

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

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(Received 26 February 2002; published 10 June 2002)

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 Cabibbo-Kobayashi-Maskawa (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.

DOI: 10.1103/PhysRevD.65.114006

PACS number(s): 12.38.Bx, 12.39.Hg, 12.60.Jv

I. INTRODUCTION

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

$$\text{BR}(B \rightarrow X_s \gamma) = (3.23 \pm 0.41) \times 10^{-4}. \quad (1)$$

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

$$\text{BR}(B \rightarrow X_s \gamma)_{SM} = (3.60 \pm 0.30) \times 10^{-4}. \quad (2)$$

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

$$2 \times 10^{-4} < \text{BR}(B \rightarrow X_s \gamma) < 4.5 \times 10^{-4}. \quad (3)$$

Flavor changing neutral currents (FCNCs) are forbidden at the tree level in the SM. The first SM contributions to the process $b \rightarrow s \gamma$ appear at the 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 far beyond the SM scenarios, such as two-Higgs-doublet models [6], left-right symmetric models (LRM's) [7], and the minimal supersymmetric standard model (MSSM) [8]. Although attempts have been made to reconcile $b \rightarrow s \gamma$ 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 [10–13]. 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 $SU(2)_L \times U(1)_Y$ [10]. LRSUSY was originally seen as a natural way to suppress rapid proton decay [12] and has recently received renewed attention for providing small neutrino masses and lepton radiative decays [14]. In addition to 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 $SO(10)$ [15]. 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 $M_{SUSY} \approx 10^{10-12}$ GeV, or at $M_{SUSY} \approx 1$ TeV, the difference having implications for gauge unification [16].

In this paper we study all contributions of the LRSUSY model to the branching ratio of $b \rightarrow s \gamma$ at the one-loop level. The decay $b \rightarrow s \gamma$ can be mediated by left-handed and right-handed W bosons and charged Higgs bosons as in the non-supersymmetric case, but also by charginos, neutralinos, and gluinos. The structure of LRSUSY provides a significant contribution to the decay $b \rightarrow s \gamma$ 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 structures. In Sec. III, we give the supersymmetric contributions in LRSUSY to the decay $b \rightarrow s \gamma$. 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.

II. 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

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right-handed fermions and their corresponding left- and right-handed scalar partners (sleptons and squarks) [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 bidoublets $\Phi_u(\frac{1}{2}, \frac{1}{2}, 0)$ 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 seesaw mechanism) for the right-handed neutrino and a small one for the left-handed neutrino [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 ensure cancellation of the anomalies which would otherwise occur in the fermionic sector. The superpotential for the LRSUSY model is

$$\begin{aligned} W_{LRSUSY} = & \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q^c + \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L^c \\ & + i(\mathbf{h}_{LR} L^T \tau_2 \Delta_L L + \mathbf{h}_{LR} L^{cT} \tau_2 \Delta_R L^c) \\ & + M_{LR} [\text{Tr}(\Delta_L \delta_L + \Delta_R \delta_R)] \\ & + \mu_{ij} \text{Tr}(\tau_2 \Phi_i^T \tau_2 \Phi_j) + W_{NR}, \end{aligned} \quad (4)$$

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

The neutral Higgs fields acquire nonzero vacuum expectation values (VEV's) through spontaneous symmetry breaking:

$$\langle \Delta \rangle_{L,R} = \begin{pmatrix} 0 & 0 \\ v_{L,R} & 0 \end{pmatrix}$$

and

$$\langle \Phi \rangle_{u,d} = \begin{pmatrix} \kappa_{u,d} & 0 \\ 0 & \kappa'_{u,d} e^{i\omega} \end{pmatrix}.$$

$\langle \Phi \rangle$ causes the mixing of W_L and W_R bosons with CP -violating phase ω , which is set to zero in this analysis. The nonzero Higgs VEV's break 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 κ_u and κ_d at the second stage of breaking.

In the supersymmetric sector of the model there are six singly charged charginos, corresponding to $\tilde{\lambda}_L$, $\tilde{\lambda}_R$, $\tilde{\phi}_u$,

$\tilde{\phi}_d$, $\tilde{\Delta}_L^\pm$, and $\tilde{\Delta}_R^\pm$. The model also has 11 neutralinos, corresponding to $\tilde{\lambda}_Z$, $\tilde{\lambda}_{Z'}$, $\tilde{\lambda}_V$, $\tilde{\phi}_{1u}^0$, $\tilde{\phi}_{2u}^0$, $\tilde{\phi}_{1d}^0$, $\tilde{\phi}_{2d}^0$, $\tilde{\Delta}_L^0$, $\tilde{\Delta}_R^0$, $\tilde{\delta}_L^0$, and $\tilde{\delta}_R^0$. Although Δ_L is not necessary for symmetry breaking [13], and is introduced only for preserving left-right symmetry, both Δ_L^{--} ($\tilde{\Delta}_L^{--}$) and its right-handed counterparts Δ_R^{--} ($\tilde{\Delta}_R^{--}$) play very important roles in lepton phenomenology of the LRSUSY model. The doubly charged Higgs bosons 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

$$\begin{aligned} \mathcal{L}_Y = & \mathbf{h}_u \bar{Q}_L \Phi_u Q_R + \mathbf{h}_d \bar{Q}_L \Phi_d Q_R \\ & + \mathbf{h}_\nu \bar{L}_L \Phi_u L_R + \mathbf{h}_e \bar{L}_L \Phi_d L_R + \text{H.c.}, \\ \mathcal{L}_M = & i \mathbf{h}_{LR} (L_L^T C^{-1} \tau_2 \Delta_L L_L + L_R^T C^{-1} \tau_2 \Delta_R L_R) \\ & + \text{H.c.}, \end{aligned} \quad (5)$$

where \mathbf{h}_u , \mathbf{h}_d , \mathbf{h}_ν , and \mathbf{h}_e are the Yukawa couplings for the up and down quarks and neutrino and electron, respectively, and \mathbf{h}_{LR} is the coupling for the triplet Higgs bosons. LR symmetry requires all \mathbf{h} matrices to be Hermitian in generation space and the \mathbf{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 $b \rightarrow s \gamma$.

A. Charginos

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

$$\mathcal{L}_C = -\frac{1}{2} (\psi^+, \psi^-) \begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix} \begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + \text{H.c.}, \quad (6)$$

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

$$X = \begin{pmatrix} M_L & 0 & g_L \kappa_u & 0 & 0 \\ 0 & M_R & g_R \kappa_u & 0 & 0 \\ 0 & 0 & 0 & -\mu & 0 \\ g_L \kappa_d & g_R \kappa_d & -\mu & 0 & 0 \\ 0 & \sqrt{2} g_R v_R & 0 & 0 & -\mu \end{pmatrix}, \quad (7)$$

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

$$\chi_i^+ = V_{ij} \psi_j^+, \quad \chi_i^- = U_{ij} \psi_j^-, \quad i, j = 1, \dots, 5, \quad (8)$$

with V and U unitary matrices satisfying

$$U^* X V^{-1} = M_D, \quad (9)$$

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

$$V X^\dagger X V^{-1} = U^* X X^\dagger (U^*)^{-1} = M_D^2. \quad (10)$$

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

B. Neutralinos

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

$$\mathcal{L}_N = -\frac{1}{2} \psi^{0T} Y \psi^0 + \text{H.c.}, \quad (11)$$

where

$$\psi^0 = (-i\lambda_L^3, -i\lambda_R^3, -i\lambda_V, \tilde{\phi}_{u1}^0, \tilde{\phi}_{u2}^0, \tilde{\phi}_{d1}^0, \tilde{\phi}_{d2}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0)^T,$$

and

$$Y = \begin{pmatrix} M_L & 0 & 0 & \frac{g_L \kappa_u}{\sqrt{2}} & 0 & 0 & -\frac{g_L \kappa_d}{\sqrt{2}} & 0 & 0 \\ 0 & M_R & 0 & \frac{g_R \kappa_u}{\sqrt{2}} & 0 & 0 & -\frac{g_R \kappa_d}{\sqrt{2}} & -\sqrt{2} g_{RV_R} & 0 \\ 0 & 0 & M_V & 0 & 0 & 0 & 0 & 2\sqrt{2} g_{V_R} & 0 \\ \frac{g_L \kappa_u}{\sqrt{2}} & \frac{g_R \kappa_u}{\sqrt{2}} & 0 & 0 & 0 & 0 & -\mu & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 & 0 \\ -\frac{g_L \kappa_d}{\sqrt{2}} & -\frac{g_R \kappa_d}{\sqrt{2}} & 0 & -\mu & 0 & 0 & 0 & 0 & 0 \\ 0 & -\sqrt{2} g_{RV_R} & \sqrt{2} g_{V_R} & 0 & 0 & 0 & 0 & 0 & -\mu \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu & 0 \end{pmatrix}. \quad (12)$$

The mass eigenstates are defined by

$$\chi_i^0 = N_{ij} \psi_j^0 \quad (i, j = 1, 2, \dots, 9), \quad (13)$$

where N is a unitary matrix chosen such that

$$N^* Y N^{-1} = N_D, \quad (14)$$

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

$$N Y^\dagger Y N^{-1} = N_D^2, \quad (15)$$

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

$$\psi^{0'} = (-i\lambda_\gamma, -i\lambda_{Z_L}, -i\lambda_{Z_R}, \tilde{\phi}_{u1}^0, \tilde{\phi}_{u2}^0, \tilde{\phi}_{d1}^0, \tilde{\phi}_{d2}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0)^T, \quad (16)$$

with

$$\begin{aligned}
\lambda_\gamma &= \lambda_L^3 \sin \theta_W + \lambda_R^3 \sin \theta_W + \lambda_V \sqrt{\cos 2\theta_W}, \\
\lambda_{Z_L} &= \lambda_L^3 \cos \theta_W - \lambda_R^3 \sin \theta_W \tan \theta_W - \lambda_V \sqrt{\cos 2\theta_W} \tan \theta_W, \\
\lambda_{Z_R} &= \lambda_R^3 \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} - \lambda_V \tan \theta_W.
\end{aligned} \tag{17}$$

Then the mass matrix Y would be replaced by a matrix Y' found to be

$$Y' = \begin{pmatrix} m_{\tilde{\gamma}} & 0 & 0 & F & 0 & 0 & -F & -2\sqrt{2}ev_R & 0 \\ 0 & m_{\tilde{Z}_L} & 0 & A_1 & 0 & 0 & A_2 & D & 0 \\ 0 & 0 & m_{\tilde{Z}_R} & E & 0 & 0 & C & B & 0 \\ F & A_1 & E & 0 & 0 & 0 & -\mu & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & -\mu & 0 & 0 & 0 & 0 \\ -F & A_2 & C & -\mu & 0 & 0 & 0 & 0 & 0 \\ -2\sqrt{2}ev_R & D & B & 0 & 0 & 0 & 0 & 0 & -\mu \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu & 0 \end{pmatrix}, \tag{18}$$

with

$$\begin{aligned}
m_{\tilde{\gamma}} &= M_V \cos 2\theta_W + (M_L + M_R) \sin^2 \theta_W, \\
m_{\tilde{Z}_L} &= M_V \cos 2\theta_W \tan^2 \theta_W + M_L \cos^2 \theta_W \\
&\quad + M_R \sin^2 \theta_W \tan^2 \theta_W, \\
m_{\tilde{Z}_R} &= M_V \tan^2 \theta_W + M_R (1 - \tan^2 \theta_W), \\
A_1 &= \frac{g}{\sqrt{2} \cos \theta_W} (\kappa_u \cos^2 \theta_W - \kappa_d \sin \theta_W), \\
A_2 &= \frac{g}{\sqrt{2} \cos \theta_W} (\kappa_d \cos^2 \theta_W - \kappa_u \sin \theta_W), \\
B &= -\sqrt{2} g \frac{1 - 2 \tan^2 \theta_W}{\sqrt{1 - \tan^2 \theta_W}} v_R, \\
C &= -\frac{g \sqrt{\cos 2\theta_W}}{\sqrt{2} \cos \theta_W} \kappa_u, \\
D &= -\frac{\sqrt{2} g \sqrt{\cos 2\theta_W}}{\cos \theta_W} v_R, \\
E &= \frac{g \sqrt{\cos 2\theta_W}}{\sqrt{2} \cos \theta_W} \kappa_d, \\
F &= \frac{e}{\sqrt{2}} (\kappa_u + \kappa_d).
\end{aligned} \tag{19}$$

The unitary matrix N would be replaced by a new matrix N' , given by

$$\begin{aligned}
N'_{j1} &= N_{j1} \sin \theta_W + N_{j2} \sin \theta_W + N_{j3} \sqrt{\cos 2\theta_W}, \\
N'_{j2} &= N_{j1} \cos \theta_W - N_{j2} \sin \theta_W \tan \theta_W \\
&\quad - N_{j3} \sqrt{\cos 2\theta_W} \tan \theta_W, \\
N'_{j3} &= N_{j2} \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} - N_{j3} \tan \theta_W, \\
N'_{jk} &= N_{jk} (k = 4, 5, \dots, 9).
\end{aligned} \tag{20}$$

Similarly, N can be expressed in terms of N' by

$$\begin{aligned}
N_{j1} &= N'_{j1} \sin \theta_W + N'_{j2} \cos \theta_W, \\
N_{j2} &= N'_{j1} \sin \theta_W - N'_{j2} \sin \theta_W \tan \theta_W \\
&\quad - N'_{j3} \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W}, \\
N_{j3} &= N'_{j1} \sqrt{\cos 2\theta_W} - N'_{j2} \tan \theta_W \sqrt{\cos 2\theta_W} \\
&\quad - N_{j3} \tan \theta_W, \\
N_{jk} &= N'_{jk} (k = 4, 5, \dots, 9).
\end{aligned} \tag{21}$$

C. Squarks

(19) In the interaction basis $(\tilde{q}_L^i, \tilde{q}_R^i)$, the squared-mass matrix for a squark of flavor f has the form:

$$\mathcal{M}_f^2 = \begin{pmatrix} m_{f,LL}^2 + F_{f,LL} + D_{f,LL} & (m_{f,LR}^2) + F_{f,LR} \\ (m_{f,LR}^2)^\dagger + F_{f,RL} & m_{f,RR}^2 + F_{f,RR} + D_{f,RR} \end{pmatrix}. \quad (22)$$

The F terms are diagonal in the flavor space, $F_{f,LL,RR} = m_f^2$, $(F_{dLR})_{ij} = -\mu(m_{d_i} \tan \beta) \mathbf{1}_{ij}$, $(F_{uLR})_{ij} = -\mu(m_{u_i} \cot \beta) \mathbf{1}_{ij}$. The D terms was also flavor diagonal:

$$D_{fLL} = M_Z^2 \cos 2\beta (T_f^3 - Q_f \sin^2 \theta_W) \mathbf{1},$$

$$D_{fRR} = M_Z^2 \cos 2\beta Q_f \sin^2 \theta_W \mathbf{1}. \quad (23)$$

The term $(m_{f,LL,RR}^2)_{ij} = m_{\tilde{Q}_{L,R}}^2 \delta_{ij}$, $m_{fLR}^2 = A_f^* m_f$. To reduce the number of free parameters, we consider the parameters to be universal, with $(m_{\tilde{Q}_{L,R}}^2)_{ij} = m_0^2 \delta_{ij}$, $A_{d,ij} = A \delta_{ij}$, and $A_{u,ij} = A \delta_{ij}$.

The squared-mass matrix for U type squarks reduces to

$$\mathcal{M}_{U_k}^2 = \begin{pmatrix} m_0^2 + M_Z^2 (T_u^3 - Q_u \sin^2 \theta_W) \cos 2\beta & m_{u_k} (A - \mu \cot \beta) \\ m_{u_k} (A - \mu \cot \beta) & m_0^2 + M_Z^2 Q_u \sin^2 \theta_W \cos 2\beta \end{pmatrix} \quad (24)$$

with the diagonal F terms absorbed into the m_0^2 . For D type squarks

$$\mathcal{M}_{D_k}^2 = \begin{pmatrix} m_0^2 + M_Z^2 (T_d^3 - Q_d \sin^2 \theta_W) \cos 2\beta & m_{d_k} (A - \mu \tan \beta) \\ m_{d_k} (A - \mu \tan \beta) & m_0^2 + M_Z^2 Q_d \sin^2 \theta_W \cos 2\beta \end{pmatrix}. \quad (25)$$

The corresponding mass eigenstates are defined as

$$\tilde{q}_{L,R} = \Gamma_{QL,R}^\dagger \tilde{q}, \quad (26)$$

where $\Gamma_{QL,R}^\dagger$ are 6×3 mixing matrices. In the universal case, there are no intergenerational mixings for squarks and the only source of flavor mixing comes from the CKM matrix. We will analyze 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 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

$$\delta_{d,LL,ij} = \frac{(m_{d,LL}^2)_{ij}}{m_0^2}, \quad \delta_{d,RR,ij} = \frac{(m_{d,RR}^2)_{ij}}{m_0^2},$$

$$\delta_{d,LR,ij} = \frac{(m_{d,LR}^2)_{ij}}{m_0^2}, \quad \delta_{d,RL,ij} = \frac{(m_{d,RL}^2)_{ij}}{m_0^2}. \quad (27)$$

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-squark mass matrix only.

III. SUPERSYMMETRIC CONTRIBUTIONS TO $b \rightarrow s \gamma$

The low-energy effective Hamiltonian responsible for the B -meson decay rates at the scale μ can be written as

$$\mathcal{H}_{eff} = -\frac{4G_F}{\sqrt{2}} K_{tb} K_{ts}^* \sum_i C_i(\mu) Q_i(\mu). \quad (28)$$

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

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

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

$$Q_8 = \frac{g_s}{16\pi^2} m_b(\mu) \bar{s} \sigma_{\mu\nu} G_a^{\mu\nu} T^a P_R b,$$

$$Q'_8 = \frac{g_s}{16\pi^2} m_b(\mu) \bar{s} \sigma_{\mu\nu} G_a^{\mu\nu} T^a P_L b. \quad (29)$$

The Wilson coefficients $C'_{7,8}$ are initially evaluated at the electroweak or soft supersymmetry breaking scale, then evolved down to the scale μ . 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,

$$M_{\gamma_R} = A_g^R + A_{\tilde{\chi}^-}^R + A_{\tilde{\chi}^0}^R, \quad (30)$$

with the gluino, chargino, and neutralino contributions given by

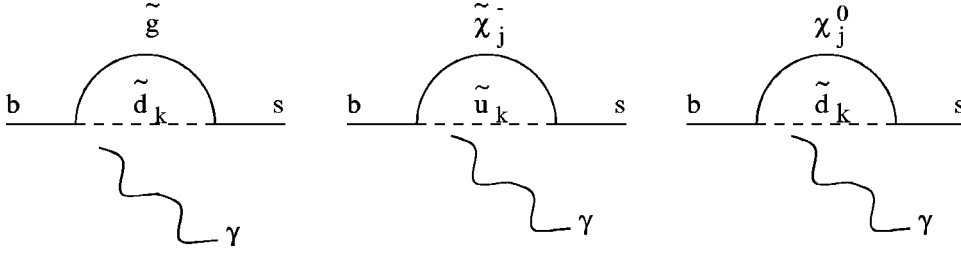


FIG. 1. The Feynman diagrams contributing to the decay $b \rightarrow s \gamma$. The outgoing photon line can be attached in all possible ways.

$$A_g^R = -\frac{\pi\alpha_s}{\sqrt{2}G_F} Q_d C(R) \sum_{k=1}^6 \frac{1}{m_{\tilde{d}_k}^2} \left\{ \Gamma_{DL}^{kb} \Gamma_{DL}^{*ks} F_2(x_{\tilde{g}\tilde{d}_k}) - \frac{m_{\tilde{g}}}{m_b} \Gamma_{DR}^{kb} \Gamma_{DL}^{*ks} F_4(x_{\tilde{g}\tilde{d}_k}) \right\}, \quad (31)$$

$$A_{\tilde{\chi}^-}^R = -\frac{\pi\alpha_w}{\sqrt{2}G_F} \sum_{j=1}^5 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_k}^2} \left\{ (G_{UL}^{jkb} - H_{UR}^{jkb})(G_{UL}^{*jks} - H_{UR}^{*jks}) [F_1(x_{\tilde{\chi}_j^- \tilde{u}_k}) + Q_u F_2(x_{\tilde{\chi}_j^- \tilde{u}_k})] \right. \\ \left. + \frac{m_{\tilde{\chi}_j^-}}{m_b} (G_{UR}^{jkb} - H_{UL}^{jkb})(G_{UL}^{*jks} - H_{UR}^{*jks}) [F_3(x_{\tilde{\chi}_j^- \tilde{u}_k}) + Q_u F_4(x_{\tilde{\chi}_j^- \tilde{u}_k})] \right\}, \quad (32)$$

$$A_{\tilde{\chi}^0}^R = -\frac{\pi\alpha_w}{\sqrt{2}G_F} Q_d \sum_{j=1}^9 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_k}^2} \left\{ (\sqrt{2}G_{0DL}^{jkb} - H_{0DR}^{jkb})(\sqrt{2}G_{0DL}^{*jks} - H_{0DR}^{*jks}) F_2(x_{\tilde{\chi}_j^0 \tilde{d}_k}) \right. \\ \left. + \frac{m_{\tilde{\chi}_j^0}}{m_b} (\sqrt{2}G_{0DR}^{jkb} - H_{0DL}^{jkb})(\sqrt{2}G_{0DL}^{*jks} - H_{0DR}^{*jks}) F_4(x_{\tilde{\chi}_j^0 \tilde{d}_k}) \right\}, \quad (33)$$

and for the decay of b_R :

$$M_{\gamma_L} = A_g^L + A_{\tilde{\chi}^-}^L + A_{\tilde{\chi}^0}^L, \quad (34)$$

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

$$A_g^L = -\frac{\pi\alpha_s}{\sqrt{2}G_F} Q_d C(R) \sum_{k=1}^6 \frac{1}{m_{\tilde{d}_k}^2} \left\{ \Gamma_{DR}^{kb} \Gamma_{DR}^{*ks} F_2(x_{\tilde{g}\tilde{d}_k}) - \frac{m_{\tilde{g}}}{m_b} \Gamma_{DL}^{kb} \Gamma_{DR}^{*ks} F_4(x_{\tilde{g}\tilde{d}_k}) \right\}, \quad (35)$$

$$A_{\tilde{\chi}^-}^L = -\frac{\pi\alpha_w}{\sqrt{2}G_F} \sum_{j=1}^5 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_k}^2} \left\{ (G_{UR}^{jkb} - H_{UL}^{jkb})(G_{UR}^{*jks} - H_{UL}^{*jks}) [F_1(x_{\tilde{\chi}_j^- \tilde{u}_k}) + Q_u F_2(x_{\tilde{\chi}_j^- \tilde{u}_k})] \right. \\ \left. + \frac{m_{\tilde{\chi}_j^-}}{m_b} (G_{UL}^{jkb} - H_{UR}^{jkb})(G_{UR}^{*jks} - H_{UL}^{*jks}) [F_3(x_{\tilde{\chi}_j^- \tilde{u}_k}) + Q_u F_4(x_{\tilde{\chi}_j^- \tilde{u}_k})] \right\}, \quad (36)$$

$$A_{\tilde{\chi}^0}^L = -\frac{\pi\alpha_w}{\sqrt{2}G_F} Q_d \sum_{j=1}^9 \sum_{k=1}^6 \frac{1}{m_{\tilde{u}_k}^2} \left\{ (\sqrt{2}G_{0DR}^{jkb} - H_{0DL}^{jkb})(\sqrt{2}G_{0DR}^{*jks} - H_{0DL}^{*jks}) F_2(x_{\tilde{\chi}_j^0 \tilde{d}_k}) \right. \\ \left. + \frac{m_{\tilde{\chi}_j^0}}{m_b} (\sqrt{2}G_{0DL}^{jkb} - H_{0DR}^{jkb})(\sqrt{2}G_{0DR}^{*jks} - H_{0DL}^{*jks}) F_4(x_{\tilde{\chi}_j^0 \tilde{d}_k}) \right\}, \quad (37)$$

where the vertex mixing matrices G , H , G_0 , and H_0 are defined in the Appendix. The convention $x_{ab} = m_a^2/m_b^2$ 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 experimen-

tal 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 [20]. There is no mixing between left- and right-handed contributions.

$$A^\gamma(m_b) = \eta^{-16/23} \left\{ A^\gamma(M_W) + A_0^\gamma \left[\frac{116}{135} (\eta^{28/23} - 1) + \frac{58}{189} (\eta^{28/23} - 1) \right] \right\}, \quad (38)$$

where $\eta = \alpha_s(m_b)/\alpha_s(M_W)$ and $A_0^\gamma = (\pi\alpha_w/2\sqrt{2}G_F)(1/M_W^2)$. 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

$$\Gamma(b \rightarrow s \gamma) = \frac{m_b^5 G_F^2 |K_{tb} K_{ts}^*|^2 \alpha}{32\pi^4} (\hat{M}_{\gamma L}^2 + \hat{M}_{\gamma R}^2), \quad (39)$$

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

$$\text{BR}(b \rightarrow s \gamma) = \frac{\Gamma(b \rightarrow s \gamma)}{\Gamma_{SL}} \text{BR}_{SL}, \quad (40)$$

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

$$\Gamma_{SL} = \frac{m_b^5 G_F^2 |K_{cb}|^2}{192\pi^3} g(z), \quad (41)$$

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

IV. NUMERICAL RESULTS

We are interested in analyzing 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 [21]. There are some extra constraints in the nonsupersymmetric sector of the theory, requiring the FCNC Higgs boson Φ_d to be heavy, but no such constraints exist in the Higgsino sector [22]. We choose the gluino mass $m_{\tilde{g}} = 300$ GeV, and the left-handed gaugino mass $M_L = 500$ GeV. The mass of the lightest bottom squark will be in the 150–200 GeV range. We include in our numerical estimates the nonsupersymmetric LRM results, and for this we constrain the lightest Higgs boson mass to be 115 GeV [23]. 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 μ 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 the 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 families.

We now proceed to discuss both these scenarios in turn.

A. 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 can 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 the 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 many fewer parameters at mass scales $\mu = M_{GUT}$, and then use renormalization group equations to run them down to the low-energy scale which is relevant to phenomenology. But, to maintain 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 the 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 $M_{0UL} = M_{0UR} = M_{0DL} = M_{0DR} = m_0$. 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). The experimental bounds are shown 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 ratio 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

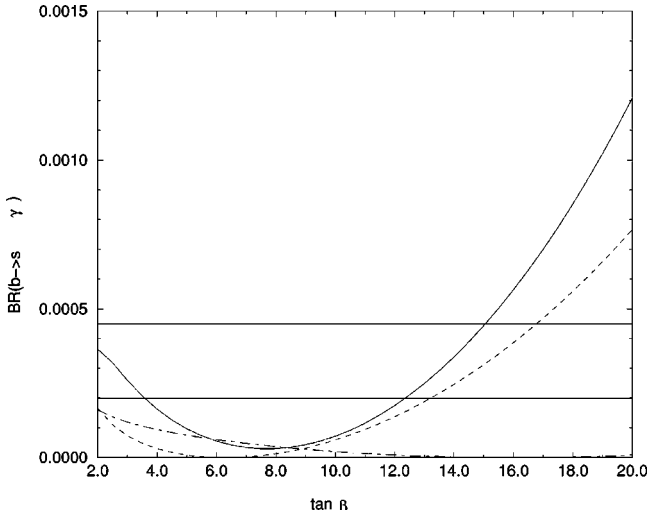


FIG. 2. Supersymmetric contributions to $\text{BR}(b \rightarrow s \gamma)$ as a function of $\tan \beta$, obtained when $m_{\tilde{g}} = 300$ GeV, $\mu = 100$ GeV, $M_L = M_R = 500$ GeV, and $m_0 = 200$ (300) GeV [dashed (dot-dashed)]. The full contribution for $m_0 = 200$ GeV is also shown (solid). The range of acceptable values of branching ratios is given.

$\tan \beta$ are allowed only for a heavier supersymmetric mass spectrum. For $\text{sgn}(\mu) < 0$, the range of acceptable intermediate values of $\tan \beta$ increases. For example, if $\mu = -100$ GeV, the larger range $\tan \beta = 22\text{--}33$ is allowed. We note that constraints on supersymmetry with large $\tan \beta$ were studied in Ref. [24] where $\tan \beta$ -enhanced chargino and charged Higgs contributions were resummed to all orders in perturbation theory.

In Fig. 2, we also present the curve corresponding to $m_0 = 300$ GeV with the other parameters the same. It can be seen that the supersymmetric contributions to the branching ratio are very small as the masses of scalar quarks are larger than the present experimental limits.

As a general feature of the LRSUSY branching ratio, in a large region of parameter space, the chargino contribution is comparable to that of 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 [25]. 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 Figs. 4 and 5, we investigate the dependence of the branching ratio of $b \rightarrow s \gamma$ on the sign and magnitude of the Higgsino mixing parameter μ . There have been indications that the new accurate measurements of the anomalous magnetic moment of the muon restrict the μ parameter to be positive, while $b \rightarrow s \gamma$ favors a negative sign. For a light squark-gaugino scenario, and low $\tan \beta$, the bound on $b \rightarrow s \gamma$ is satisfied for either sign of the μ parameter. In Fig. 4 one could see that a restricted region of intermediate values for μ with $\text{sgn}(\mu) > 0$ is allowed by the experimental constraints on $b \rightarrow s \gamma$, in the 225–325 GeV region. The parameter space is less restrictive for μ negative, to -175 GeV.

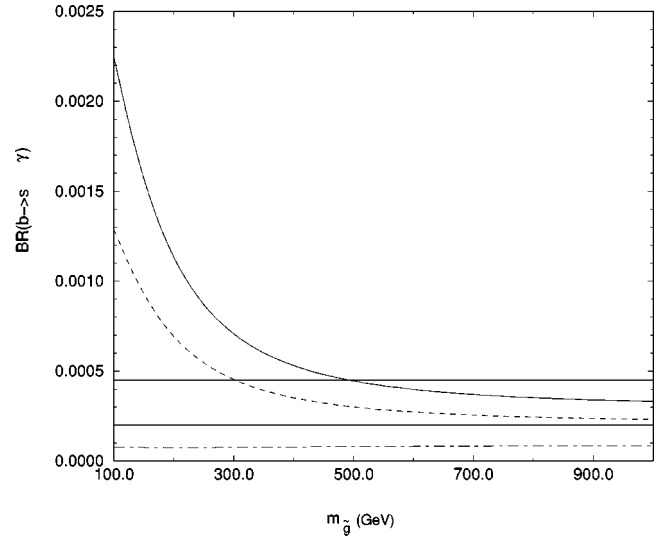


FIG. 3. Supersymmetric contributions to $\text{BR}(b \rightarrow s \gamma)$ as a function of the mass of the gluino $m_{\tilde{g}}$, obtained when $\tan \beta = 5$, $\mu = 100$ GeV, $M_L = M_R = 500$ GeV, and $m_0 = 100$ (300) GeV [dashed (dot-dashed)]. The full contribution for $m_0 = 100$ GeV is also shown (solid). The range of acceptable values of branching ratios is given.

Note that for $\mu \rightarrow 0$ the branching ratio drops outside the allowed range. This phenomenon occurs because the mixing term obtained from flipping chirality on the gaugino leg decouples. We reject such small values of the μ 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

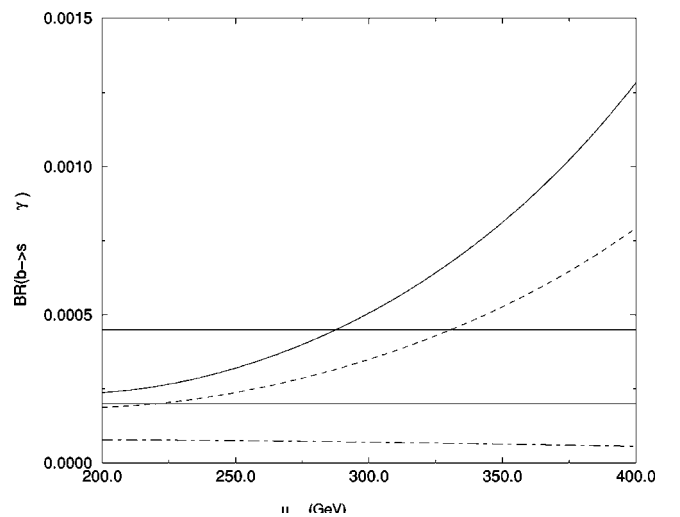


FIG. 4. Supersymmetric contributions to $\text{BR}(b \rightarrow s \gamma)$ as a function of μ , obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV, and $m_0 = 100$ (300) GeV [dashed (dot-dashed)]. The full contribution $m_0 = 100$ GeV is also shown (solid). The range of acceptable values of branching ratios is given.

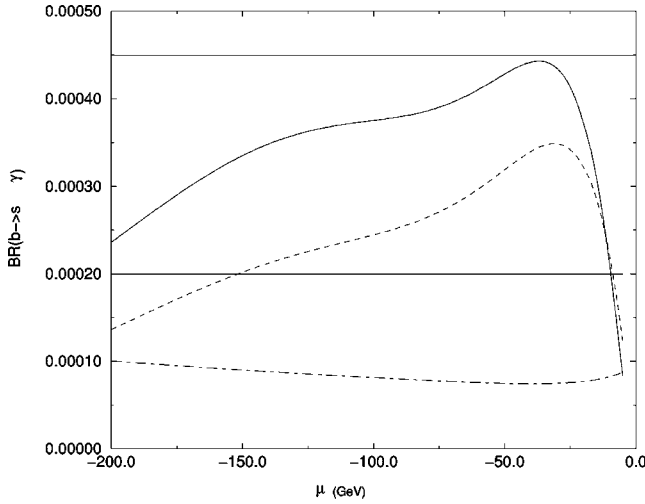


FIG. 5. Supersymmetric contributions to $BR(b \rightarrow s \gamma)$ as a function of μ , obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV, and $m_0 = 150$ (300) GeV [dashed (dot-dashed)]. The full contribution $m_0 = 150$ GeV is also shown (solid). The range of acceptable values of branching ratios is given.

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. Compared 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-handed symmetry is broken at the same scale as supersymmetry; we expect in those cases to have

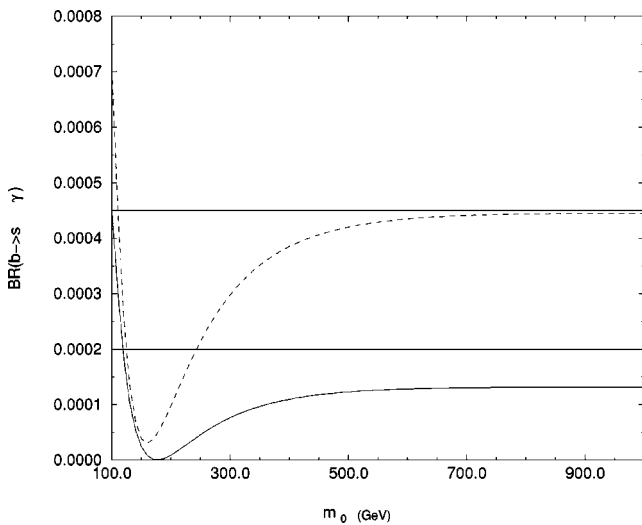


FIG. 6. Supersymmetric contributions to $BR(b \rightarrow s \gamma)$ as a function of m_0 , obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $M_L = M_R = 500$ GeV, and $\mu = 100$ GeV (dashed). The full contribution is also shown (solid). The range of acceptable values of branching ratios is given.

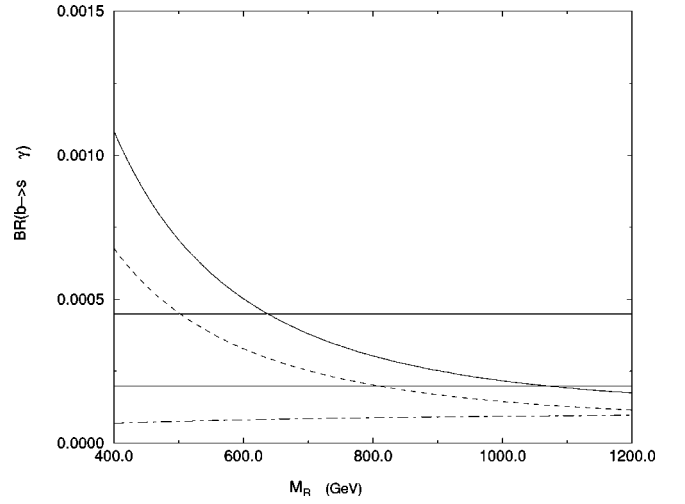


FIG. 7. Supersymmetric contributions to $BR(b \rightarrow s \gamma)$ as a function of M_R , obtained when $\tan \beta = 5$, $m_{\tilde{g}} = 300$ GeV, $\mu = 100$ GeV, $m_0 = 100$ (300) GeV [dashed (dot-dashed)], and $M_L = M_R$ is assumed. The full contribution $m_0 = 100$ GeV is also shown (solid). The range of acceptable values of branching ratios is given.

approximately $M_L = M_R$ [26]. For light squarks, Higgsinos, and gluinos, the gaugino mass must be heavy, in the 600–800 GeV range.

B. 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 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 α_s/α and $m_{\tilde{g}}/m_b$ enhancements, respectively. In this case only the gluino scenario is analyzed in the MSSM, and found to be dominated by δ_{23}^{RL} [20]. In LRSUSY, the situation is different: the chargino graph contribution is comparable to that of the gluino for a large range of gaugino masses.

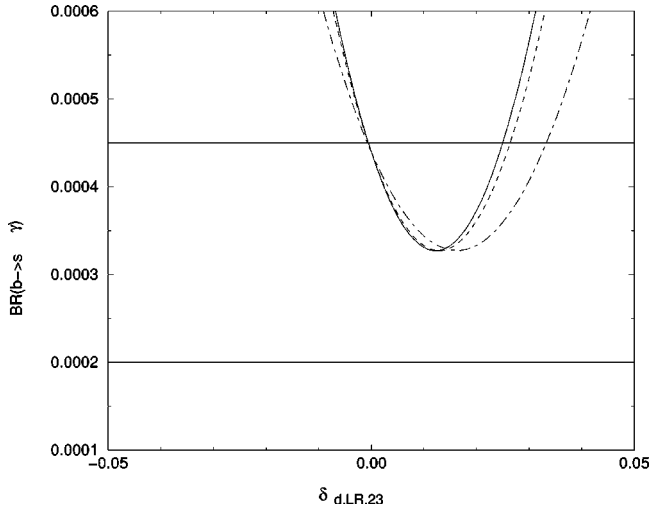


FIG. 8. Dependence of $\text{BR}(b \rightarrow s \gamma)$ on $\delta_{d,LR,23}$, obtained when $\tan \beta = 5$, $\mu = 500$ GeV, and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_g^2/m_0^2$, 0.16 (solid), 0.64 (dashed), and 1.44 (dot-dashed). m_0 is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

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 that 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 $b \rightarrow s \gamma$ 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 $x = m_g^2/m_0^2$. 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 $b \rightarrow s \gamma$ can be realized in the exchange particle loop. Comparison with the MSSM [20,27] 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 shown in Fig. 9. MSSM results for $b \rightarrow s \gamma$ 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, essentially 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 Figs. 10 and 11 we plot the dependence of the branching ratio of $b \rightarrow s \gamma$ 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 are allowed for large

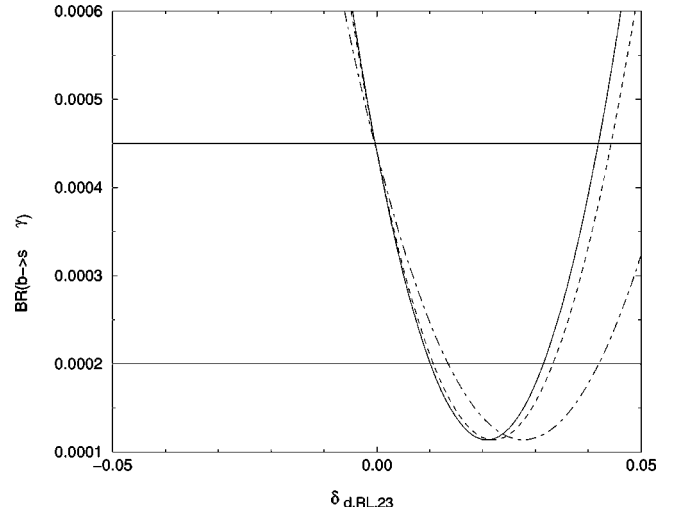


FIG. 9. Dependence of $\text{BR}(b \rightarrow s \gamma)$ on $\delta_{d,RL,23}$, obtained when $\tan \beta = 5$, $\mu = 500$ GeV, and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_g^2/m_0^2$, 0.16 (solid), 0.64 (dashed), and 1.44 (dot-dashed). m_0 is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

gluino masses), quite different from in the MSSM, where values centered around $\delta_{d,LL,23} = 0$ were favored [20]. 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 increase 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. [28], 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 com-

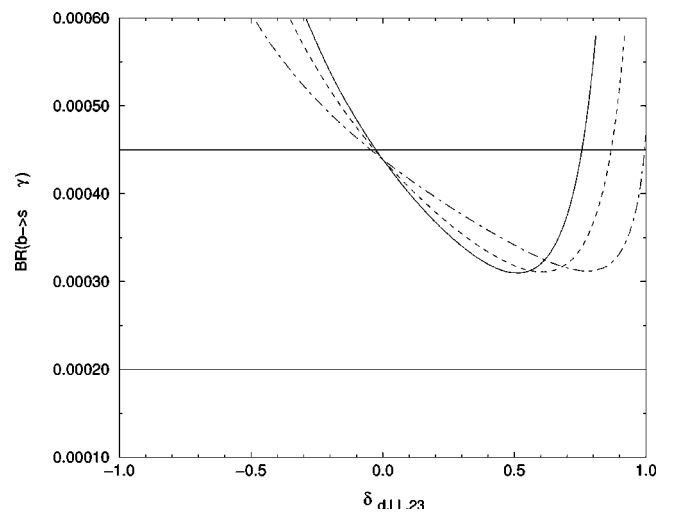


FIG. 10. Dependence of $\text{BR}(b \rightarrow s \gamma)$ on $\delta_{d,LL,23}$, obtained when $\tan \beta = 5$, $\mu = 400$ GeV, and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_g^2/m_0^2$, 0.16 (solid), 0.64 (dashed), and 1.44 (dot-dashed). m_0 is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

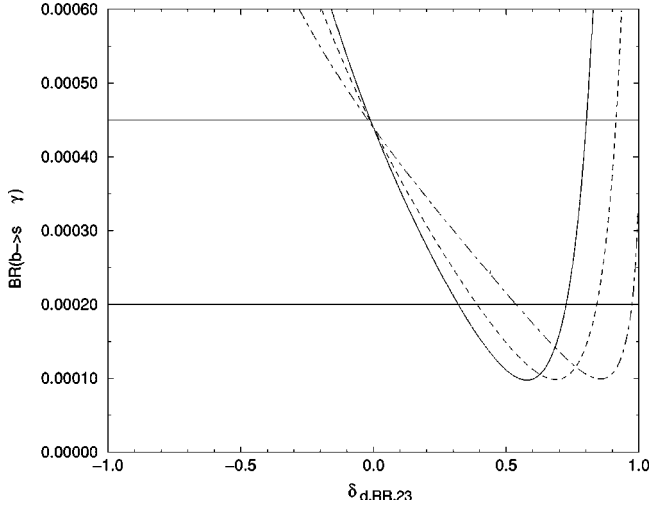


FIG. 11. Dependence of $BR(b \rightarrow s \gamma)$ on $\delta_{d,RR,23}$, obtained when $\tan \beta = 5$, $\mu = 500$ GeV, and $M_L = M_R = 500$ GeV. The different lines correspond to different values of $x = m_g^2/m_0^2$, 0.16 (solid), 0.64 (dashed), and 1.44 (dot-dashed). m_0 is fixed to be 500 GeV. The range of acceptable values of branching ratios is given.

patible with our analysis even though only the gluino-mediated contribution to the decay was considered there.

V. 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 analyzed 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 the MSSM. If we allow for soft supersymmetry intergenerational mixing in the squark sector, new sources of flavor violation can occur; we refer to that case as the unconstrained LRSUSY model and we analyze it and compare it to the 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 μ parameter is favored to be such that $\mu/M_{L,R} \sim \mathcal{O}(1)$, and a larger region of parameter space satisfies the experimental constraints for $\text{sgn}(\mu) < 0$ than for $\text{sgn}(\mu) > 0$. A small range of low or intermediate values of $\tan \beta$ is 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 ≥ 300 GeV. For a gluino mass of order 300 GeV and very light squarks ($m_{\tilde{q}} = 100$ 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 generations in the down-squark mass mixing matrix, the branching ratio is dominated by the internal chirality flipping diagrams, as in

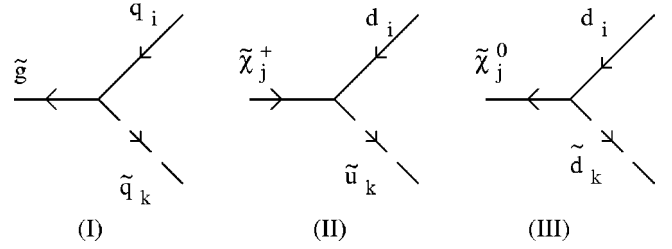


FIG. 12. The Feynman rules used in the calculation.

the MSSM. Here, however, the chargino graphs are comparable to the gluino contributions. The model puts stricter constraints on the chirality flipping mass mixings $\delta_{d,LR,23}$ and $\delta_{d,RL,23}$ than on the chirality conserving flavor mixing parameters $\delta_{d,LL,23}$ and $\delta_{d,RR,23}$. The difference between LRSUSY and the MSSM is quite striking in restrictions on the chirality conserving $\delta_{d,LL,23}$ and $\delta_{d,RR,23}$. As opposed to the MSSM where both negative and positive values of these parameters are allowed, LRSUSY severely restricts the range of 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 distinguishably different range of parameters.

ACKNOWLEDGMENTS

This work was funded by NSERC of Canada (SAP0105354).

APPENDIX

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

$$(I) = -i g_s \sqrt{2} T_{\alpha\beta}^a (\Gamma_{QL}^{ki} P_L - \Gamma_{QR}^{ki} P_R), \quad (A1)$$

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

$$(II) = -i g C^{-1} [(G_{UL}^{ki} - H_{UR}^{ki}) P_L + (G_{UR}^{ki} - H_{UL}^{ki}) P_R], \quad (A2)$$

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

$$G_{UL}^{ki} = V_{j1}^* (K_{CKM})_{il} (\Gamma_{UL})_{kl},$$

$$G_{UR}^{ki} = U_{j2} (K_{CKM})_{il} (\Gamma_{UR})_{kl},$$

$$\begin{aligned}
 H_{UL}^{jki} &= \frac{1}{\sqrt{2}m_W} \left(\frac{m_{u_l}}{\sin \beta} U_{j3} + \frac{m_{d_l}}{\cos \beta} U_{j4} \right) \\
 &\quad \times (K_{CKM})_{il} (\Gamma_{UL})_{kl}, \\
 H_{UR}^{jki} &= \frac{1}{\sqrt{2}m_W} \left(\frac{m_{u_l}}{\sin \beta} V_{j3}^* + \frac{m_{d_l}}{\cos \beta} V_{j4}^* \right) \\
 &\quad \times (K_{CKM})_{il} (\Gamma_{UR})_{kl}. \quad (\text{A3})
 \end{aligned}$$

Finally, the contribution from the third graph is

$$\begin{aligned}
 (\text{III}) &= -ig [(\sqrt{2}G_{0DL}^{jki} + H_{0DR}^{jki}) P_L \\
 &\quad - (\sqrt{2}G_{0DR}^{jki} - H_{0DL}^{jki}) P_R], \quad (\text{A4})
 \end{aligned}$$

where the neutralino-quark-squark mixing matrices G_0 and H_0 are defined as

$$\begin{aligned}
 G_{0DL}^{jki} &= \left[\sin \theta_W Q_d N'_{j1} + \frac{1}{\cos \theta_W} (T_d^3 - Q_d \sin^2 \theta_W) N'_{j2} \right. \\
 &\quad \left. - \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} \frac{Q_u + Q_d}{2} N'_{j3} \right] (K_{CKM})_{il} (\Gamma_{DL})_{kl}, \\
 G_{0DR}^{jki} &= - \left[\sin \theta_W Q_d N'_{j1} - \frac{Q_d \sin^2 \theta_W}{\cos \theta_W} N'_{j2} \right. \\
 &\quad \left. + \frac{\sqrt{\cos 2\theta_W}}{\cos \theta_W} (T_d^3 - Q_d \sin^2 \theta_W) N'_{j3} \right] \\
 &\quad \times (K_{CKM})_{il} (\Gamma_{DR})_{kl}, \\
 H_{0DL}^{jki} &= \frac{1}{\sqrt{2}m_W} \left(\frac{m_{u_l}}{\sin \beta} N'_{j5} + \frac{m_{d_l}}{\cos \beta} N'_{j7} \right) \\
 &\quad \times (K_{CKM})_{il} (\Gamma_{DL})_{kl}, \\
 H_{0DR}^{jki} &= \frac{1}{\sqrt{2}m_W} \left(\frac{m_{u_l}}{\sin \beta} N'_{j5}^* + \frac{m_{d_l}}{\cos \beta} N'_{j7}^* \right) \\
 &\quad \times (K_{CKM})_{il} (\Gamma_{DR})_{kl}. \quad (\text{A5})
 \end{aligned}$$

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