generated and are related to the R-parity violating couplings in Eq. (1), let us integrate the relevant RG equations [13] in a crude way to get, e.g.,

\[
\xi_i \sim \frac{3\lambda_{333} h_b}{8\pi^2} \ln \frac{M_m}{M_Z} \tag{18}
\]

where \( h_b \) is the b-quark Yukawa coupling. In our basis, rotating away \( \epsilon_i \) terms induces additional contributions to trilinear couplings, in other words, e.g., \( \epsilon_i h_b \) has to be added to
\( \lambda'_{33} \). From Eq. (18), one can see that the smallness of \( \xi_i \sim 10^{-6} \) is due to both the loop factor involving the Yukawa coupling \( h_b \), and small input couplings \( \lambda', \lambda \) or \( \epsilon_i \) which would originate from a horizontal symmetry \[14,15\]. As the tree mass contains the logarithmic enhancement factor, it is typically much larger than the loop mass. It has been recently found that the loop mass can also be enhanced for large tan \( \beta \) to explain the solar neutrino mass scale \[16,17\]. On the other hand, there could be a cancellation in \( \xi_i \) (4) to suppress the tree mass substantially \[13\]. In this case, the tree mass could be comparable to the loop mass, and thus our formulae (14) are invalidated and become complicated functions of the input parameters \( \lambda, \lambda' \) or \( \epsilon_i \).

Another immediate consequence of the mixing \( \Theta \) is proton and neutron decay when combined with the B violating couplings \( \lambda'' \). The bounds from proton longevity are well studied when both \( \lambda' \) and \( \lambda'' \) are present. The strongest bounds come from tree level diagrams with exchanges of a right-handed squark \( \tilde{d}_R \) \[18\]

\[
\lambda'_{ijk} \lambda''_{1lk} \lesssim 10^{-24} \quad j = 1, 2, \quad (19)
\]

for a squark mass of 1 TeV. The diagrams for the proton decay at one loop level involve the other couplings and give rise to mild constraints \[19\]

\[
\lambda' \lambda'' \lesssim 10^{-9}. \quad (20)
\]

New bounds on \( \lambda'' \) combined with the R-parity violating bilinear terms were first derived in Ref. \[20\] by using a mass insertion due to \( L_i H_2 \) terms in the superpotential, which yields

\[
\epsilon_i \lambda''_{12} \lesssim 10^{-21}. \quad (21)
\]

To take into account the effect of the R-parity violating bilinear terms consistently, one has to consider the full mixing between leptons and neutralinos or charginos discussed above, and derive the bounds on the combination \( \lambda'' \Theta \). It turns out then that there are some more new diagrams at tree level contributing to the proton or neutron decay and provide more stringent bounds.

To derive the four-fermion effective Lagrangian responsible for the proton decay to a neutrino resulting from the exchange of right-handed squarks, let us first write down explicitly the interaction Lagrangian between neutralino, quark and right-handed squark

\[
\left\{ \bar{u} \left[ \sqrt{2} g Q_u \tan \theta_W P_L \tilde{B} - h_u^i P_R \tilde{H}_2^0 \right] \tilde{u}_R \\
+ \bar{d} \left[ \sqrt{2} g Q_d \tan \theta_W P_L \tilde{B} - h_d^i P_R \tilde{H}_1^0 \right] \tilde{d}_R \right\} + \text{h.c.}, \quad (22)
\]

where \( P_{L,R} = \frac{1 \mp \gamma_5}{2} \), and \( h_u^i = g m_u^i / \sqrt{2} M_W \sin \beta \), \( h_d^i = g m_d^i / \sqrt{2} M_W \cos \beta \) are the up and down type quark Yukawa couplings respectively. The R-parity violating interactions involving \( \lambda'' \) are

\[
\lambda''_{ijk} \left[ \bar{u}^c_R d_R^j \tilde{d}_R^k + \bar{u}^c_R c_R^j \tilde{u}_R^k + \bar{d}_R^c d_R^j \tilde{d}_R^k \right]. \quad (23)
\]

From the above two equations, one obtains the four-fermion interactions describing the proton decay to a neutrino,
Here we estimate the size of the hadronic coefficient $\eta$ from Ref. [21], where $m_\eta$ is a parameter for the size of the hadronic coefficient $\eta$.

Note that this bound corresponds to those from $\Theta^{ij}_\eta(N_{ij}P_L P) K^{-}$ which would be suppressed comparing with those from $A$ because of the quark mass suppressions in $B$ and $C$ unless $\Delta \bar{m}^2$ is very small and $\tan \beta$ is very large. Let us make some more remarks on the terms in Eq. (25). The $B$ term is nothing but the term arising from the trilinear coupling $\lambda_{j11}$ or $\lambda_{j22}$ defined in the mass basis of the fields. In other words, $\lambda_{j11}$ is the induced coupling of the size $\epsilon_j h_d^d$ or $\Theta^{ij}_\eta h_d^d$. On the other hand, the $A$ and $C$ terms cannot be related to trilinear couplings $\lambda'$ and thus provide the genuine new contributions of the bilinear $R$-parity violation to proton decay. In a naive mass insertion method worked out in Ref. [20], only the $C$ term can be obtained.

From the mixing of the neutralinos with the neutrinos, we obtain from Eq. (25) the decay rate of the proton into $K^+$ and $\nu_3$ neglecting $B$ and $C$ as follows

$$\Gamma = \frac{|A|^2 (m_P^2 - m_{K^+}^2) \tilde{b}_1^2 m_\nu}{32\pi m_P^2 F_N},$$

where $m_P$ denotes the proton mass, etc. From the unobservation of the decay $P \to K\nu$ [22], we obtain a very stringent bound:

$$\lambda''_{112} < 6 \times 10^{-19} x \quad \text{where} \quad x = \left(\frac{0.05 \text{ eV}}{m_\nu}\right)^{1/2} \left(\frac{F_N}{300 \text{ GeV}}\right)^{1/2} \left(\frac{\tilde{m}}{300 \text{ GeV}}\right)^2 \left(\frac{\bar{m}}{\Delta \bar{m}}\right).$$

Note that this bound corresponds to that with $\lambda' \sim 10^{-7}$ in Eq. (14) or $\epsilon_1 \sim 10^{-4}$ in Eq. (21).

We can also find that the chargino–charged lepton mixing $\Theta^{L,R}$ allows for more tree level decay of neutron or proton to charged leptons involving trilinear couplings other than $\lambda''_{112}$. 

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This can be seen by combining the $\lambda''$ couplings (23) and the interaction of chargino, quark and right-handed squark described by

$$V_{ij}h^d_{1i}H^*_1 P_L u^j d^*_R + V_{ij}h^d_{1i}H^*_2 P_L d^j u^*_R + h.c., \quad (28)$$

where $V_{ij}$ are the CKM matrix elements. The relevant four-fermion operator involving charginos takes the form;

$$\frac{\lambda''_{ijk} V_{ip} h^u_{j'}}{m_{1i}} \bar{d}^j P_R d^k H^*_2 P_L d_p + \frac{\lambda''_{ijk} V_{ip} h^d_{j'}}{m_{1i}} \bar{u}^i P_R d^j H^*_1 P_L d_p, \quad (29)$$

from which one finds the corresponding effective couplings for nucleons and mesons;

$$\mathcal{D} \Theta^D_{2j}(t_P L N)K^- + \mathcal{E} \Theta^L_{ij}(t_P L P)\pi^0 + \mathcal{F} \Theta^L_{ij}(t_P L P)K^0,$$

where

$$\mathcal{D} = \frac{\lambda''_{212} V_{21} h^u_3 \eta}{m_{1i}^2} + \frac{\lambda''_{312} V_{31} h^u_3 \eta}{m_{1i}^2},$$

$$\mathcal{E} = \frac{\lambda''_{113} V_{11} h^u_3 \eta}{m_{1i}^2}, \quad \mathcal{F} = \frac{\lambda''_{123} V_{31} h^u_3 \eta}{m_{1i}^2}. \quad (30)$$

As in the previous case, the $\mathcal{E}$ or $\mathcal{F}$ term is again related to the $\lambda'$ coupling of the size $\epsilon_{ij} V_{31} h^d_3$ or $\Theta^L_{ij} V_{31} h^d_3$. Note however that the $\mathcal{D}$ term gives rise to a new tree level contribution involving $\lambda''_{e12}$. Unobservation of the above decays [22] leads to the very severe constraints on $\lambda''_{212}$, $\lambda''_{312}$, $\lambda''_{113}$, and $\lambda''_{123}$ as given in Table I. There, the normalization factors $w, y, z$ are given by

$$w = \frac{\tilde{m}_{b_R}^2 \mu F_C}{(m_{\nu F_N})^{1/2}} \left( \frac{\xi}{\xi_2} \right) \cot \beta$$

$$y, z = \frac{\tilde{m}_{e_R}^2 \mu F_C^2}{(m_{\nu F_N})^{1/2} M^2} \left( \frac{\xi}{\xi_2} \right) \sin 2\beta \quad (31)$$

where $m_{\nu_{eq}}$ is normalized by $5 \times 10^{-2}$ eV and the other mass parameters by 300 GeV as in Eq. (27). We considered the nucleon decay into muon for which $\xi/\xi_2 \sim 1$. It would be worth emphasizing again that one obtains very strong bounds on $\lambda''_{212}$, $\lambda''_{312}$;

$$\lambda''_{212} < 2 \times 10^{-13} y$$

$$\lambda''_{312} < 3 \times 10^{-14} z. \quad (32)$$

which arise from the charged lepton and chargino mixing.

Now taking into account the one loop diagrams with radiative corrections to the $\lambda''$ vertex, we can constrain the remaining $\lambda''$’s [15]. Relative to the tree diagram, one loop diagrams involving $\lambda''_{ijk}$ will be suppressed by the factor $\zeta_{ijk}$, more explicitly, $(A_{ijk}''/\lambda''_{ijk}) = \zeta_{ijk}(A_{tree}/\lambda''_{122})$ where $A_{loop}$ and $A_{tree}$ stand for the loop and tree amplitudes, respectively. The one loop diagrams with the charged higgs boson give
\begin{align*}
\zeta_{ij1} & \approx \frac{g^2 m_i^u m_d^j}{16\pi^2 M_W^2 \sin 2\beta} V_{i2}^* V_{1j}, \\
\zeta_{ij2} & \approx \frac{g^2 m_i^d m_d^j}{16\pi^2 M_W^2 \sin 2\beta} V_{i1}^* V_{1j}.
\end{align*}

\text{(33)}

We derive the above suppression factors using the following up-type quark, down-type quark
and charged higgs interaction Lagrangian:

\begin{align*}
V_{ij} H^+ & \left[ h_i^u \cot \beta \bar{u}_i P_L d_j + h_j^d \tan \beta \bar{u}_i P_R d_j \right] + \text{h.c.}
\end{align*}

\text{(34)}

With this one loop suppression factors the upper bounds on $\lambda''_{ijk}$ are given by

\begin{align*}
\lambda''_{ijk} & < \frac{6 \times 10^{-19}}{\zeta_{ijk}} \left( \frac{5 \times 10^{-2} \text{ eV}}{m_\nu} \right)^{1/2} \\
& \left( \frac{F_N}{300 \text{ GeV}} \right)^{1/2} \left( \frac{\tilde{m}}{300 \text{ GeV}} \right)^2 \left( \frac{\tilde{m}^2}{b_1 \Delta \tilde{m}^2} \right) \sin 2\beta.
\end{align*}

\text{(35)}

From these one-loop suppression factors, we obtain the upper bounds on all $\lambda''_{ijk}$’s listed in

Table I. The normalization factor $x'$ is given by $x' = x \sin 2\beta$. Note that the bounds from

1-loop diagrams with the neutrino (charged lepton)–neutralino (chargino) mixing are more

stringent than any other existing bounds; $\lambda' \lambda'' < 10^{-9}$ \cite{19}, $\lambda_{33} \lambda''_{33} < 10^{-11}$ \cite{20,19}, or the

single bounds $\lambda''_{33} < 10^{-7}$ in gauge mediated SUSY breaking models \cite{23}.

In conclusion, we investigated the consequences of R-parity violation when the bilinear

soft breaking terms are the leading source for the neutrino masses and mixing implied by

the recent Super-Kamiokande data on atmospheric neutrinos. Non-zero sneutrino VEVs

lead to the mixing between neutralinos (charginos) and neutrinos (charged leptons) which

have the factorization property. We show that, due to this property, neutrino mixing angles

can be directly measured in colliders by comparing the branching ratios of the LSP. Taking

into account the particle-particle mixing properly, we are able to find some more new

contributions to proton and neutron decay processes which yield very strong upper bounds

on the B violating couplings.
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TABLES

TABLE I. Constraints on \( \lambda''_{ijk} \) obtained by combining the proton decay limits and the Super-Kamiokande neutrino data. The normalization factors \( x, w, y, z, x' \) can be found in the text.

| Coupling | Constraint |
|----------|------------|
| \( \lambda''_{112} \) | \( 6 \times 10^{-19} x \) |
| \( \lambda''_{113} \) | \( 3 \times 10^{-15} w \) |
| \( \lambda''_{123} \) | \( 7 \times 10^{-15} w \) |
| \( \lambda''_{212} \) | \( 2 \times 10^{-13} y \) |
| \( \lambda''_{212} \) | \( 3 \times 10^{-14} z \) |
| \( \lambda''_{213} \) | \( 1 \times 10^{-10} x' \) |
| \( \lambda''_{223} \) | \( 5 \times 10^{-10} x' \) |
| \( \lambda''_{313} \) | \( 2 \times 10^{-11} x' \) |
| \( \lambda''_{323} \) | \( 1 \times 10^{-10} x' \) |