$b \rightarrow s\gamma$ in the left-right supersymmetric model

Mariana Frank $^1$ and Shuquan Nie $^2$

Department of Physics, Concordia University, 1455 De Maisonneuve Blvd. W. Montreal, Quebec, Canada, H3G 1M8

The rare decay $b \rightarrow s\gamma$ is studied in the left-right supersymmetric model. We give explicit expressions for all the amplitudes associated with the supersymmetric contributions coming from gluinos, charginos and neutralinos in the model to one-loop level. The branching ratio is enhanced significantly compared to the standard model and minimal supersymmetric standard model values by contributions from the right-handed gaugino and squark sector. We give numerical results coming from the leading order contributions. If the only source of flavor violation comes from the CKM matrix, we constrain the scalar fermion-gaugino sector. If intergenerational mixings are allowed in the squark mass matrix, we constrain such supersymmetric sources of flavor violation. The decay $b \rightarrow s\gamma$ sets constraints on the parameters of the model and provides distinguishing signs from other supersymmetric scenarios.

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$^1$Email: mfrank@vax2.concordia.ca
$^2$Email: sxnie@alcor.concordia.ca
1 Introduction

The experimental and theoretical investigation of the inclusive decay $B \to X_s \gamma$ is an important benchmark for physics beyond the Standard Model (SM). Experimentally, the inclusive decay $B \to X_s \gamma$ has been measured at ALEPH [1], BELLE [2] and CLEO [3], with BABAR measurements keenly awaited, giving the following weighted average:

$$BR(B \to X_s \gamma) = (3.23 \pm 0.41) \times 10^{-4}.$$  \hspace{1cm} (1)

This present experimental average is in good agreement with the next-to-leading order predictions in the SM [4]:

$$BR(B \to X_s \gamma)_{SM} = (3.35 \pm 0.30) \times 10^{-4},$$ \hspace{1cm} (2)

But this value still allows a large acceptable range for the inclusive decay [5]:

$$2 \times 10^{-4} < BR(B \to X_s \gamma) < 4.5 \times 10^{-4}. \hspace{1cm} (3)$$

Flavor changing neutral currents (FCNC) are forbidden at the tree level in the SM. The first SM contributions to the process $b \to s \gamma$ appear at one-loop level through the Cabibbo-Kobayashi-Maskawa (CKM) flavor mixing. Despite small uncertainties in the theoretical evaluation of the branching ratio, agreement between experiment and theory is impressive, and this fact is used to set constraints on the parameters for beyond the SM scenarios, such as two-Higgs-Doublet-Models (2HDM) [6], left-right symmetric models (LRM) [7], and minimal supersymmetric standard model (MSSM) [8]. Although attempts have been made to reconcile $b \to s \gamma$ with right-handed $b$-quark decays [9], a complete analysis for a fully left-right supersymmetric model is still lacking.

The Left-Right Supersymmetric (LRSUSY) model is perhaps the most natural extension of the MSSM [10, 11, 12, 13]. Left-right supersymmetry is based on the group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, which would then break spontaneously to $SU(2)_L \times U(1)_Y$.
LRSUSY was originally seen as a natural way to suppress rapid proton decay and has recently received renewed attention for providing small neutrino masses and lepton radiative decays. Besides being a plausible symmetry itself, LRSUSY models have the added attractive features that they can be embedded in a supersymmetric grand unified theory such as \(SO(10)\). Another support for left-right theories is provided by building realistic brane worlds from Type I strings. This involves left-right supersymmetry, with supersymmetry broken either at the string scale \(M_{SUSY} \approx 10^{10-12} \, \text{GeV}\), or at \(M_{SUSY} \approx 1 \, \text{TeV}\), the difference having implications for gauge unification.

In this paper we study all contributions of the LRSUSY model to the branching ratio of \(b \to s\gamma\) at one-loop level. The decay \(b \to s\gamma\) can be mediated by left-handed and right-handed W bosons and charged Higgs bosons as in nonsupersymmetric case, but also by charginos, neutralinos and gluinos. The structure of the LRSUSY provides a significant contributions to the decay \(b \to s\gamma\) from the right-handed squarks and an enlarged gaugino-Higgsino sector with right-handed couplings, which is not as constrained as the right-handed gauge sector in left-right symmetric models. We anticipate that these would contribute a large enhancement of the decay rate and would constrain some of the parameters of the model.

The paper is organized as follows. We describe the structure of the model in Sec. II, with particular emphasis on the gaugino-Higgsino and squark structure. In Sec. III, we give the supersymmetric contributions in LRSUSY to the decay \(b \to s\gamma\). We confront the calculation with experimental results in Sec. IV, where we present the numerical analysis to constrain the parameters of the model for two scenarios: one with CKM flavor mixing only, the other including supersymmetric soft breaking flavor violation. We reach our conclusions in Sec. V.
2 The Model

The LRSUSY electroweak symmetry group, $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, has matter doublets for both left- and right-handed fermions and their corresponding left- and right-handed scalar partners (sleptons and squarks) \cite{12}. In the gauge sector, corresponding to $SU(2)_L$ and $SU(2)_R$, there are triplet gauge bosons $(W^+, W^-, W^0)_L, (W^+, W^-, W^0)_R$, respectively, and a singlet gauge boson $V$ corresponding to $U(1)_{B-L}$, together with their superpartners. The Higgs sector of this model consists of two Higgs bi-doublets, $\Phi_u(\frac{1}{2}, \frac{1}{2}, 0)$ and $\Phi_d(\frac{1}{2}, \frac{1}{2}, 0)$, which are required to give masses to the up and down quarks. The spontaneous symmetry breaking of the group $SU(2)_R \times U(1)_{B-L}$ to the hypercharge symmetry group $U(1)_Y$ is accomplished by giving vacuum expectation values to a pair of Higgs triplet fields $\Delta_L(1, 0, 2)$ and $\Delta_R(0, 1, 2)$, which transform as the adjoint representation of $SU(2)_R$. The choice of two triplets (versus four doublets) is preferred because with this choice a large Majorana mass can be generated (through the see-saw mechanism) for the right-handed neutrino and a small one for the left-handed neutrino \cite{11}. In addition to the triplets $\Delta_{L,R}$, the model must contain two additional triplets, $\delta_L(1, 0, -2)$ and $\delta_R(0, 1, -2)$, with quantum number $B - L = -2$, to insure cancellation of the anomalies which would otherwise occur in the fermionic sector. The superpotential for the LRSUSY model is:

$$W_{LRSUSY} = h_i^Q \tau_2 \Phi_i \tau_2 Q^c + h_i^L L^T \tau_2 \Phi_i \tau_2 L^c + i(h_{LR} L^T \tau_2 \Delta_L L + h_{LR} c^T \tau_2 \Delta_R L^c) + M_{LR} [Tr(\Delta_L \delta_L + \Delta_R \delta_R)] + \mu_{ij} Tr(\tau_2 \Phi_i^T \tau_2 \Phi_j) + W_{NR}$$  \hspace{1cm} (4)

where $W_{NR}$ denotes (possible) non-renormalizable terms arising from higher scale physics or Planck scale effects \cite{17}. The presence of these terms insures that, when the SUSY breaking scale is above $M_{W_R}$, the ground state is R-parity conserving \cite{18}.

The neutral Higgs fields acquire non-zero vacuum expectation values ($VEV's$) through
spontaneous symmetry breaking:

\[
\langle \Delta \rangle_{L,R} = \begin{pmatrix} 0 & 0 \\ v_{L,R} & 0 \end{pmatrix}, \quad \text{and} \quad \langle \Phi \rangle_{u,d} = \begin{pmatrix} \kappa_{u,d} & 0 \\ 0 & \kappa'_{u,d} e^{i\omega} \end{pmatrix}.
\]

\(\langle \Phi \rangle\) causes the mixing of \(W_L\) and \(W_R\) bosons with \(CP\)-violating phase \(\omega\). The non-zero Higgs VEV’s breaks both parity and \(SU(2)_R\). In the first stage of breaking, the right-handed gauge bosons, \(W_R\) and \(Z_R\) acquire masses proportional to \(v_R\) and become much heavier than the SM (left-handed) gauge bosons \(W_L\) and \(Z_L\), which pick up masses proportional to \(\kappa_u\) and \(\kappa_d\) at the second stage of breaking.

In the supersymmetric sector of the model there are six singly-charged charginos, corresponding to \(\tilde{\lambda}_L, \tilde{\lambda}_R, \tilde{\phi}_u, \tilde{\phi}_d, \tilde{\Delta}^\pm_L, \) and \(\tilde{\Delta}^\pm_R\). The model also has eleven neutralinos, corresponding to \(\tilde{\lambda}_Z, \tilde{\lambda}_Z', \tilde{\lambda}_V, \tilde{\phi}^0_{1u}, \tilde{\phi}^0_{2u}, \tilde{\phi}^0_{1d}, \tilde{\phi}^0_{2d}, \tilde{\Delta}^0_L, \tilde{\Delta}^0_R, \tilde{\delta}^0_L, \) and \(\tilde{\delta}^0_R\). Although \(\Delta_L\) is not necessary for symmetry breaking \[13\], and is introduced only for preserving left-right symmetry, both \(\Delta^-_L(\tilde{\Delta}^-_L)\) and its right-handed counterparts \(\Delta^-_R(\tilde{\Delta}^-_R)\) play very important roles in lepton phenomenology of the LRSUSY model. The doubly charged Higgs and Higgsinos do not affect quark phenomenology, but the neutral and singly charged components do, through mixings in the chargino and neutralino mass matrices. We include only the \(\tilde{\Delta}_R\) contribution in the numerical analysis.

The supersymmetric sources of flavor violation in the LRSUSY model come from either the Yukawa potential or the trilinear scalar coupling.

The interaction of fermions with scalar (Higgs) fields has the following form:

\[
\mathcal{L}_Y = h_u \overline{Q}_L \Phi_u Q_R + h_d \overline{Q}_L \Phi_d Q_R + h_\nu \overline{T}_L \Phi_u L_R + h_e \overline{T}_L \Phi_d L_R + H.c.;
\]

\[
\mathcal{L}_M = i h_{LR} (L^T_L C^{-1} \tau_2 \Delta_L L_L + L^T_R C^{-1} \tau_2 \Delta_R L_R) + H.c. \tag{5}
\]

where \(h_u, h_d, h_\nu\) and \(h_e\) are the Yukawa couplings for the up and down quarks and neutrino and electron, respectively, and \(h_{LR}\) is the coupling for the triplet Higgs bosons. LR symmetry requires all \(h\)-matrices to be Hermitean in generation space and \(h_{LR}\) matrix
to be symmetric. We present below the gaugino-Higgsino as well as the sfermion structure of the model, before proceeding with calculation of the branching ratio of $b \to s\gamma$.

## 2.1 Charginos

The terms relevant to the masses of charginos in the Lagrangian are:

$$L_C = -\frac{1}{2}(\psi^+, \psi^-) \begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix} \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + H.c.,$$  

(6)

where $\psi^+ = (-i\lambda_L^+, -i\lambda_R^+, \tilde{\phi}_{u1}^+, \tilde{\phi}_{d1}^+, \tilde{\Delta}_R^T)^T$ and $\psi^- = (-i\lambda_L^-, -i\lambda_R^-, \tilde{\phi}_{u2}^-, \tilde{\phi}_{d2}^-, \tilde{\delta}_R^T)^T$, and:

$$X = \begin{pmatrix} M_L & 0 & g_L \kappa_u & 0 & 0 \\ 0 & M_R & g_R \kappa_u & 0 & 0 \\ 0 & 0 & 0 & -\mu & 0 \\ g_L \kappa_d & g_R \kappa_d & -\mu & 0 & 0 \\ 0 & \sqrt{2} g_{RU} R & 0 & 0 & -\mu \end{pmatrix},$$  

(7)

where we have taken, for simplification, $\mu_{ij} = \mu$. The chargino mass eigenstates $\chi_i$ are obtained by:

$$\chi_i^+ = V_{ij} \psi_j^+, \chi_i^- = U_{ij} \psi_j^-, i, j = 1, \ldots, 5,$$  

(8)

with $V$ and $U$ unitary matrices satisfying:

$$U^* X V^{-1} = M_D,$$  

(9)

where $M_D$ is a diagonal matrix with non-negative entries. Positive square roots of the eigenvalues of $X^TX$ ($XX^T$) will be the diagonal entries of $M_D$ such that:

$$V X^T X V^{-1} = U^* X X^T (U^*)^{-1} = M_D^2,$$  

(10)

The diagonalizing matrices $U^*$ and $V$ are obtained by computing the eigenvectors corresponding to the eigenvalues of $X^TX$ and $XX^T$, respectively.
2.2 Neutralinos

The terms relevant to the masses of neutralinos in the Lagrangian are:

\[ \mathcal{L}_N = -\frac{1}{2} \psi^0_T Y \psi^0 + H.c. , \]  

(11)

where \( \psi^0 = (-i\lambda_{L}^0, -i\lambda_{R}^0, -i\lambda_{V}, \tilde{\phi}_{u1}^0, \tilde{\phi}_{u2}^0, \tilde{\phi}_{d1}^0, \tilde{\phi}_{d2}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0)^T \), and:

\[
Y = \begin{pmatrix}
M_L & 0 & 0 & \frac{g_L \kappa_u}{\sqrt{2}} & 0 & 0 & -\frac{g_R \kappa_d}{\sqrt{2}} & 0 & 0 \\
0 & M_R & 0 & \frac{g_R \kappa_u}{\sqrt{2}} & 0 & 0 & -\frac{g_L \kappa_d}{\sqrt{2}} & -\sqrt{2}g_R v_R & 0 \\
0 & 0 & M_V & 0 & 0 & 0 & 0 & 2\sqrt{2}g_V v_R & 0 \\
\frac{g_L \kappa_u}{\sqrt{2}} & \frac{g_R \kappa_u}{\sqrt{2}} & 0 & 0 & 0 & 0 & 0 & -\mu & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
-\frac{g_L \kappa_d}{\sqrt{2}} & -\frac{g_R \kappa_d}{\sqrt{2}} & 0 & -\sqrt{2}g_R v_R & \sqrt{2}g_V v_R & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix} . \]  

(12)

The mass eigenstates are defined by:

\[ \chi_i^0 = N_{ij} \psi_j^0 \ (i, j = 1, 2, \ldots 9) , \]  

(13)

where \( N \) is a unitary matrix chosen such that:

\[ N^* Y N^{-1} = N_D , \]  

(14)

and \( N_D \) is a diagonal matrix with non-negative entries. To determine \( N \), we take the square of Eq. (14) obtaining:

\[ N Y^T N^{-1} = N_D^2 , \]  

(15)

which is similar to Eq. (10).

We found it convenient to define the neutralino states in terms of the photino and left and right zino states:

\[ \psi^0' = (-i\lambda_{\gamma}, -i\lambda_{Z_L}, -i\lambda_{Z_R}, \tilde{\phi}_{u1}^0, \tilde{\phi}_{u2}^0, \tilde{\phi}_{d1}^0, \tilde{\phi}_{d2}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0)^T , \]  

(16)
with:

\[
\lambda_\gamma = \lambda_L^3 \sin \theta_W + \lambda_R^3 \sin \theta_W + \lambda_V \sqrt{\cos 2\theta_W}
\]

\[
\lambda_{Z_L} = \lambda_L^3 \cos \theta_W - \lambda_R^3 \sin \theta_W \tan \theta_W - \lambda_V \sqrt{\cos 2\theta_W} \tan \theta_W
\]

\[
\lambda_{Z_R} = \lambda_R^3 \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} - \lambda_V \tan \theta_W
\]

Then the mass matrix \( Y \) would be replaced by a matrix \( Y' \) found to be:

\[
Y' = \begin{pmatrix}
 m_\gamma & 0 & 0 & 0 & F & 0 & 0 & -F & -2\sqrt{2}e v_R & 0 \\
 0 & m_{\tilde{Z}_L} & 0 & A_1 & 0 & 0 & A_2 & D & 0 \\
 0 & 0 & m_{\tilde{Z}_R} & E & 0 & 0 & C & B & 0 \\
 F & A_1 & E & 0 & 0 & 0 & -\mu & 0 & 0 \\
 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 & 0 \\
 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 & 0 \\
 -F & A_2 & C & -\mu & 0 & 0 & 0 & 0 & 0 \\
 -2\sqrt{2}e v_R & D & B & 0 & 0 & 0 & 0 & 0 & -\mu \\
 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu & 0
\end{pmatrix}
\]

(18)

with:

\[
m_\gamma = M_V \cos 2\theta_W + (M_L + M_R) \sin^2 \theta_W,
\]

\[
m_{\tilde{Z}_L} = M_V \cos 2\theta_W \tan^2 \theta_W + M_L \cos^2 \theta_W + M_R \sin^2 \theta_W \tan^2 \theta_W,
\]

\[
m_{\tilde{Z}_R} = M_V \tan^2 \theta_W + M_R(1 - \tan^2 \theta_W),
\]

\[
A_1 = \frac{g}{\sqrt{2} \cos \theta_W} (\kappa_u \cos^2 \theta_W - \kappa_d \sin \theta_W),
\]

\[
A_2 = \frac{g}{\sqrt{2} \cos \theta_W} (\kappa_d \cos^2 \theta_W - \kappa_u \sin \theta_W),
\]

\[
B = -\sqrt{2}g \frac{1 - 2 \tan^2 \theta_W}{\sqrt{1 - \tan^2 \theta_W}} v_R,
\]

\[
C = -\frac{g \sqrt{\cos 2\theta_W}}{\sqrt{2} \cos \theta_W} \kappa_u,
\]

\[
D = -\frac{\sqrt{2} g \sqrt{\cos 2\theta_W}}{\cos \theta_W} v_R,
\]

\[
E = \frac{g \sqrt{\cos 2\theta_W}}{\sqrt{2} \cos \theta_W} \kappa_d,
\]

\[
F = \frac{e}{\sqrt{2}} (\kappa_u + \kappa_d),
\]

(19)
The unitary matrix \( N \) would be replaced by a new matrix: \( N' \) given by:

\[
\begin{align*}
N'_{j1} &= N_{j1} \sin \theta_W + N_{j2} \sin \theta_W + N_{j3} \sqrt{\cos 2\theta_W} \\
N'_{j2} &= N_{j1} \cos \theta_W - N_{j2} \sin \theta_W \tan \theta_W - N_{j3} \sqrt{\cos 2\theta_W} \tan \theta_W \\
N'_{j3} &= N_{j2} \sqrt{\cos 2\theta_W} - N_{j3} \tan \theta_W \\
N'_{jk} &= N_{jk}, \quad (k = 4, 5, \ldots 9).
\end{align*}
\]

Similarly \( N \) can be expressed in terms of \( N' \) by:

\[
\begin{align*}
N_{j1} &= N'_{j1} \sin \theta_W + N'_{j2} \cos \theta_W \\
N_{j2} &= N'_{j1} \sin \theta_W - N'_{j2} \sin \theta_W \tan \theta_W - N'_{j3} \sqrt{\cos 2\theta_W} \tan \theta_W \\
N_{j3} &= N'_{j1} \sqrt{\cos 2\theta_W} - N'_{j2} \tan \theta_W \sqrt{\cos 2\theta_W} - N'_{j3} \tan \theta_W \\
N_{jk} &= N'_{jk}, \quad (k = 4, 5, \ldots 9).
\end{align*}
\]

### 2.3 Squarks

In the interaction basis, \((\tilde{q}^i_L, \tilde{q}^i_R)\), the squared-mass matrix for a squark of flavor \( f \) has the form:

\[
M^2_f = \begin{pmatrix}
(m^2_{f,LL} + F_{f,LL} + D_{f,LL} (m^2_{f,LR})^2 + F_{f,LR}) + F_{f,RR} + D_{f,RR} \\
(m^2_{f,RR})^2 + F_{f,RR} + D_{f,RR}
\end{pmatrix}.
\]

The F-terms are diagonal in the flavor space, \( F_{f,LL,RR} = m^2_f \), \( F_{d LR} = -\mu(m_d \tan \beta) \mathbf{1}_{ij} \), \( F_{u LR} = -\mu(m_u \cot \beta) \mathbf{1}_{ij} \). The D-terms are also flavor-diagonal:

\[
\begin{align*}
D_{f,LL} &= M^2_Z \cos 2\beta (T_f^3 - Q_f \sin^2 \theta_W) \mathbf{1} \\
D_{f,RR} &= M^2_Z \cos 2\beta Q_f \sin^2 \theta_W \mathbf{1}
\end{align*}
\]

The term \((m^2_{f,LL,RR})_{ij} = m^2_{Q_{L,R}} \delta_{ij}, m^2_f L R = A_f m_f \). To reduce the number of free parameters, we consider the parameters to be universal, with: \((m^2_{Q_{L,R}})_{ij} = m^2_0 \delta_{ij}, A_{d,ij} = A \delta_{ij} \) and \( A_{u,ij} = A \delta_{ij} \).
The squared-mass matrix for U-type squarks reduces to:

$$\mathcal{M}_{U_{k}}^{2} = \begin{pmatrix} m_{0}^{2} + M_{Z}^{2} (T_{u}^{3} - Q_{u} \sin^{2} \theta_{W}) \cos 2\beta & m_{u_{k}} (A - \mu \cot \beta) \\ m_{u_{k}} (A - \mu \cot \beta) & m_{0}^{2} + M_{Z}^{2} Q_{u} \sin^{2} \theta_{W} \cos 2\beta \end{pmatrix}.$$  \hspace{1cm} (24)

with the diagonal F-terms absorbed into the $m_{0}^{2}$. For D-type squarks:

$$\mathcal{M}_{D_{k}}^{2} = \begin{pmatrix} m_{0}^{2} + M_{Z}^{2} (T_{d}^{3} - Q_{d} \sin^{2} \theta_{W}) \cos 2\beta & m_{d_{k}} (A - \mu \tan \beta) \\ m_{d_{k}} (A - \mu \tan \beta) & m_{0}^{2} + M_{Z}^{2} Q_{d} \sin^{2} \theta_{W} \cos 2\beta \end{pmatrix}.$$  \hspace{1cm} (25)

The corresponding mass eigenstates are defined as:

$$\tilde{q}_{L,R} = \Gamma_{Q}^{\dagger} q_{L,R}$$  \hspace{1cm} (26)

where $\Gamma_{Q}^{\dagger} q_{L,R}$ are $6 \times 3$ mixing matrices. In the universal case, there is no intergenerational mixings for squarks and the only source of flavor mixing comes from the CKM matrix.

We will analyse this case first. Next we will look at the case in which mixing in the squark sector is permitted and consider the effect of intergenerational mixings on the rate of the process $b \rightarrow s \gamma$. As it is generally done in the mass insertion approximation method [19], where the off-diagonal squark mass matrix elements are assumed to be small and their higher orders can be neglected, we use normalized parameters:

$$\delta_{d,LL,ij} = \frac{(m_{d,LL}^{2})_{ij}}{m_{0}^{2}}, \quad \delta_{d,RR,ij} = \frac{(m_{d,RR}^{2})_{ij}}{m_{0}^{2}},$$

$$\delta_{d,LR,ij} = \frac{(m_{d,LR}^{2})_{ij}}{m_{0}^{2}}, \quad \delta_{d,RL,ij} = \frac{(m_{d,RL}^{2})_{ij}}{m_{0}^{2}}.$$  \hspace{1cm} (27)

In our analysis, we use the mass eigenstate formalism, which is valid no matter how large the intergenerational mixings are. We assume significant mixing between the second and third generations in the down-squarks mass matrix only.

### 3 Supersymmetric contributions to $b \rightarrow s \gamma$

The low-energy effective Hamiltonian responsible for the B meson decay rates at the scale $\mu$ can be written as:

$$\mathcal{H}_{\text{eff}} = -\frac{4G_{F}}{\sqrt{2}} K_{tb} K_{ts}^{*} \sum_{i} C_{i}(\mu) Q_{i}(\mu).$$  \hspace{1cm} (28)
Figure 1: The Feynman diagrams contributing to the decay $b \rightarrow s\gamma$. The outgoing photon line can be attached in all possible ways.

The operators relevant to the process $b \rightarrow s\gamma$ in LRSUSY are:

\[
Q_7 = \frac{e}{16\pi^2} m_b(\mu) \bar{s}\sigma_{\mu\nu} P_R b F^{\mu\nu},
\]

\[
Q'_7 = \frac{e}{16\pi^2} m_b(\mu) \bar{s}\sigma_{\mu\nu} P_L b F^{\mu\nu},
\]

\[
Q_8 = \frac{gs}{16\pi^2} m_b(\mu) \bar{s}\sigma_{\mu\nu} G_\mu^{\alpha T a} P_R b,
\]

\[
Q'_8 = \frac{gs}{16\pi^2} m_b(\mu) \bar{s}\sigma_{\mu\nu} G_\mu^{\alpha T a} P_L b
\]

and the Wilson coefficients $C'_{7,8}$ are initially evaluated at the electroweak or soft supersymmetry breaking scale, then evolved down to the scale $\mu$. The Feynman diagrams contributing to this decay in LRSUSY are illustrated in Fig. 1.

The matrix elements responsible for the $b \rightarrow s\gamma$ decay acquire the following contributions from the supersymmetric sector of the model. For $b_L$ decay:

\[
M_{\gamma R} = A_{\tilde{g}}^R + A_{\tilde{\chi}^0}^R + A_{\tilde{\chi}^-}^R
\]

with the gluino, chargino and neutralino contributions given by:

\[
A_{\tilde{g}}^R = -\frac{\alpha_s}{\sqrt{2}G_F} Q_d C(R) \sum_{k=1}^6 \frac{1}{m_{\tilde{d}_k}^2} \{ \Gamma_{DL}^{k b} \Gamma_{DL}^{k s} F_2(x_{\tilde{g}\tilde{d}_k}) - \frac{m_{\tilde{g}}}{m_b} \Gamma_{DR}^{k b} \Gamma_{DL}^{k s} F_4(x_{\tilde{g}\tilde{d}_k}) \}
\]

\[
A_{\tilde{\chi}^-}^R = -\frac{\alpha_w}{\sqrt{2}G_F} \sum_{j=1}^5 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_j}^2} \{ (G_{UL}^{j kb} - H_{UL}^{j kb}) (G_{UL}^{s jks} - H_{UL}^{s jks}) [F_1(x_{\tilde{\chi}^- j\tilde{u}_k}) + Q_u F_2(x_{\tilde{\chi}^- j\tilde{u}_k})] 
\]

\[
+ \frac{m_{\tilde{\chi}^-}}{m_b} (G_{UR}^{j kb} - H_{UL}^{j kb}) (G_{UL}^{s jks} - H_{UL}^{s jks}) [F_3(x_{\tilde{\chi}^- j\tilde{u}_k}) + Q_u F_4(x_{\tilde{\chi}^- j\tilde{u}_k})] \}
\]

\[
A_{\tilde{\chi}^0}^R = -\frac{\alpha_w}{\sqrt{2}G_F} Q_d \sum_{j=1}^5 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_j}^2} \{ (\sqrt{2}G_{ODL}^{j kb} - H_{0DR}^{j kb}) (\sqrt{2}G_{ODL}^{s jks} - H_{0DR}^{s jks}) F_2(x_{\tilde{\chi}^0 j\tilde{d}_k}) 
\]

\[
+ \frac{m_{\tilde{\chi}^0}}{m_b} (\sqrt{2}G_{0DR}^{j kb} - H_{0DR}^{j kb}) (\sqrt{2}G_{ODL}^{s jks} - H_{0DR}^{s jks}) F_4(x_{\tilde{\chi}^0 j\tilde{d}_k}) \}
\]
and, for the decay of $b_R$:

$$M_{\gamma L} = A_{\tilde{g}}^L + A_{\tilde{\chi}^0}^L + A_{\tilde{\chi}}^L$$  \hspace{1cm} (34)$$

again, with the following gluino, chargino and neutralino contributions:

$$A_{\tilde{g}}^L = -\frac{\pi \alpha_s}{\sqrt{2} G_F} Q_d C(R) \sum_{k=1}^{6} \frac{1}{m^2_{\tilde{d}_k}} \{ \Gamma_{DR}^{\tilde{g}} \Gamma_{DR}^{\tilde{g} s k} F_2(x_{\tilde{g} \tilde{d}_k}) - m_{\tilde{g}} \Gamma_{DL}^{\tilde{g} s k} F_4(x_{\tilde{g} \tilde{d}_k}) \}$$  \hspace{1cm} (35)$$

$$A_{\tilde{\chi}^-}^L = -\frac{\pi \alpha_w}{\sqrt{2} G_F} \sum_{j=1}^{5} \sum_{k=1}^{6} \frac{1}{m^2_{\tilde{\chi}^- k}} \{ (G_{UR}^{j k}) - H_{UL}^{j k}) (G_{UR}^{j k} + H_{UL}^{j k}) [F_1(x_{\tilde{\chi}^- \tilde{u}_k}) + Q_u F_2(x_{\tilde{\chi}^- \tilde{u}_k})]$$

$$+ \frac{m_{\tilde{\chi}^-}}{m_b} (G_{UL}^{j k} - H_{UL}^{j k}) (G_{UL}^{j k} + H_{UL}^{j k}) [F_3(x_{\tilde{\chi}^- \tilde{u}_k}) + Q_u F_4(x_{\tilde{\chi}^- \tilde{u}_k})] \}$$  \hspace{1cm} (36)$$

$$A_{\tilde{\chi}^0}^L = -\frac{\pi \alpha_w}{\sqrt{2} G_F} Q_d \sum_{j=1}^{9} \sum_{k=1}^{6} \frac{1}{m^2_{\tilde{\chi}^0 k}} \{ (\sqrt{2} G_{0 DR}^{j k} - H_{0 DL}^{j k}) (\sqrt{2} G_{0 DR}^{j k} - H_{0 DL}^{j k}) [F_2(x_{\tilde{\chi}^0 \tilde{d}_k})$$

$$+ \frac{m_{\tilde{\chi}^0}}{m_b} (\sqrt{2} G_{0 DL}^{j k} - H_{0 DR}^{j k}) (\sqrt{2} G_{0 DL}^{j k} - H_{0 DR}^{j k}) F_4(x_{\tilde{\chi}^0 \tilde{d}_k}) \}$$  \hspace{1cm} (37)$$

where vertex mixing matrices $G$, $H$, $G_0$ and $H_0$ are defined in the Appendix. The convention $x_{ab} = m^2_a/m^2_b$ is used. $C(R) = 4/3$ is the quadratic Casimir operator of the fundamental representation of $SU(3)_C$.

In order to compare the results obtained with experimental branching ratios, QCD corrections must be taken into account. We assume below the SM renormalization group evolution pattern; supersymmetric estimates exist for the gluino contributions only [22].

There is no mixing between left and right-handed contributions.

$$A^\gamma(m_b) = \eta^{-16/23} \{ A^\gamma(M_W) + A^\gamma_0 \left[ \frac{116}{135} (\eta^{28/23} - 1) + \frac{58}{189} (\eta^{28/23} - 1) \right] \},$$  \hspace{1cm} (38)$$

where $\eta = \alpha_s(m_b)/\alpha_s(M_W)$ and $A^\gamma_0 = \frac{\pi \alpha_w}{2 \sqrt{2} G_F M_W} \frac{1}{M_W}$. We choose the renormalization scale to be $\mu = m_b = 4.2$ GeV.

The inclusive decay width for the process $b \rightarrow s\gamma$ is given by:

$$\Gamma(b \rightarrow s\gamma) = \frac{m^5_b G^2_F |K_{ib}K^*_{ls}|^2}{32\pi^4} \left( \hat{M}_{\gamma L}^2 + \hat{M}_{\gamma R}^2 \right) \left( \eta^2_{\gamma L} + \eta^2_{\gamma R} \right),$$  \hspace{1cm} (39)$$

where the hat means evolving down to the decay scale $\mu = m_b$. The branching ratio can
be expressed as

$$BR(b \to s\gamma) = \frac{\Gamma(b \to s\gamma)}{\Gamma_{SL}} BR_{SL},$$

(40)

where the semileptonic branching ratio $$BR_{SL} = BR(b \to c e\bar{\nu}) = (10.49 \pm 0.46)\%$$ and:

$$\Gamma_{SL} = \frac{m_{\tilde{e}}^2 G_F^2 |K_{cb}|^2}{192\pi^3} g(z),$$

(41)

where $$z = m_c^2/m_b^2$$ and $$g(z) = 1 - 8z + 8z^3 - z^4 - 12z^2\log z$$.

4 Numerical results

We are interested in analysing the case in which the supersymmetric partners have masses around the weak scale, so we will assume relatively light superpartner masses. We diagonalize the neutralino and chargino mass matrices numerically and we require in all calculations that the masses of gluinos, charginos, neutralinos and squarks be above their experimental bounds. There are some extra constraints in the non-supersymmetric sector of the theory, requiring the FCNC Higgs boson $$\Phi_d$$ to be heavy, but no such constraints exist in the Higgsino sector [23]. We choose the gluino mass $$m_{\tilde{g}} = 300$$ GeV, and left-handed gaugino masses of $$M_L = 500$$ GeV. The mass of the lightest bottom squark will be in the 150 – 200 GeV range. We include in our numerical estimates the non-supersymmetric LRM results, and for this we constrain the lightest Higgs mass to be 115 GeV [24]. We analyze our results for low and moderate values of $$\tan \beta$$, although we study the dependence of the branching ratio on $$\tan \beta$$. We also investigate the dependence of the branching ratio on both positive and negative values of the $$\mu$$ parameter.

As a first step, we assume the only source of flavor violation to come from the CKM matrix. This scenario is related to the minimal flavor violation scenario in supergravity. This restricted possibility of flavor violation will set important constraints on the parameter space of LRSUSY.
We then allow, in the second stage of our investigation, for new sources of flavor violation coming from the soft breaking terms. In MSSM, this scenario is known as the unconstrained MSSM and there the gluino contribution dominates. This is not so in LRSUSY, where the chargino contribution is important for low to intermediate values of $M_R$, the right-handed gaugino mass parameter. We restrict all allowable LL, LR, RL and RR sflavor mixing, assuming them to be dominated by mixings between the second and third squark family.

We now proceed to discuss both these scenarios in turn.

4.1 The constrained LRSUSY

By the constrained LRSUSY model, we mean the scenario in which the only source of flavor violation comes from the quark sector, through the CKM matrix, which we assume to be the same for both the left and right handed sectors.

Before any meaningful numerical results be obtained, explicit values for the parameters in the model must be specified. There are many parameters in the model such that it is hard, if not impossible, to get an illustrative presentation of calculation results. If LRSUSY is embedded in a supersymmetric grand unification theory such as $SO(10)$, there exist some relationships among the parameters at the unification scale $M_{GUT}$. We can generally choose specific values for much less parameters at mass scales $\mu = M_{GUT}$, then use renormalization group equations to run them down to the low energy scale which is relevant to phenomenology. But, for maintaining both simplicity and generality, we can present an analysis in which LRSUSY is not embedded into another group. Then we can choose all parameters as independently free parameters, with the numerical results confronting with experiments directly.

To make the results tractable, we assume all trilinear scalar couplings in the soft supersymmetry breaking Lagrangian as $A_{ij} = A\delta_{ij}$ and $\mu_{ij} = \mu\delta_{ij}$, and we fix $A$ to be
50 GeV in all the analysis. We also set a common mass parameter for all the squarks \( M_{0 UL} = M_{0 UR} = M_{0 DL} = M_{0 DR} = m_0 \). The general range of these parameters is as discussed in the previous paragraphs. We also take \( K_{CKM}^L = K_{CKM}^R \). This choice is conservative, and much larger values of mixing matrix elements are allowed in scenarios that attempt to explain the decay properties of the \( b \) quark as being saturated by the right-handed \( b \) [9]. Our choice does not favor one handedness over the other, and has the added advantage that no new mixing angles are introduced in the quark matrices.

We investigate first the dependence of the branching ratio on the values of \( \tan \beta \) in Fig. 2. The solid line represents the supersymmetric contribution, the dashed the total contribution (including the SM and LRM). All the graphs show the experimental bounds as horizontal lines in the figures. The universal squark mass is around 200 GeV, and the chargino and neutralino are light. The choice of parameters puts stringent restrictions on the allowed values for \( \tan \beta \): either very low (\( \tan \beta = 2-4 \)), or intermediate in a very small range (\( \tan \beta = 12-14 \)) values are allowed. As \( \tan \beta \) becomes large, the branching ration is almost linearly proportional to \( \tan \beta \). For larger values of \( \tan \beta \) the branching ratio will exceed the acceptable range easily. In our analysis, larger values of \( \tan \beta \) are allowed only for a heavier supersymmetric mass spectra. For \( \text{sign}(\mu) < 0 \), the range of acceptable intermediate values of \( \tan \beta \) increases. For example, if \( \mu = -100 \) GeV, the larger range \( \tan \beta = 22 - 33 \) is allowed. We note that constraints on supersymmetry with large \( \tan \beta \) were studied in Ref. [25] where \( \tan \beta \)-enhanced chargino and charged Higgs contributions were resummed to all order in perturbation theory.

As a general feature of the LRSUSY branching ratio, in a large region of parameter space, the chargino contribution is comparable to the gluino, while the neutralino contribution is always smaller. We investigate the dependence of the branching ratio on the gluino mass, for a light squark scenario. The chargino and neutralino masses are light and \( \mu/M_{L,R} \sim \mathcal{O}(1) \), a scenario favored by recent analyses of the anomalous magnetic
moment of the muon [26]. We present the results in Fig. 3. The gluino is constrained to be heavier than 300 GeV, albeit for a very light supersymmetric spectrum, close to experimental limits.

In the next two figures, we investigate the dependence of the branching ratio of $b \to s\gamma$ to the sign and magnitude of the Higgsino mixing parameter $\mu$. There have been indications that the new accurate measurement of the anomalous magnetic moment of the muon restrict the $\mu$ parameter to be positive, while $b \to s\gamma$ favors a negative sign. For a light squark-gaugino scenario, and low $\tan \beta$, the bound on $b \to s\gamma$ is satisfied for either sign of the $\mu$ parameter. In Fig. 4 one could see that a restricted region of intermediate values for $\mu$, with $\text{sign}(\mu) > 0$ is allowed by the experimental constraints on $b \to s\gamma$, in the 225-325 GeV region. The parameter space is less restrictive for $\mu$ negative, to -175 GeV. Note that for $\mu \to 0$ the branching ratio drops outside the allowed range. This phenomenon occurs because the mixing term obtained from flipping chirality on
Figure 3: Supersymmetric contributions to BR($b \to s\gamma$) as a function of the mass of the gluino $m_{\tilde{g}}$, obtained when $\tan \beta = 5$, $\mu = 100$ GeV, $M_L = M_R = 500$ GeV and $m_0 = 100$ GeV. The full contributions is also shown(dashed). The range of acceptable values of branching ratios is given.

Figure 4: Supersymmetric contributions to BR($b \to s\gamma$) as a function of $\mu$, obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV and $m_0 = 100$ GeV. The full contributions is also shown(dashed). The range of acceptable values of branching ratios is given.
Figure 5: Supersymmetric contributions to BR($b \rightarrow s\gamma$) as a function of $\mu$, obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV and $m_0 = 150$ GeV. The full contributions is also shown (dashed). The range of acceptable values of branching ratios is given.

the gaugino leg decouples. We reject such small values of the $\mu$ parameter because the chargino and neutralino masses are smaller than the existing experimental bounds.

The branching ratio for $b \rightarrow s\gamma$ is sensitive to the universal scalar mass $m_0$ in the region of small masses only. For $m_0 \geq 400$ GeV, the branching ratio reaches its QCD-corrected value and is stable against further variations in the scalar mass. In this scenario, the neutralinos and gluinos are light and $\tan \beta = 5$. This situation is not unlike the dependence of the SM contribution on the $t$ quark mass [4, 8]. This dependence is shown in Fig. 6.

In all the previous figures we set the left and right handed gaugino masses to the same value. This allowed a large contribution to the decay ratio of $b \rightarrow s\gamma$ to come from the right-handed sector. We investigate in Fig. 7 the dependence of the branching ratio on the gaugino mass. As opposed to the right-handed gauge sector, the restriction on the right handed gaugino scale is not as severe. There exist scenarios in which the right
Figure 6: Supersymmetric contributions to BR($b \rightarrow s\gamma$) as a function of $m_0$, obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV and $\mu = 100$ GeV. The full contributions is also shown (dashed). The range of acceptable values of branching ratios is given.

handed symmetry is broken at the same scale as supersymmetry; we expect in those cases to have approximately $M_L = M_R$ [27]. For light squarks, Higgsinos and gluinos, the gaugino mass must be heavy, in the 600-800 GeV range.

4.2 The unconstrained LRSUSY

When supersymmetry is softly broken, there is no reason to expect that the soft parameters would be flavor blind, or that they would violate flavor in the same way as in the SM. Yukawa couplings generally form a matrix in the generation space, and the off-diagonal elements will lead naturally to flavor changing radiative decays. Neutrino oscillations, in particular, indicate strong flavor mixing between the second and third neutrino generations, and various analyses have been carried out assuming the same for the charged sleptons. In the quark/squark sector, the kaon system strongly limits mixings between the first and the second generations; but constraints for the third generation are much
Figure 7: Supersymmetric contributions to BR($b \rightarrow s\gamma$) as a function of $M_R$, obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $m_0 = 100$ GeV $\mu = 100$ GeV and $M_L = M_R$ is assumed. The full contributions is also shown (dashed). The range of acceptable values of branching ratios is given.

Weaker, and expected to come from $b \rightarrow s\gamma$. The unconstrained LRSUSY model, similar to the unconstrained MSSM, allows for new sources of flavor violation between the second and third families only, both chirality conserving (LL and RR) and chirality flipping (LR and RL). We will assume that intergenerational mixing occurs in the down squark mass matrix only and that the up type squark mass matrix is diagonal.

With the definition of the mass insertion as in Eq. (27), we can investigate the effect of intergenerational mixing on the $b \rightarrow s\gamma$ decays. In the MSSM, the branching ratio is dominated in this case by the gluino diagram, in particular by the chirality flip part of the gluino contribution, due to the $\alpha_s/\alpha$ and $m_{\tilde{g}}/m_b$ enhancements, respectively. In this case only the gluino scenario is analysed in the MSSM, and found to be dominated by $\delta_{23}^{RL}$ [22]. In LRSUSY, the situation is different: the chargino graph contribution is comparable to the gluino for a large range of gaugino masses.

We keep our analysis general, but to show our results, we select only one possible
source of flavor violation in the squark sector at a time, and assume the others vanish. All
diagonal entries in the squark mass matrix are set equal and we study the branching ratio
as a function of their common value $m_0^2$ and the relevant off-diagonal element. In Fig. 8
we show the dependence of $b \to s\gamma$ as a function of $\delta_{d,LR,23}$ when this is the only source of
flavor violation. The horizontal lines represent the range of values allowed experimentally
for the branching ratio. The ratio is plotted as a function of different values for the ratio
$x = m_2^2/m_0^2$. Fixing $m_0 = 500$ GeV, this corresponds to gluino masses of 200 GeV, 400
GeV and 600 GeV respectively. Negative values of $\delta_{d,LR,23}$ are more constrained than
positive values, but in any case $\delta_{d,LR,23} \leq 4\%$. This flavor violating parameter is strongly
constrained because through the $\delta_{d,LR,23}$ term, the helicity flip needed for $b \to s\gamma$ can be
realized in the exchange particle loop. Comparison with the MSSM [22, 28] shows that
this parameter is more constrained in LRSUSY, but only slightly.

The situation is very different when the only source of flavor violation is $\delta_{d,RL,23}$, as
Figure 9: Dependence of BR($b \to s \gamma$) on $\delta_{d,RL,23}$, obtained when $\tan \beta = 5$, $\mu = 500$ GeV and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_b^2/m_0^2$, 0.16(solid), 0.64(dashed) and 1.44(dot-dashed). $m_0$ is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

shown in Fig. 4. MSSM results for $b \to s \gamma$ are symmetric around $\delta_{d,RL,23} = 0$ and the experimental bounds are satisfied for any small values of $\delta_{d,RL,23}$. In LRSUSY, practically no negative values of $\delta_{d,RL,23}$ satisfy the bounds, and this flavor violating parameter is less restricted than $\delta_{d,LR,23}$ for the same values of the squark and gluino masses.

In Fig. 10 and Fig. 11 we plot the dependence of the branching ratio of $b \to s \gamma$ on the chirality conserving mixings $\delta_{d,LL,23}$ and $\delta_{d,RR,23}$ respectively, with the proviso that these are the only off-diagonal matrix elements in the squark mass matrix squared. Although the restriction is not as pronounced as the one for chirality flipping parameters, nonetheless the parameter $\delta_{d,LL,23}$ is more restricted if it is negative (to 50%) than if positive (where almost all values allowed for large gluino masses), quite different than in the MSSM, where values centered around $\delta_{d,LL,23} = 0$ were favored [22]. The same is true for the parameter $\delta_{d,RR,23}$ which in MSSM was restricted slightly only for $\pm 100\%$ values, but in LRSUSY regions of restrictions are centered around 50% , and increasing
with gluino mass for fixed scalar mass; again, any negative values are ruled out by the experimental bounds in the parameter region considered.

In Ref. [29], a detailed analysis of FCNC and CP constraints on these parameters was presented. For the decay $b \rightarrow s\gamma$, only poor constraints on $\delta_{d,LL,23}$ existed, while $\delta_{d,LR,23}$ was found to be constrained strongly. This is compatible with our analysis even though only the gluino-mediated contribution to the decay was considered there.

## 5 Conclusions

We have presented a detailed and complete analysis of all one-loop contributions to the branching ratio of $b \rightarrow s\gamma$ in the LRSUSY model. We analysed separately the case in which the only source of flavor violation comes from the quark sector (CKM matrix). We refer to that case as the constrained LRSUSY model, in analogy with MSSM. If we allow for soft-supersymmetry intergenerational mixing in the squark sector, new sources
Figure 11: Dependence of $\text{BR}(b \to s\gamma)$ on $\delta_{d,RR,23}$, obtained when $\tan \beta = 5$, $\mu = 500$ GeV and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_{\tilde{g}}^2/m_0^2$, 0.16(solid), 0.64(dashed) and 1.44(dot-dashed). $m_0$ is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

flavor violation can occur; we refer to that case as the unconstrained LRSUSY and we analyse it and compare it to MSSM under similar conditions.

The model contains too many parameters to allow for a precise restriction on any single one. However as a general feature, some constraints arise for low squark masses. In the constrained LRSUSY case, for intermediate gluino-neutralino masses, the $\mu$ parameter is favored to be such that $\mu/M_{L,R} \sim O(1)$, and a larger region of parameter space satisfies the experimental constraints for $\text{sign}(\mu) < 0$ than for $\text{sign}(\mu) > 0$. A small range of low or intermediate values of $\tan \beta$ are allowed for such a choice: for larger values of $\tan \beta$ the gaugino, Higgsino and squark masses must be higher. The branching ratio is relatively insensitive to values of squark masses above 500 GeV, where the branching ratio becomes equal to its QCD-corrected value. For a light neutralino-chargino scenario, the mass of the gluino must be $\geq 300$ GeV. For a gluino mass of order 300 GeV and very light squarks ($m_{\tilde{t}} = 100$ GeV), the left and/or right gaugino masses must be in the
600-800 GeV range.

For the unconstrained LRSUSY model, assuming flavor mixing only between the second and third generation in the down squark mass mixing matrix, the branching ratio is dominated by the internal chirality flipping diagrams, as in MSSM. Here however, the chargino graphs are comparable to the gluino contributions. The model puts stricter constrains on the chirality flipping mass mixings $\delta_{d,LR,23}$ and $\delta_{d,RL,23}$ than the chirality conserving flavor mixing parameters $\delta_{d,LL,23}$ and $\delta_{d,RR,23}$. The difference between LR-SUSY and MSSM is quite striking in restrictions on the chirality conserving $\delta_{d,LL,23}$ and $\delta_{d,RR,23}$. As opposed to MSSM where both negative and positive values of these parameters are allowed, LRSUSY severely restricts the range of the negative values. This is understood as a consequence of the left-right structure of the gauge-gaugino sector. For $\delta_{d,RR,23}$, there seems to be a small range of disallowed values in a narrow range around 50%. If the dominant sources of flavor violation come from chirality conserving sflavor mixing, the MSSM and LRSUSY allow for a distinguishingly different range of parameters.

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The relevant Feynman rules used in the calculation are listed in this appendix.

The three vertices of gluino-quark-squark, chargino-quark-squark and neutralino-quark-squark interactions are represented in Fig. 12. From the first graph:

\[(I) = -ig_s \sqrt{2} T^a_{\alpha\beta} (\Gamma^k_{QL} P_L - \Gamma^k_{QR} P_R),\]  

(42)

where \( P_{L,R} = (1 \pm \gamma_5)/2, \) and \( T^a \) are SU(3) color generators normalized to \( Tr(T^a T^b) = \delta^{ab}/2; \) and \( \Gamma_{QL,R} \) are mixing matrices for scalar quarks. From the second graph:

\[(II) = -ig C^{-1} [(G^j_{UL} - H^j_{UR}) P_L + (G^j_{UR} - H^j_{UL}) P_R],\]  

(43)

where \( C \) is the charge conjugation operator (in spinor space) and the chargino-quark-squark mixing matrices \( G \) and \( H \) are defined as:

\[G^j_{UL} = V^*_{j1}(K_{CKM})_{il}(\Gamma_{UL})_{kl}\]
\[G^j_{UR} = U_{j2}(K_{CKM})_{il}(\Gamma_{UR})_{kl}\]
\[H^j_{UL} = \frac{1}{\sqrt{2m_W}} \left( \frac{m_{ui}}{\sin \beta} U_{j3} + \frac{m_{di}}{\cos \beta} U_{j4} \right) (K_{CKM})_{il}(\Gamma_{UL})_{kl}\]
\[H^j_{UR} = \frac{1}{\sqrt{2m_W}} \left( \frac{m_{ui}}{\sin \beta} V^*_{j3} + \frac{m_{di}}{\cos \beta} V^*_{j4} \right) (K_{CKM})_{il}(\Gamma_{UR})_{kl}.\]  

(44)

Finally the contribution from the third graph is:

\[(III) = -ig \left( \sqrt{2} G^j_{0DL} + H^j_{0DR} \right) P_L - \left( \sqrt{2} G^j_{0DR} - H^j_{0DL} \right) P_R,\]  

(45)
where the neutralino-quark-squark mixing matrices $G_0$ and $H_0$ are defined as

\[
G_{0DL}^{jki} = \frac{1}{\cos \theta_W} \left[ \sin \theta_W Q_d N'_{j1} + \frac{1}{2} Q_d \left( T^3_d - Q_d \sin^2 \theta_W \right) N'_{j2} \right] \\
- \frac{\sqrt{\cos 2 \theta_W}}{\cos \theta_W} \frac{Q_u + Q_d}{2} N'_{j3} \right] (K_{CKM})_{il} (\Gamma_{DL})_{kl} \\
G_{0DR}^{jki} = -\left[ \sin \theta_W Q_d N'_{j1} - \frac{Q_d \sin^2 \theta_W}{\cos \theta_W} N'_{j2} \right] \\
+ \frac{\sqrt{\cos 2 \theta_W}}{\cos \theta_W} \left( T^3_d - Q_d \sin^2 \theta_W \right) N'_{j3} \right] (K_{CKM})_{il} (\Gamma_{DR})_{kl}
\]

\[
H_{0DL}^{jki} = \frac{1}{\sqrt{2m_W}} \left( m_{ul} N'_{j5} + m_{dl} \sin \beta N'_{j7} \right) (K_{CKM})_{il} (\Gamma_{DL})_{kl} \\
H_{0DR}^{jki} = \frac{1}{\sqrt{2m_W}} \left( m_{ul} N'_{j5} + m_{dl} \cos \beta N'_{j7} \right) (K_{CKM})_{il} (\Gamma_{DR})_{kl}
\]

(46)

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