Lepton flavor violation in supersymmetric $B - L$ extension of the standard model

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**ABSTRACT:** Supersymmetric $B - L$ extension of the Standard Model (SM) is one of the best candidate for physics beyond the SM that accounts for TeV scale seesaw mechanism and provides an attractive solution for the Higgs naturalness problem. We analyze the charged lepton flavor violation (LFV) in this class of models. We show that due to the smallness of Dirac neutrino Yukawa coupling, the decay rates of $l_i \to l_j \gamma$ and $l_i \to 3l_j$, generated by the renormalization group evolution of soft SUSY breaking terms from GUT to seesaw scale, are quite suppressed. Therefore, this model is free from the stringent LFV constraints usually imposed on the supersymmetric seesaw model. We also demonstrate that the right-sneutrino is a long-lived particle and can be pair produced at the LHC through the $B - L$ gauge boson. Then, they decay into same-sign dilepton, with a total cross section of order $\mathcal{O}(1) \text{pb}$. This signal is one of the striking signatures of supersymmetric $B - L$ extension of the SM.

**KEYWORDS:** Supersymmetry, $B - L$, TeV scale seesaw, Lepton flavor violation, same-sign dilepton.
1. Introduction

If neutrinos were massless, the Lepton Flavor (LF) in the charged sector of the SM would be conserved. The observed neutrino oscillations are evidences for neutrino masses, which may entail that lepton flavor is no longer conserved. Nevertheless, LFV is almost forbidden in the SM with massive neutrinos. The processes of charged LFV are suppressed by tiny ratio of neutrino masses to the W-boson mass. For instance, the branching ratio of decay $\mu \rightarrow e\gamma$ is of order $10^{-43}(m_\nu/1\text{eV})^4$, which is far from the experimental reach.

In fact, there is no fundamental reason that implies the conservation of LF in the SM. LF is an accidental symmetry at low energy, and it may be violated beyond the SM. Indeed, several SM extensions, like grand unified field theory (GUT), technicolor, and supersymmetry, indicate the possibility of large LFV. Therefore, a signal of LFV in charged lepton sector would be a clear hint for physics beyond the SM. The present experimental limits [1] are:

$$BR(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11},$$
$$BR(\tau \rightarrow \mu\gamma) < 6.8 \times 10^{-8},$$
$$BR(\tau \rightarrow e\gamma) < 1.1 \times 10^{-7},$$
$$BR(\mu \rightarrow 3e) < 1.0 \times 10^{-12}.$$  (1.1)

The MEG experiment at PSI [2] is expected to reach the limit of $10^{-13}$ for the branching ratio of $\mu \rightarrow e\gamma$ and $\mu \rightarrow 3e$ processes. This will be a very serious test for physics beyond the SM.

Supersymmetry is an attractive candidate for new physics at TeV scale that provides an elegant solution for the SM gauge hierarchy problem and stabilize the SM Higgs mass at
the electroweak scale. In SUSY models, new particles and new interactions are introduced that lead to potentially large LFV rates. Therefore, searches for LFV in charged sector may probe the pattern of SUSY breaking and constrain its origin \cite{3}. Furthermore, seesaw mechanism is an interesting solution to the problem of the small neutrino masses. In what is called type I seesaw mechanism, SM singlets (right-handed neutrinos) with mass of order $\mathcal{O}(10^{14})$ GeV are introduced. It turns out that the combination of these two interesting ideas of SUSY and seesaw implies sizable rates for LFV, even when SUSY breaking terms are assumed to be completely flavor blind. Consequently, the SUSY spectrum should be pushed up to few TeV’s \cite{4}. In this case, there will be no hope to probe SUSY particles at LHC. Also, with very heavy right-handed neutrino, there is no way to test the seesaw mechanism directly at the LHC. Therefore, TeV scale seesaw mechanism was well motivated and has been recently considered as an alternative paradigm \cite{5}.

The TeV scale right-handed neutrino can be naturally implemented in supersymmetric $B-L$ extension of the SM (SUSY $B-L$), which is based on the gauge group $G_{B-L} \equiv SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ \cite{6}. In this model, three SM singlet fermions arise quite naturally due to the $U(1)_{B-L}$ anomaly cancellation conditions. These particles are accounted for right-handed neutrinos, and hence a natural explanation for the seesaw mechanism is obtained \cite{3,4,5}. The masses of these right-handed neutrinos are of order the $B-L$ breaking scale. In SUSY $B-L$ model, the $B-L$ Higgs potential receives large radiative corrections that induce spontaneous $B-L$ symmetry breaking at TeV scale, in analogy to the electroweak symmetry breaking in MSSM \cite{6}. In this case, to fulfill the experimental measurements for the light-neutrino masses, with TeV scale right-handed neutrinos, the Dirac neutrino masses should be order $\mathcal{O}(10^{-4})$ GeV, i.e., they have to be as light as the electron.

In this paper we analyze the LFV in SUSY $B-L$ model. We show that due to the smallness of Dirac neutrino Yukawa couplings, the decay rate of of $l_i \rightarrow l_j \gamma$ and $l_i \rightarrow 3l_j$ are quite suppressed. Hence, the predictions of SUSY $B-L$ for the branching ratio of these processes remain identical to the MSSM ones. Also, we study the pair production of right-sneutrinos at the LHC and show that they are long-lived particles. The decay of these right-sneutrinos leads to a very interesting signal of the same-sign dilepton with possible different lepton flavors. We demonstrate that the cross section of this event is of order $\mathcal{O}(1)$ pb. Therefore, it is quite accessible at the LHC and can be considered as indisputable evidence for SUSY $B-L$ model.

The paper is organized as follows. In section 2 we present the main features of the SUSY $B-L$ extension of the SM. In particular, we analyze the spontaneous $B-L$ breaking at TeV scale by large radiative corrections to the $B-L$ Higgs potential. Section 3 is devoted for the LFV in SUSY $B-L$. We start with the conventional $l_i \rightarrow l_j$ transitions, then we study the same-sign dilepton event which is a clean signal for large right-sneutrino mixing. Finally we give our conclusions in section 4.

2. Supersymmetric $B-L$ extension of the SM

In this section we analyze the minimal supersymmetric version of the $B-L$ extension of
The SM based on the gauge group $G_{B-L} = SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$. This SUSY $B - L$ is a natural extension of the MSSM with three right-handed neutrinos to account for measurements of light neutrino masses. The particle content of the SUSY $B - L$ is the same content as the MSSM with the following extra particles: three chiral right-handed superfields ($N_i$), vector superfield necessary to gauge the $U(1)_{B-L}$ ($Z_{B-L}$), and two chiral SM-singlet Higgs superfields ($\chi_1, \chi_2$ with $B - L$ charges $Y_{B-L} = -2$ and $Y_{B-L} = +2$ respectively). As in MSSM, the introduction of a second Higgs singlet ($\chi_2$) is necessary in order to cancel the $U_{B-L}$ anomalies produced by the fermionic member of the first Higgs ($\chi_1$) superfield.

The interactions between Higgs and matter superfields are described by the superpotential

$$W_{B-L} = W_{MSSM} \left( Y_{ij} L_i H_2 N_j^c + (Y_N)_{ij} N_i^c N_j^c \chi_1 + \mu_1 H_2 + \mu' \chi_1 \chi_2 \right).$$  \hspace{1cm} (2.1)$$

Note that $Y_{B-L}$ for leptons and Higgs are given by $\tilde{Y}$:

$$Y_{B-L}(L) = Y_{B-L}(E) = Y_{B-L}(N) = -1, \quad Y_{B-L}(H_1) = Y_{B-L}(H_2) = 0. \hspace{1cm} (2.2)$$

It also remarkable that due to the $B - L$ gauge symmetry, the $R$-parity violating terms are now forbidden. These terms violate baryon and lepton number explicitly and lead to proton decay at unacceptable rates. On the other hand, the relevant soft SUSY breaking terms, assuming certain universality of soft SUSY breaking terms at GUT scale are in general given by

$$-L_{soft}^{B-L} = -L_{soft}^{MSSM} + \tilde{m}_N^{ij} \tilde{N}_i \tilde{N}_j + m_\chi^2 |\chi_1|^2 + m_\chi^2 |\chi_2|^2$$

$$+ \left[ Y_{ij} \tilde{L}_i \tilde{N}_j^c H_u + Y_{ij} \tilde{N}_i \tilde{N}_j^c \chi_1 + B \mu' \chi_1 \chi_2 + \frac{1}{2} M_{B-L} \tilde{Z}_{B-L} \tilde{Z}_{B-L} + h.c \right], \hspace{1cm} (2.3)$$

where $(Y_N)^{ij} \equiv (Y_N A_N)_{ij}$ is the trilinear associated with Majorana neutrino Yukawa coupling. We now show how the $B - L$ breaking scale can be related to the scale of SUSY breaking, as emphasized in Ref. [8]. The scalar potential for the Higgs fields $H_{1,2}$ and $\chi_{1,2}$ is given by

$$V(H_1, H_2, \chi_1, \chi_2) = \frac{1}{2} g^2 (H_1^* \frac{\tau^a}{2} H_1 + H_2^* \frac{\tau^a}{2} H_2)^2 + \frac{1}{8} g^2 |H_2|^2 - |H_1|^2|^2$$

$$+ \frac{1}{2} g^2 |\chi_1|^2 + m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_3^2 (H_1 H_2 + h.c)$$

$$+ \mu_1^2 |\chi_1|^2 + \mu_2^2 |\chi_2|^2 - \mu_3^2 (\chi_1 \chi_2 + h.c), \hspace{1cm} (2.4)$$

where

$$m_i^2 = m_0^2 + \mu^2, \quad i = 1, 2 \quad m_3^2 = -B \mu, \hspace{1cm} (2.5)$$

$$\mu_i^2 = m_0^2 + \mu^2, \quad i = 1, 2 \quad \mu_3^2 = -B \mu', \hspace{1cm} (2.6)$$

As can be seen from Eq.(2.4), the potential $V(H_1, H_2, \chi_1, \chi_2)$ is factorizable. It can be written as $V(H_1, H_2) + V(\chi_1, \chi_2)$ where $V(H_1, H_2)$ is the usual MSSM scalar potential.
which leads to the radiative electroweak symmetry breaking. As is known, due to the running from GUT to weak scale with large top Yukawa coupling, \( m_2^2 \) receives negative contributions that radiatively breaks the electroweak symmetry. Therefore, we will focus here on the new potential \( V(\chi_1, \chi_2) \) to analyze the possibility of breaking \( B - L \) at TeV scale, through the soft SUSY breaking terms. This potential is given by

\[
V(\chi_1, \chi_2) = \frac{1}{2} g''^2 (|\chi_2|^2 - |\chi_1|^2)^2 + \mu_1^2 |\chi_1|^2 + \mu_2^2 |\chi_2|^2 - \mu_3^2 (\chi_1 \chi_2 + h.c). \tag{2.7}
\]

It should be noted that \( V(\chi_1, \chi_2) \) is quite similar to the MSSM Higgs potential which spontaneously breaks the electroweak symmetry. Therefore it is expected that the \( B - L \) symmetry breaking approach is going to be the same as the well known procedure of electroweak symmetry breaking in MSSM. The minimization of \( V(\chi_1, \chi_2) \) leads to the following condition:

\[
v^2 = (v_1^2 + v_2^2) = \frac{(\mu_1^2 - \mu_2^2) - (\mu_1^2 + \mu_2^2) \cos 2\theta}{2g''^2 \cos 2\theta}, \tag{2.8}
\]

where \( \langle \chi_1 \rangle = v_1 \) and \( \langle \chi_2 \rangle = v_2 \). The angle \( \theta \) is defined as \( \tan \theta = v_1/v_2 \). The minimization conditions also leads to

\[
\sin 2\theta = \frac{2\mu_3^2}{\mu_1^2 + \mu_2^2}. \tag{2.9}
\]

After \( B - L \) breaking, the \( Z_{B-L} \) gauge boson acquires a mass \([5]\): \( M_{Z_{B-L}}^2 = 4g''^2 v^2 \). The high energy experimental searches for an extra neutral gauge boson impose lower bounds on this mass. The stringent constraint on \( U(1)_{B-L} \) obtained from LEP II result, which implies \([15]\)

\[
\frac{M_{Z_{B-L}}}{g''} > 6 \text{TeV}. \tag{2.10}
\]

The discovery potential for \( Z_{B-L} \) at the LHC has been analyzed through its decay into an electron–positron pair \([16]\) and into 3 leptons \([17]\). It was shown that \( \mathcal{O}(1) \) TeV \( Z_{B-L} \) can be easily probed at the LHC with an integrated luminosity of order \( \sim 100 \text{ fb}^{-1} \).

For a given \( M_{Z_{B-L}} \), the minimization condition (2.8) can be used to determine the supersymmetric parameter \( \mu^2 \), up to a sign. One finds

\[
\mu^2 = \frac{m_{\chi_2}^2 - m_{\chi_1}^2 \tan^2 \theta}{\tan^2 \theta - 1} - \frac{1}{4} M_{Z_{B-L}}^2. \tag{2.11}
\]

In order to ensure that the potential \( V(\chi_1, \chi_2) \) is bounded from below, one must require

\[
\mu_1^2 + \mu_2^2 > 2|\mu_3^2|. \tag{2.12}
\]

This is the stability condition for the potential. Also, to avoid that \( \langle \chi_1 \rangle = \langle \chi_2 \rangle = 0 \) be a local minimum we have to require

\[
\mu_1^2 \mu_2^2 < \mu_3^4. \tag{2.13}
\]

It is clear that with positive values of \( \mu_1^2 \) and \( \mu_2^2 \), given in Eq. (2.6), one can not simultaneously fulfill the above conditions. However, as pointed out in Ref. [8], the renormalization
group evolutions of the scalar masses $m_{\chi_1}^2$ and $m_{\chi_2}^2$ of Higgs singlets $\chi_1$ and $\chi_2$ are different. Therefore, at TeV scale the mass $m_{\chi_1}^2$ becomes negative, whereas $m_{\chi_2}^2$ remains positive. In this case, both of the electroweak, $B-L$ and SUSY breakings are linked at scale of $O$(TeV).

In this regards, the observed light-neutrino masses can be obtained if the neutrino Yukawa couplings, $Y_\nu$, are of order $O(10^{-6})$ [3, 7], which are close to the order of magnitude of the electron Yukawa coupling. The LHC discovery for TeV right-handed neutrino in $B-L$ extension of the SM has been studied in Ref. [18]. It was shown that the production rate of the right-handed neutrinos is quite large over a significant range of parameter space. Searching for the right-handed neutrinos is accessible via four lepton channel, which is a very clean signal at LHC, with negligibly small SM background.

With TeV scale right-sneutino, the low-energy sneutrino mass matrix is given by $12 \times 12$ hermitian matrix [19]. However the mixing between left- and right- sneutrinos is quite suppressed since it is proportional to Yukawa coupling $Y_\nu \lesssim O(10^{-6})$. A large mixing between the right-sneutrinos and anti- right-sneutrinos is quite plausible, since it is given in terms the Yukawa $Y_N \sim O(1)$. Therefore, one can focus on the right-sneutrino sector and study the possible oscillation between sneutrino and anti-sneutrino. The right-sneutrino mass matrix in the $(\tilde{N}_c, \tilde{N}_c^*)$ basis can be written as

$$M^2 \simeq \begin{pmatrix} \tilde{m}_N^2 + M_N^2 & -v'_1 (Y_N^d)^* + v'_2 Y_N^d \mu' \\ -v'_1 (Y_N^d)^* + v'_2 Y_N^d \mu' & \tilde{m}_N^2 + M_N^2 \end{pmatrix},$$

(2.14)

As can be seen from the above expression, the off-diagonal elements could be of the same order as the diagonal ones. Therefore, a large mixing can be obtained. In this case, the $(6 \times 6)$ right-sneutrino mass matrix is diagonalized by unitary matrix $X_{\tilde{\nu}}$:

$$X_{\tilde{\nu}} M^2 X_{\tilde{\nu}}^\dagger = M_{\text{diag}}^2.$$  

(2.15)

Hence,

$$\tilde{\nu}_R_i = (X_{\tilde{\nu}})_{ij} \tilde{N}_j, \quad i, j = 1, 2, \ldots, 6.$$  

(2.16)

Finally, we consider the neutral gaugino and Higgsino sector which is going to be modified by the new $B-L$ gaugino and the fermionic partners of the singlet scalar $\chi_{1,2}$. In the weak interaction basis defined by $\psi^0 = (\tilde{B}^0, \tilde{W}^3_1, \tilde{H}^0_1, \tilde{H}^0_2, \tilde{\chi}_1^0, \tilde{\chi}_1^0, \tilde{Z}_{B-L}^0)$, the neutral fermion mass matrix is given by the following $7 \times 7$ matrix [24]:

$$M_n = \begin{pmatrix} M_4 & O \\ O & M_3 \end{pmatrix},$$

(2.17)

where the $M_4$ is the MSSM-type neutralino mass matrix and $M_3$ is the additional neutralino mass matrix with $3 \times 3$:

$$M_3 = \begin{pmatrix} 0 & -\mu' & -2g'' v' \sin \theta \\ -\mu' & 0 & 2g'' v' \cos \theta \\ -2g'' v' \sin \theta + 2g'' v' \cos \theta & 2g'' v' \cos \theta & M_{1/2} \end{pmatrix}.$$  

(2.18)
In case of real mass matrix, one diagonalizes the matrix $\mathcal{M}_n$ with a symmetric mixing matrix $V$ such as

$$V \mathcal{M}_n V^T = \text{diag}(m_{\chi^0_k}), \quad k = 1, \ldots, 7.$$  \hfill (2.19)

In this aspect, the lightest neutralino (LSP) has the following decomposition

$$\chi_1^0 = V_{11} \tilde{B} + V_{12} \tilde{W}^3 + V_{13} \tilde{H}^0_d + V_{14} \tilde{H}^0_u + V_{15} \tilde{\chi}_1 + V_{16} \tilde{\chi}_2 + V_{17} \tilde{Z}_{B-L}.$$  \hfill (2.20)

The LSP is called pure $\tilde{Z}_{B-L}$ if $V_{17} \sim 1$ and $V_{1i} \sim 0$, $i = 1, \ldots, 6$ and pure $\tilde{\chi}_{1(2)}$ if $V_{15(6)} \sim 1$ and all the other coefficients are close to zero \cite{20}. It is worth noting that the MSSM chargino mass matrix remains intact in this type of models, since there is no any new charged fermion have been introduced.

3. LFV in SUSY $B-L$ model

In MSSM, the SUSY contributions to the decay channels of $l_i \rightarrow l_j \gamma$ are dominated by one loop diagrams with neutralino-slepton and chargino-sneutrino exchanges. It turns out that the experimental limit on $\mu \rightarrow e\gamma$ induces stringent constraints on the transitions between first and second generations. Applying the $\mu \rightarrow e\gamma$ constraints on the neutralino contribution leads to the following upper bounds of the slepton mass insertions:

$$(\delta_{LR}^{ij})_{12} \lesssim 10^{-6}, \quad (\delta_{LL}^{ij})_{12} \lesssim 10^{-3}.$$  \hfill (3.1)

Note that due to the $SU(2)_L$ gauge invariance, one gets the following relation between slepton and sneutrino mass insertions: $(\delta_{LR}^{ij})_{ij} \simeq (\delta_{LL}^{ij})_{ij}$. For $(\delta_{LL}^{ij})_{12} \simeq 10^{-3}$, the chargino contribution to $\mu \rightarrow e\gamma$ is automatically below the current experimental limit. These bounds generally impose very stringent constraints on the soft SUSY breaking terms, known as SUSY flavor problem.

It is also worth mentioning that $\mu \rightarrow 3e$ and $\mu \rightarrow e$ conversion in nuclei, $i.e.$, $\mu + N \rightarrow e + N$ are considered as another source of probing possible SUSY effects. The relation between these two processes and $\mu \rightarrow e\gamma$ is, in general, model independent. However, in SUSY framework, where these processes are generated by the photon penguin, $Z$-penguin and box diagrams, one usually finds $\text{BR}(\mu \rightarrow 3e) \sim \text{BR}(\mu \rightarrow e) \sim \mathcal{O}(10^{-3}) \times \text{BR}(\mu \rightarrow e\gamma)$.

In this respect, it seems the present limit on $\mu \rightarrow 3e$ and $\mu \rightarrow e$ conversion are less sensitive than the current bound on $\mu \rightarrow e\gamma$. However, future experiments would reach the limit of $10^{-17}$ for the branching ratio of $\mu \rightarrow e$ conversion and $10^{-16}$ for $\text{BR}(\mu \rightarrow 3e)$, while $\text{BR}(\mu \rightarrow e\gamma)$ may approach $10^{-14}$ at most. These search limits will be powerful tools to probe SUSY at scale of order several TeV. Therefore, in case of negative measurements for all these LFV processes, a very sever constraint is expected to be imposed on the SUSY parameter space.

In minimal supersymmetric seesaw model (which consists of MSSM and right-handed neutrinos), sizable rates for LFV may be obtained through slepton flavor mixing induced radiatively by the large neutrino mixing during the evolution from the grand unification (GUT) scale down to right-handed neutrino scale. In this case, even if universal soft SUSY
breaking parameters are assumed, one finds that the slepton mass matrix receives flavor dependent radiative corrections and the lepton mass insertions are given by

\[(\delta_{LL}^l)_{12} \sim \frac{m_0^2}{m^2} (Y^+_{\nu} Y_{\nu})_{12}, \quad (\delta_{LR}^l)_{12} \sim \frac{m_e A_0}{m^2} (Y^+_{\nu} Y_{\nu})_{12}.\]  

(3.2)

For neutrino Yukawa couplings of order one, the above contribution could enhance the lepton mass insertion significantly. In this case, the upper bound given in Eq.(3.1), in particular \((\delta_{LL}^l)_{12} < 10^{-3}\) is violated unless the slepton masses are quite heavy. It has been explicitly checked that if the neutrino Yukawa coupling is of the form: \(Y_{\nu} = U_{MNS} Y_{\nu}^{\text{diag}}\), then the predicted SUSY contribution to \(\mu \to e\gamma\) is enhanced significantly and exceeds the current experimental limits for most of the parameter space \([9]\). In this respect, the new upper bound \(BR(\mu \to e\gamma) < 10^{-13}\) from MEG experiment might impose a lower bound on the SUSY spectrum of order few TeV, which will be unaccessible at LHC. Therefore, LFV is a serious test for the large scale seesaw mechanism within the SUSY framework.

It is therefore of considerable interest to study TeV scale seesaw, which can easily overcome the LFV problem in SUSY seesaw models. As shown in the previous section, SUSY \(B-L\) extension of the SM is natural framework for implementing TeV scale seesaw. In this class of models, the sever constraints from charged LFV processes are relaxed.

### 3.1 \(l_i \to l_j \gamma\) processes

In SUSY \(B-L\), there are two additional one-loop diagrams contributing to the decay \(l_i \to l_j \gamma\) with \(B-L\) neutralino and chargino exchange, as shown in Fig. 1. In the \(\tilde{Z}_{B-L}\) neutralino contribution, the sleptons are running in the loop. While the chargino diagram involves both left- and right- sneutrinos. It is worth noting that these new contributions are similar to the usual MSSM contribution where neutralino and slepton or chargino and sneutrino are running in the loop. Therefore, the model independent bound on the mass insertions in Eq.(3.1) remains valid for \(x = (m_{\tilde{Z}_{B-L}}/m_l)^2 \approx 1\). Moreover, since the soft SUSY breaking terms are now evolving from GUT to TeV scale, a factor of order \(\mathcal{O}(10)\) is obtained form the \(\ln(M_{\text{GUT}}/M_R)\) in the slepton/sneutrino mass corrections. However, as one can see from Fig. 1, these contributions are proportional to the square of Dirac neutrino Yukawa. Therefore, they are expected to be quite small.

\(^{1}\text{See also Ref.}\ [10]\)
In fact, the $\tilde{Z}_{B-L}$ neutralino contribution is proportional to $B-L$ gauge coupling squared time the mass insertion $(\delta_{LL}^i)_{ab}$, which is proportional to $Y_\nu^2$. As emphasized above, in TeV scale seesaw, the Dirac neutrino Yukawa coupling $Y_\nu$ is of order $O(10^{-6})$. Therefore, the new neutralino amplitude is of order $O(10^{-12})$, which leads to a negligible contribution.

The chargino contribution may be dominated by the mixing between left- and right-sneutrinos. Note that in large scale SUSY seesaw, the right-handed (s)neutrinos are decoupled, hence the chargino contribution is associated with the left-sneutrino only with mass insertion $(\delta_{LL}^i)_{ij}$ correlated with the constrained $(\delta_{LL}^i)_{ij}$. Thus, in SUSY $B-L$, the new chargino contribution is given it terms of $SU(2)$ gauge coupling $g_2$, Dirac neutrino Yukawa coupling $Y_\nu$, and the mass insertion $(\delta_{LR}^i)_{12}$. Nevertheless $(\delta_{LR}^i)_{12}$ is given by

$$
(\delta_{LR}^i)_{12} \simeq (Y_\nu)_{12} \frac{v v^\prime}{m^2}.
$$

For $Y_\nu \sim O(10^{-6})$, the mass insertion $(\delta_{LR}^i)_{ab}$ is of order $O(10^{-7})$, hence the chargino contribution is also quite negligible, of order $O(10^{-14})$. Thus, one can conclude that the LFV associated to $l_i \rightarrow l_j \gamma$ processes, which is generated by RGE from GUT to seesaw scale, is very tiny in SUSY $B-L$ model.

### 3.2 Same-sign and different flavor dilepton signal at the LHC

It is important to note that within MSSM or SUSY seesaw model, another test for LFV at the LHC may be provided by generating final state with different lepton flavors. For example, $\mu^+$ and $e^-$ can be generated at the final state as follows: $q\bar{q} \rightarrow \tilde{l}_i^+ \tilde{l}_i^- \rightarrow \mu^+e^- + 2\chi^0$. However, the cross section of this process is proportional to mass insertion $(\delta_{LL}^i)_{12}$. Therefore, it is typically less than 1 fb, for slepton mass of order 200 GeV. Note that the dilepton associated with this process has opposite sign of electric charges. In MSSM or SUSY seesaw model, same-sign dilepton may be generated only through the gluino and/or squark production followed by several cascade decays.

Now we consider the same-sign and different flavor dilepton production mediated by right-handed neutrino and right sneutrino at the LHC in SUSY $B-L$ model. In particular, we work out the following two processes, shown in Fig. 2.
(i) $pp \to \tilde{Z}_{B-L} \to \tilde{\nu}_R \nu_R \to l_j^- \chi^+ + l_k^- \chi^+ \to l_j^- l_k^- + \text{jets} + \text{missing energy}$.  

(ii) $pp \to \tilde{Z}_{B-L} \to \nu_R, \nu_R \to l_j^- W^+ + l_k^- W^+ \to l_j^- l_k^- + \text{jets}$.  

Here, the following remarks are in order. (i) In the first process the LFV is obtained through the right-sneutrino mixing matrix: $(X_{\tilde{\nu}})_{ij}$. While in the second channel, the neutrino mixing matrix $U_{\nu}$ is responsible for such flavor violation. (ii) Both couplings of $\tilde{\nu}_R - l^- - \chi^+$ and $\nu_R - l^- - W^+$ interactions are suppressed by the mixing between left- and right-sneutrino, which is given by: $\sim m_D/M_R \simeq \mathcal{O}(10^{-7})$. However, the sneutrino coupling has another suppression factor $\sim \mathcal{O}(0.1)$, due to the chargino diagonalizing matrix, $U_\chi$. (iii) One can show that the decay width of right-sneutrino $\Gamma_{\tilde{\nu}_R}$ and right-handed neutrino $\Gamma_{\nu_R}$ are given by

$$
\Gamma_{\tilde{\nu}_R} \sim \frac{1}{8\pi} \frac{|U_{\tilde{\nu}_j} Y_{\nu_k}|^2}{m_{\tilde{\nu}_R}}, \quad \Gamma_{\nu_R} \sim \frac{1}{8\pi} \frac{|Y_{\nu}|^2}{m_{\nu_R}}. \quad (3.4)
$$

It is clear that $\Gamma_{\tilde{\nu}_R} < \Gamma_{\nu_R}$, therefore the right-sneutrino is a long-lived particle more than the right-handed neutrino. In this case, it is expected that the right-sneutrino will have interesting features at the LHC. Accordingly, we will focus our discussion on the case of right-sneutrino pair production.  

From the interaction terms in the SUSY $B - L$ Lagrangian, one finds that the dominant production for the right-sneutrino is through the exchange of $Z_{B-L}$ and its decay is dominated by chargino channel, so that $BR(\tilde{\nu}_R \to l^- \chi^+) \simeq 1$. In general, the amplitude of the process $q\bar{q} \to \tilde{\nu}_R \nu_R \to l^- \chi^+ + l^- \chi^+$, through the s-channel, is given by [21]

$$
\mathcal{M}_{ij} = \sum_k \mathcal{M}_P \frac{i}{q^2 - m_{\tilde{\nu}_k}^2 + i m_{\tilde{\nu}_k} \Gamma_{\tilde{\nu}_k}} (X_{\tilde{\nu}_j})_{kl} \mathcal{M}_D \frac{i}{q^2 - m_{\nu_{R_k}}^2 + i m_{\nu_{R_k}} \Gamma_{\nu_{R_k}}} (X_{\nu_R})_{kj} A_{kl}(q^2) \times \left[ \text{production cross section} \right]
\times \left[ \text{decay branching ratio} \right], \quad (3.5)
$$

where $\mathcal{M}_P$ is the production amplitude for $q\bar{q} \to \tilde{\nu}_{R_k} \nu_{R_k}$ and $\mathcal{M}_D$ is the decay amplitude for $\tilde{\nu}_{R_k}$. As emphasized in Ref.[21], the total cross section $\sigma_{ij} = \sigma(q\bar{q} \to \tilde{\nu}_{R_k} \nu_{R_k} \to l^- \chi^+ + \text{jets} + \text{missing energy})$ can be written as

$$
\sigma_{ij} = \int d^2 q \sum_{kl} (X_{\tilde{\nu}_j}^*)_{kl} (X_{\tilde{\nu}_k})_{ij} A_{kl}(q^2) \times \left[ \text{production cross section} \right]
\times \left[ \text{decay branching ratio} \right], \quad (3.6)
$$

where $A_{kl}(q^2)$ is the product of right-sneutrino propagators:

$$
A_{kl}(q^2) = \frac{i}{q^2 - m_{\tilde{\nu}_k}^2 + i m_{\tilde{\nu}_k} \Gamma_{\tilde{\nu}_k}} \frac{i}{q^2 - m_{\nu_{R_l}}^2 + i m_{\nu_{R_l}} \Gamma_{\nu_{R_l}}} \delta \left( q^2 - m_{\tilde{\nu}_R}^2 \right). \quad (3.7)
$$

Assuming that the off-diagonal elements of the right-sneutrino mass matrix is less than the average right-sneutrino mass, which is quite natural assumption and always valid in standard SUSY breaking mechanisms. Also since the decay width $\Gamma_{\tilde{\nu}_R} \ll m_{\tilde{\nu}_R}$, one can approximate the right-sneutrino propagator as follows [21]:

$$
A_{kl}(q^2) = \frac{1}{1 + i \Delta M_{\tilde{\nu}_R} / \Gamma_{\tilde{\nu}_R}} \frac{\pi}{m_{\tilde{\nu}_R} \Gamma_{\nu_{R_l}}} \delta \left( q^2 - m_{\tilde{\nu}_R}^2 \right). \quad (3.8)
$$
Thus, the total cross section $\sigma_{ij}$ can be written as [22]

$$\sigma_{ij} \approx \frac{|(\Delta m^2_{\tilde{\nu}_R})_{ij}|^2}{m^2_{\tilde{\nu}_R} \Gamma_{\tilde{\nu}_R}^2} \sigma(q\bar{q} \rightarrow \tilde{\nu}_R\tilde{\nu}_R),$$

(3.9)

From Eq.(3.4), the right-sneutrino decay width is of order $\Gamma_{\tilde{\nu}_R} \lesssim \mathcal{O}(10^{-14})$ GeV$^{-1}$. Therefore, with $m_{\tilde{\nu}_R} \sim \mathcal{O}(1)$ TeV and $\Delta M_{\tilde{\nu}}/m_{\tilde{\nu}_R} \sim \mathcal{O}(10^{-2})$, one finds $^2$

$$\sigma_{ij} \approx 10^{10} \sigma(q\bar{q} \rightarrow \tilde{\nu}_R\tilde{\nu}_R),$$

(3.10)

with

$$\sigma(q\bar{q} \rightarrow \tilde{\nu}_R\tilde{\nu}_R) \simeq \frac{g^4_{B-L} m^2_q}{6\pi \left(s^2 - m^2_{Z_{B-L}}\right)^2} \sqrt{1 - \left(\frac{2m_{\tilde{\nu}_R}}{s}\right)^2} \left[1 - \left(\frac{2m_q}{s}\right)^2\right].$$

(3.11)

It is remarkable that for gauge coupling $g'' \sim \mathcal{O}(0.1)$ and $m_{Z_{B-L}} \sim \mathcal{O}(1)$ TeV one finds that the cross section is given by

$$\sigma(q\bar{q} \rightarrow l_i^+ l_j^- + \text{jet + missing energy}) \simeq 10^{-10} \text{ GeV}^{-2} \simeq \mathcal{O}(1) \text{ pb}.$$  

(3.12)

For these values of cross section, the same-sign dilepton signal can be easily probed at the LHC. This event will be a clear hint for sizeable LFV at the LHC, which is more significant than the bounds obtained from the rare decays, $l_i \rightarrow l_j \gamma$. Since the SM background of the same-sign dilepton is negligibly small, the discovery of this process would be undoubted signal for SUSY $B - L$ model.

4. Conclusions

We have analyzed the LFV in supersymmetric $B - L$ extension of the standard model. In this model, $B - L$ symmetry is radiatively broken at TeV scale. Therefore, it is a natural framework for TeV scale seesaw mechanism with Dirac neutrino Yukawa coupling of order $\mathcal{O}(10^{-6})$. We have shown that because of the smallness of Dirac neutrino Yukawa couplings, the decay rates of $l_i \rightarrow l_j \gamma$ and $l_i \rightarrow 3l_j$, generated by the RGE from GUT to seesaw scale, are quite suppressed. In this case, the LFV constraints imposed on this class of models remain as in the MSSM. Also, we have studied another possibility for LFV at the LHC, which associated with the same-sign dilepton, produced through the decay of the long lived pair of right-sneutrinos. We have shown that the total cross section of the process: $q\bar{q} \rightarrow Z_{B-L} \rightarrow \tilde{\nu}_R\tilde{\nu}_R \rightarrow l_j^+ l_k^- + \text{jets + missing energy}$ is of order $\mathcal{O}(1)$ pb. Therefore, it is experimentally accessible at the LHC, with negligibly small SM background. The probe of this signal will provide indisputable evidence for SUSY $B - L$ extension of the SM and also for right sneutrino-anti-sneutrino oscillation.

\[ ^2 \text{More detailed analysis for these processes at LHC, with specific models of supersymmetry breaking, will be considered elsewhere.} \]
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