$t \to cg, c\gamma, cZ$ in the left-right supersymmetric model

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We analyze top-quark flavor violating decays into a charm quark and a gluon, photon, or Z boson in a supersymmetric model incorporating left-right symmetry. We include loop calculations involving contributions from scalar quarks, gluinos, charginos, and neutralinos. We perform the calculations first assuming the minimal (flavor-diagonal) scalar quark scenario and then allowing for arbitrary mixing between the second and the third generation of scalar quarks, in both the up and the down sectors. In each case we present separately the contributions from gluino, chargino, and neutralino loops and compare their respective strengths. In the flavor-diagonal case, the branching ratio cannot exceed 10^{-5} (10^{-6}) for the gluon (photon/Z boson); while for the unconstrained (flavor-nondiagonal) case the same branching ratios can reach almost 10^{-4} for the gluon, 10^{-6} for the photon, and 10^{-5} for the Z boson, all of which are slightly below the expected reach of LHC.

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I. INTRODUCTION

Flavor changing interactions, in general, and of top quark, in particular, are a promising test ground for new physics. In the standard model (SM) processes such as $t \to cg, \gamma, Z$ are absent at tree level and highly suppressed by the Glashow-Iliopoulos-Maiani mechanism at one loop. The branching ratios predicted in the SM [1] are of order of $10^{-10}; 10^{-12}; 10^{-13} - 10^{-12}$ for the decays $t \to cg; c\gamma; cZ$ and thus far from present and future reaches of e^+e^- colliders or the LHC. Models beyond the SM predict branching ratios which are orders of magnitude larger than this, and, thus, an experimental signal can be interpreted as a signal for new physics. An added bonus of top-quark phenomenology is that it provides a window into the electroweak breaking mechanism, since the top mass is close to the scale of symmetry breaking.

Previous analyses have explored several possibilities of enhancing flavor violating decays, mostly within supersymmetry. $t \to cV$ ($V = g, \gamma, Z$) have also been explored in two Higgs doublet models [1,2], technicolor models [3], top-color assisted technicolor models [4], models with extra vector singlets [5], and supersymmetry with [6–12], and without [13], R parity. While supersymmetry appears to provide the most convincing scenario for new physics, most analyses concentrate on the minimal supersymmetric standard model (MSSM). Analyses have been performed initially in the constrained MSSM, where flavor changing interactions are driven by flavor universal supersymmetry breaking soft terms including left, right, and intergenerational scalar quark mixing. Li *et al.* evaluated one-loop supersymmetry (SUSY) QCD and electroweak contributions [6] and obtained branching ratios of 10^{-6} for $t \to cg$ and 10^{-8} for $t \to c\gamma, Z$. Couture *et al.* included the left-handed squark mixing in their calculation [7] and obtained branching ratios of 10^{-5} for the gluon, 10^{-7} for the photon, and 10^{-6} for the Z boson. Inclusion of the right-handed squark mixing [8] did not change branching ratios significantly. Lopez *et al.* further refined the calculations in the flavor-diagonal case by including contributions from the neutralino quark-squark loops; in this case the branching ratios obtained were 10^{-5} for $t \to cg$ and 10^{-7} for $t \to c\gamma, Z$. In the nonuniversal case several authors [10] found branching ratios of 10^{-5} for $t \to cg$ and 10^{-6} for $t \to c\gamma, Z$.

In the unconstrained MSSM, assumptions about universality of soft supersymmetry breaking terms and new sources of flavor violation are included in the scalar quark matrices. Liu *et al.* looked at the fully unconstrained MSSM and found decay branching ratios which can reach 10^{-4}, 10^{-6} for $t \to cg$ and $t \to c\gamma, Z$, respectively; while Delépine and Khalil [12] found, for the unconstrained MSSM with a light stop, branching ratios of 10^{-5} for the gluon decay and 10^{-6} for the photon. Note that not all authors use the same values for relevant parameters. Experimental limits on SUSY masses have become more stringent; in particular, the gluino mass is taken to be higher in later papers, which affects the estimated branching ratios.

However, not much work has been done on flavor changing interactions in the top-quark decays in scenarios beyond the MSSM. We present here such an analysis, based on extending the SM group $SU(3)_C \times SU(2)_L \times U(1)_Y$ to a left-right symmetric group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (LRSUSY). This gauge symmetry allows for the seesaw mechanism within a supersymmetric scenario and predicts neutrino masses and mixing naturally. Flavor violation in b decays has shown possible enhancement of results over the MSSM ones: in particular, restrictions on intergenerational left-left (LL), right-right (RR), and left-right (LR) mixing of scalar quarks are more restrictive than in the case of MSSM [14].

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In this work, we present a full investigation of the flavor changing two-body decays of the top quark in LRSUSY. We investigate both the constrained case, in which the only source of flavor violation comes from the Cabibbo-Kobayashi-Maskawa (CKM) matrix in the quark sector (with the proviso that we assume it to be the same for left- and right-handed quarks), as well as in the unconstrained model, in which soft symmetry breaking parameters are allowed to induce flavor-dependent mixing in the squark mass matrix. Such a flavor-dependent scenario is motivated by neutrino oscillations, which indicate the presence of two large mixing angles (between first and second, and second and third generations). If LRSUSY occurs in nature as an intermediate symmetry in breaking down from a SUSY grand unified theory scenario, such as $SO(10)$, large neutrino mixing angles in the lepton mass matrices may appear not in the down quark mass matrices (because the right-handed charged current interaction is broken down at the unification scale), but in the squark mass matrices [15]. The new mixing might affect b [16] as well as t flavor physics.

Previous studies of unconstrained supersymmetric models have been based on either the mass insertion [17] or on the mass eigenstate method [16,18]. In the mass insertion framework [17], one chooses a basis for fermion and sfermion states in which all the couplings of these particles to neutral gauginos are flavor-diagonal. Flavor changes in the squark sector arise from the nondiagonality of the squark propagators. Off-diagonal elements mix squark flavors for both left- and right-handed squarks. In the mass eigenstate method [18], squark mass matrices are given in the super-CKM basis and are diagonalized by rotating the superfields. In this basis, the up-squark and down-squark mass matrices are correlated by this rotation and thus not independent. Potential new sources of flavor violation arise from couplings of quarks and squarks to gauginos. This method has the advantage that, when the off-diagonal elements in the squark mass matrices become large, the method is still valid, unlike the mass insertion which is a perturbation-based expansion. We will return to these considerations when performing specific evaluations and chose the mass eigenstate method for its greater flexibility.

Our paper is organized as follows: In Sec. II we present the general framework of the left-right supersymmetric model. The squark sector is then discussed in Sec. III in both the flavor-diagonal and -nondiagonal scenarios, concentrating especially on the flavor changing mixing between the second and third generations. In Sec. IV, we give the full analytical expressions contributing to $t \to cV$ ($V = g, \gamma, Z$) decays by providing expressions for the gluino, chargino, and neutralino contributions separately. The numerical analysis of the decays is given in Sec. V, where we concentrate on individual and relative contributions of gluinos, charginos, and neutralinos to the branching ratios of the decays. The discussion is carried out in the flavor-diagonal and -nondiagonal scenarios. Section VI is devoted to our summary and conclusion. The chargino and neutralino mass matrices and their diagonalization procedures are summarized in Appendix A. Appendix B includes the relevant Feynman rules.

II. DESCRIPTION OF THE LRSUSY MODEL

The minimal supersymmetric left-right model is based on the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ [19,20]. The matter fields of this model consist of three families of quark and lepton chiral superfields with the following transformations under the gauge group:

$$Q_L^i = \begin{pmatrix} u_L^i \\ d_L^i \end{pmatrix} \sim (3, 2, 1, 1/3), \quad Q_R^i = \begin{pmatrix} d_R^i \\ u_R^i \end{pmatrix} \sim (3, 1, 2, -1/3),$$
$$L_L^i = \begin{pmatrix} \nu_L^i \\ e_L^i \end{pmatrix} \sim (1, 2, 1, -1), \quad L_R^i = \begin{pmatrix} e_R^i \\ \nu_R^i \end{pmatrix} \sim (1, 1, 2, 1), \tag{2.1}$$

where the numbers in the brackets represent the quantum numbers under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The Higgs sector consists of bidoublet and triplet Higgs superfields:

$$\begin{aligned}
\Phi_i &= \begin{pmatrix} \phi_{1i}^0 & \phi_{2i}^+ \\ \phi_{1i}^- & \phi_{2i}^0 \end{pmatrix} \sim (1, 2, 2, 0), \quad (i = u, d),\\
\Delta_L &= \begin{pmatrix} \frac{\Delta_L^-}{\sqrt{2}} & \Delta_L^0 \\ \Delta_L^{--} & -\frac{\Delta_L^-}{\sqrt{2}} \end{pmatrix} \sim (1, 3, 1, -2),\\
\delta_L &= \begin{pmatrix} \frac{\delta_L^+}{\sqrt{2}} & \delta_L^{++} \\ \delta_L^0 & -\frac{\delta_L^+}{\sqrt{2}} \end{pmatrix} \sim (1, 3, 1, 2),\\
\Delta_R &= \begin{pmatrix} \frac{\Delta_R^-}{\sqrt{2}} & \Delta_R^0 \\ \Delta_R^{--} & -\frac{\Delta_R^-}{\sqrt{2}} \end{pmatrix} \sim (1, 1, 3, -2),\\
\delta_R &= \begin{pmatrix} \frac{\delta_R^+}{\sqrt{2}} & \delta_R^{++} \\ \delta_R^0 & -\frac{\delta_R^+}{\sqrt{2}} \end{pmatrix} \sim (1, 1, 3, 2).
\end{aligned} \tag{2.2}$$

The bidoublet Higgs superfields appear in all LRSUSY to implement the $SU(2)_L \times U(1)_Y$ symmetry breaking and to generate a CKM mixing matrix. Supplementary Higgs representations are needed to break left-right symmetry spontaneously: Either doublets or triplets would achieve this, but triplet Higgs bosons Δ_L and Δ_R are chosen to facilitate the seesaw mechanism [21]. In supersymmetry, additional triplet superfields δ_L, δ_R are introduced to cancel triangle gauge anomalies in the fermionic sector. The most general superpotential involving these superfields in LRSUSY is

$$W = \mathbf{Y}_Q^{(i)} Q_L^T \Phi_i i\tau_2 Q_R + \mathbf{Y}_L^{(i)} L_L^T \Phi_i i\tau_2 L_R + \mathbf{Y}_{\rm LR}(L_L^T i\tau_2 \delta_L L_L + L_R^T i\tau_2 \Delta_R L_R) + \mu_{\rm LR}[\mathrm{Tr}(\Delta_L \delta_L + \Delta_R \delta_R)] + \mu_{ij}\,\mathrm{Tr}(i\tau_2 \Phi_i^T i\tau_2 \Phi_j) + W_{\rm NR}, \qquad (2.3)$$

where $\mathbf{Y}_Q$ and $\mathbf{Y}_L$ are the Yukawa couplings for the quarks and leptons, respectively, and $\mathbf{Y}_{\rm LR}$ is the coupling for the triplet Higgs bosons. The parameters μ_{ij} and $\mu_{\rm LR}$ are the Higgs mass parameters. Left-right symmetry requires all $\mathbf{Y}$ matrices to be Hermitian in the generation space and $\mathbf{Y}_{\rm LR}$ matrix to be symmetric. Here $W_{\rm NR}$ denotes (possible) nonrenormalizable terms arising from higher scale physics or Planck scale effects [22]. The presence of these terms insures that, when the SUSY-breaking scale is above M_{W_R}, the ground state is R-parity conserving. In addition, the Lagrangian also includes soft supersymmetry breaking terms as well as F and D terms:

$$\mathcal{L}_{\rm soft} = [\mathbf{A}_Q^i \mathbf{Y}_Q^{(i)} \tilde{Q}_L^T \Phi_i i\tau_2 \tilde{Q}_R + \mathbf{A}_L^i \mathbf{Y}_L^{(i)} \tilde{L}_L^T \Phi_i i\tau_2 \tilde{L}_R + \mathbf{A}_{\rm LR}\mathbf{Y}_{\rm LR}(\tilde{L}_L^T i\tau_2 \delta_L \tilde{L}_L + \tilde{L}_R^T i\tau_2 \Delta_R \tilde{L}_R) + m_\Phi^{(ij)2}\Phi_i^\dagger \Phi_j] + [(m_{L_L}^2)_{ij}\tilde{L}_{Li}^\dagger \tilde{L}_{Lj} + (m_{L_R}^2)_{ij}\tilde{L}_{Ri}^\dagger \tilde{L}_{Rj}] - M_{\rm LR}^2[\mathrm{Tr}(\Delta_R \delta_R) + \mathrm{Tr}(\Delta_L \delta_L) + \mathrm{H.c.}] - [B\mu_{ij}\Phi_i \Phi_j + \mathrm{H.c.}] + [(m_{Q_L}^2)_{ij}\tilde{Q}_{Li}^\dagger \tilde{Q}_{Lj} + (m_{Q_R}^2)_{ij}\tilde{Q}_{Ri}^\dagger \tilde{Q}_{Rj}], \qquad (2.4)$$

where $\mathbf{A}_Q$, $\mathbf{A}_L$, $\mathbf{A}_{\rm LR}$ are trilinear scalar couplings. The LR symmetry is broken spontaneously to $U(1)_{\rm em}$ through nonzero vacuum expectation values (VEV's) of the Higgs fields. These values are

$$\langle \Phi_u \rangle = \begin{pmatrix} \kappa_u & 0 \\ 0 & 0 \end{pmatrix}, \quad \langle \Phi_d \rangle = \begin{pmatrix} 0 & 0 \\ 0 & \kappa_d \end{pmatrix}, \quad \langle \Delta_L \rangle = 0,$$

$$\langle \delta_L \rangle = 0, \quad \langle \Delta_R \rangle = \begin{pmatrix} 0 & v_{\Delta_R} \\ 0 & 0 \end{pmatrix}, \quad \langle \delta_R \rangle = \begin{pmatrix} 0 & 0 \\ v_{\delta_R} & 0 \end{pmatrix},$$

where we already set to zero the CP-violating phase in the mixing of W_L and W_R. 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_{Δ_R}, v_{δ_R} and become much heavier than the SM (left-handed) gauge bosons W_L and Z_L, which in turn pick up masses proportional to κ_u and κ_d at the second stage of breaking.

Fermionic partners of gauge and Higgs bosons mix. In LRSUSY there are six singly charged charginos, corresponding to $\tilde{\lambda}_L$, $\tilde{\lambda}_R$, $\tilde{\phi}_u$, $\tilde{\phi}_d$, $\tilde{\Delta}_L$, and $\tilde{\Delta}_R$. 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 $\tilde{\delta}_R^0$. 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.

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

The interaction of fermions with scalar (Higgs) fields relevant for the quark sector has the following form:

$$\mathcal{L}_Y = \mathbf{Y}_Q^u \bar{Q}_L \Phi_u Q_R + \mathbf{Y}_Q^d \bar{Q}_L \Phi_d Q_R + \mathrm{H.c.} \qquad (2.5)$$

We present below a detailed analysis of flavor violation in the model in both the flavor-diagonal and -nondiagonal case, before proceeding with calculation of the branching ratio of $t \to cV$.

III. FLAVOR CHANGING IN THE LRSUSY

When a basis, called the super-CKM, in which the quark states are diagonal, is used to express the squark mass matrix, flavor-nondiagonal entries naturally arise. So our analysis is focused on the effects of such scenarios in rare top-quark decays in comparison with the flavor-diagonal case.

In the interaction (flavor) basis, $(\tilde{Q}_L^i, \tilde{Q}_R^i)$, the squared-mass matrix for squarks, is

$$\mathcal{M}_f^2 = \begin{pmatrix} m_{f\rm LL}^2 + F_{f\rm LL} + D_{f\rm LL} & (m_{f\rm LR}^2) + F_{f\rm LR} \\ (m_{f\rm LR}^2)^\dagger + F_{f\rm RL} & m_{f\rm RR}^2 + F_{f\rm RR} + D_{f\rm RR} \end{pmatrix}. \qquad (3.1)$$

The F terms are diagonal in the flavor space, $F_{f\rm LL, f\rm RR} = m_f^2$, $(F_{d\rm LR})_{ij} = \mu(m_{d_i}\tan\beta)\mathbf{1}_{ij}$, $(F_{u\rm LR})_{ij} = \mu(m_{u_i}\cot\beta)\mathbf{1}_{ij}$, where $\tan\beta = \kappa_u/\kappa_d$ is defined. The D terms are also flavor-diagonal

$$D_{f\rm LL} = m_Z^2 \cos 2\beta(T_{3f} - Q_f \sin^2\theta_{\rm W})\mathbf{1}_{3\times3}, \quad D_{f\rm RR} = m_Z^2 \cos 2\beta Q_f \sin^2\theta_{\rm W}\mathbf{1}_{3\times3}. \qquad (3.2)$$

In the universal case, one has $(m_{f\rm LL, f\rm RR}^2)_{ij} = m_{\tilde{Q}_{L,R}}^2 \delta_{ij}$, $m_{f\rm LR}^2 = A_f^* m_f$. To reduce the number of free parameters, we also consider the following parameters to be universal: $(m_{\tilde{Q}_{L,R}}^2)_{ij} = \mathrm{M}_{\rm susy}^2 \delta_{ij}$, $A_{d,ij} = A\delta_{ij}$, and $A_{u,ij} = A\delta_{ij}$. Note that, due to the invariance under both $SU(2)_L$ and $SU(2)_R$, $m_{u\rm LL, u\rm RR}^2$ in the up sector cannot be specified independently from the matrix $m_{d\rm LL, d\rm RR}^2$ in the down sector. They are related as $m_{u\rm LL, u\rm RR}^2 = K_{\rm CKM}(m_{d\rm LL, d\rm RR}^2)K_{\rm CKM}^\dagger$, where $K_{\rm CKM}$ is the CKM matrix.[1]

The mass(-squared) matrix in the universal case for the U-type squarks then reduces to, in block form,

[1]For simplicity, we assume $K_{\rm CKM}^L = K_{\rm CKM}^R$, which is a conservative choice and does not require new mixing angles in the right-handed quark matrices [14].

$$\mathcal{M}^2_{U_k} = \begin{pmatrix} \mathrm{M}^2_{\mathrm{susy}} + m^2_{u_k} + m^2_Z(T_{3u} - Q_u \sin^2\theta_{\mathrm{W}})\cos 2\beta & m_{u_k}(A + \mu\cot\beta) \\ m_{u_k}(A + \mu\cot\beta) & \mathrm{M}^2_{\mathrm{susy}} + m^2_{u_k} + m^2_Z Q_u \sin^2\theta_{\mathrm{W}}\cos 2\beta \end{pmatrix},$$

and, for the D-type squarks, to

$$\mathcal{M}^2_{D_k} = \begin{pmatrix} \mathrm{M}^2_{\mathrm{susy}} + m^2_{d_k} + m^2_Z(T_{3d} - Q_d \sin^2\theta_{\mathrm{W}})\cos 2\beta & m_{d_k}(A + \mu\tan\beta) \\ m_{d_k}(A + \mu\tan\beta) & \mathrm{M}^2_{\mathrm{susy}} + m^2_{d_k} + m^2_Z Q_d \sin^2\theta_{\mathrm{W}}\cos 2\beta \end{pmatrix}.$$

The corresponding mass eigenstates are defined as

$$\begin{pmatrix} \tilde{Q}_L \\ \tilde{Q}_R \end{pmatrix} = \begin{pmatrix} \Gamma^\dagger_{QL} \\ \Gamma^\dagger_{QR} \end{pmatrix} \tilde{q}, \tag{3.3}$$

where $\Gamma_{QL,QR}$ are 6×3 mixing matrices and $\tilde{q}$ is a 6×1 column vector.

These up- and down-squark mass matrices are 6×6, but are written above in 2×2 block form. In the flavor-diagonal scenario (constrained LRSUSY), each of these blocks has no nonzero off-diagonal elements. That is, there is no intergenerational mixings for squarks and the only source of flavor mixing comes from the CKM matrix.

The way to induce off-diagonal entries radiatively is to consider the evolution of the squark and quark masses from the SUSY-breaking scale (where both quark and squark mass matrices are flavor-diagonal in the same basis), down to the electroweak scale via renormalization group equations. Experimental bounds from $D^0 - \bar{D}^0$ and $K^0 - \bar{K}^0$ data [23] involving the first generation are tightly constrained. The mixings between the second and the third generations are, on the other hand, free [23]. So in our analysis we assume significant mixing between the second and the third generations in the up- and down-squark mass matrices and neglect those involving the first generation.

The up-squark mass-squared matrix in the $(u_L, c_L, t_L, u_R, c_R, t_R)$ basis can be written as

$$\mathcal{M}^{2,FC}_{U_k} = \mathrm{M}^2_{\mathrm{susy}} \begin{pmatrix} 1 & 0 & 0 & 0 & 0 & 0 \\ 0 & 1 & (\delta^{\mathrm{LL}}_U)_{23} & 0 & (\delta^{\mathrm{LR}}_U)_{22} & (\delta^{\mathrm{LR}}_U)_{23} \\ 0 & (\delta^{\mathrm{LL}}_U)_{32} & 1 & 0 & (\delta^{\mathrm{LR}}_U)_{32} & (\delta^{\mathrm{LR}}_U)_{33} \\ 0 & 0 & 0 & 1 & 0 & 0 \\ 0 & (\delta^{\mathrm{LR}}_U)_{22} & (\delta^{\mathrm{RL}}_U)_{23} & 0 & 1 & (\delta^{\mathrm{RR}}_U)_{23} \\ 0 & (\delta^{\mathrm{RL}}_U)_{32} & (\delta^{\mathrm{LR}}_U)_{33} & 0 & (\delta^{\mathrm{RR}}_U)_{32} & 1 \end{pmatrix}, \tag{3.4}$$

where $(\delta^{\mathrm{LR}}_U)_{22} = m_c(A + \mu\cot\beta)/\mathrm{M}^2_{\mathrm{susy}}$, $(\delta^{\mathrm{LR}}_U)_{33} = m_t(A + \mu\cot\beta)/\mathrm{M}^2_{\mathrm{susy}}$. Here we assume that all diagonal elements on the main diagonal are set to a common value $\mathrm{M}^2_{\mathrm{susy}}$,[2] and the flavor changing off-diagonal entries of each block are furthermore scaled with $\mathrm{M}^2_{\mathrm{susy}}$ to make the parameters dimensionless. Thus, these off-diagonal parameters are defined as

$$(\delta^{\mathrm{LL}}_U)_{ij} = \frac{(m^2_{u\mathrm{LL}})_{ij}}{\mathrm{M}^2_{\mathrm{susy}}}, \quad (\delta^{\mathrm{LR}}_U)_{ij} = \frac{(m^2_{u\mathrm{LR}})_{ij}}{\mathrm{M}^2_{\mathrm{susy}}},$$
$$(\delta^{\mathrm{RL}}_U)_{ij} = \frac{(m^2_{u\mathrm{RL}})_{ij}}{\mathrm{M}^2_{\mathrm{susy}}}, \quad (\delta^{\mathrm{RR}}_U)_{ij} = \frac{(m^2_{u\mathrm{RR}})_{ij}}{\mathrm{M}^2_{\mathrm{susy}}}, \quad i \neq j = 2, 3. \tag{3.5}$$

[2]We comment on the effect of relaxing this condition. Basically, there are non-negligible corrections coming to the 3rd and 6th diagonal entries of both the up- and down-squark matrices. The term to be added to unity in the up sector is $m_t^2/\mathrm{M}^2_{\mathrm{susy}}$, and in the down sector $m_b^2/\mathrm{M}^2_{\mathrm{susy}}$ with a contribution from the D term. So, only in the up sector, the $M_{\mathrm{susy}} \sim m_t$ case can give sizable contributions. However, in the flavor-nondiagonal scenario, we do not consider such small M_{susy} values.

Here we still take the diagonal elements of $m_{u\mathrm{LR},u\mathrm{RL}}$ same as in the flavor-diagonal case. The down-sector case is completely analogous to the up sector.

As mentioned in the introduction, there are two approaches to compute the effects of the flavor changing parameters δ's to physical quantities. The first historical one is the mass insertion formalism [17], and the second and more recent one is the general mass eigenstate formalism [11,18]. In the mass insertion formalism, the δ terms represent mixing between chirality states of different squarks, and it is possible to compute the contributions of the first order flavor changing mass insertions perturbatively if one assumes smallness of the intergenerational mixing elements (δ's) when compared with the diagonal elements. Otherwise, second or even higher order mass insertions have to be taken into account, as is the case for kaon decays [24].

The more recent approach, the general mass eigenstate formalism, is an alternative to the mass insertion formalism. Its advantage is that it allows for large off-diagonal elements where the mass insertion method is no longer applicable. In the general mass eigenstate formalism, the mass matrix in Eq. (3.4) (and the similar one in the down sector) is diagonalized and the flavor changing parameters enter into our expressions through the matrix $\Gamma_{QL,QR}$ defined in Eq. (3.3). So, in the rare top decays $t \to cV$, the new flavor changing neutral currents show themselves in both gluino-squark-quark and neutralino-squark-quark couplings in the up-type squark loops and in the chargino-squark-quark coupling in the down-type squark loop.

After giving expressions for the effective one-loop tcV vertex in the next section, we analyze the flavor-diagonal case first. Then we consider the effect of intergenerational mixings on the rate of the process $t \to cV$. We assume significant mixing between the second and the third generations only, in both the up- and down-squark mass matrices.

IV. $t \to cV$ IN LRSUSY

In this section we present the one-loop LRSUSY effective tcV vertex by considering gluino, chargino, and neutralino loops with up and down squarks, including the effects from left-right and intergenerational squark mixing. The relevant one-loop Feynman diagrams for $t \to cV$ ($V = g, \gamma, Z$) are shown in Fig. 1. Since there are no flavor violating tcV couplings in the Lagrangian, there is no tree-level contribution. Under the assumption that the charm quark mass is negligible, the transition amplitude of the decay is given by

$$\mathcal{M}(t \to cV) = \bar{u}(p-k)A_V^{\mu}u(p)\epsilon_{\mu}(k,\lambda), \tag{4.1}$$

where p and k are the momenta of the top quark and gauge boson V, respectively; $\epsilon(k,\lambda)$ is the polarization vector of the gauge boson. The effective vertex A_V^{μ} can be expressed as

$$A_V^{\mu} = i \sum_{w=\tilde{g},\chi^+,\chi^0}\bigg(\gamma^{\mu}(P_L F_{VL}^{w} + P_R F_{VR}^{w}) + \frac{2p^{\mu}}{m_t}(P_L T_{VL}^{w} + P_R T_{VR}^{w})\bigg), \tag{4.2}$$

where $P_{L,R} = (1 \mp \gamma_5)/2$ is the left (right) chirality projection operator; $F_{VL(R)}$ and $T_{VL(R)}$ are the form factors. In order to compute these form factors, we need to diagonalize the mass matrices for both chargino and neutralino sectors, which are given in Appendix A. Certain Feynman rules of the relevant vertices are presented in Appendix B. The calculations are carried out in the 't Hooft-Feynman gauge in d dimensions. Below we give the explicit expressions of the form factors for each decay. We prefer to use the notation in Ref. [8], which makes the cancellation of infinities more obvious when we set $d = 4$. To put the form factors into compact form, we have further used the identities given in Ref. [9].

A. $t \to cg$ decay

1. Gluino contribution

$$\begin{aligned}
F_{gL}^{\tilde{g}} &= \frac{g_s^3 T^{al}}{8\pi^2}\sum_{k=1}^{6}\bigg[C_2[F]\bigg(\Gamma_{UL}^{k3}\Gamma_{UR}^{*k2}(C_{\epsilon}^{c\tilde{g}[kk]} - C_{\rm SE}^{\tilde{g}[k]}) + \Gamma_{UR}^{k3}\Gamma_{UR}^{*k2}(C_{\rm SEG}^{\tilde{g}[k]} - C_{\rm SEG}^{\tilde{g}[k]}(m_t^2=0))\frac{m_{\tilde{g}}}{m_t}\bigg) \\
&\quad + \frac{C_2[G]}{2}\bigg(\Gamma_{UL}^{k3}\Gamma_{UR}^{*k2}(-C_{\epsilon}^{c\tilde{g}[kk]} + C_{\epsilon}^{d\tilde{g}[k]} + C_{\tilde{g}}^{[k]} + C_{k^2}^{\tilde{g}[k]} + C_{t}^{\tilde{g}[k]}) - \Gamma_{UR}^{k3}\Gamma_{UR}^{*k2}C_{\tilde{g}}^{[k]}\frac{m_t}{m_{\tilde{g}}}\bigg)\bigg], \\
F_{gR}^{\tilde{g}} &= \frac{g_s^3 T^{al}}{8\pi^2}\sum_{k=1}^{6}\bigg[C_2[F]\bigg(\Gamma_{UR}^{k3}\Gamma_{UL}^{*k2}(C_{\epsilon}^{c\tilde{g}[kk]} - C_{\rm SE}^{\tilde{g}[k]}) + \Gamma_{UL}^{k3}\Gamma_{UL}^{*k2}(C_{\rm SEG}^{\tilde{g}[k]} - C_{\rm SEG}^{\tilde{g}[k]}(m_t^2=0))\frac{m_{\tilde{g}}}{m_t}\bigg) \\
&\quad + \frac{C_2[G]}{2}\bigg(\Gamma_{UR}^{k3}\Gamma_{UL}^{*k2}(-C_{\epsilon}^{c\tilde{g}[kk]} + C_{\epsilon}^{d\tilde{g}[k]} + C_{\tilde{g}}^{[k]} + C_{k^2}^{\tilde{g}[k]} + C_{t}^{\tilde{g}[k]}) - \Gamma_{UL}^{k3}\Gamma_{UL}^{*k2}C_{\tilde{g}}^{[k]}\frac{m_t}{m_{\tilde{g}}}\bigg)\bigg], \\
T_{gL}^{\tilde{g}} &= -\frac{g_s^3 T^{al}}{8\pi^2}\sum_{k=1}^{6}\bigg[C_2[F](\Gamma_{UL}^{k3}\Gamma_{UL}^{*k2}C_{\tilde{g}\rm top}^{[kk]} - \Gamma_{UR}^{k3}\Gamma_{UL}^{*k2}C_{\rm top}^{\tilde{g}[kk]}) + \frac{C_2[G]}{2}\bigg(\Gamma_{UL}^{k3}\Gamma_{UL}^{*k2}\bigg(C_{\tilde{g}t}^{[k]} - C_{\tilde{g}\rm top}^{[kk]} - C_{\tilde{g}}^{[k]}\frac{m_t}{m_{\tilde{g}}}\bigg) \\
&\quad + \Gamma_{UR}^{k3}\Gamma_{UL}^{*k2}(C_{t}^{\tilde{g}[k]} + C_{\rm top}^{\tilde{g}[kk]})\bigg)\bigg], \\
T_{gR}^{\tilde{g}} &= -\frac{g_s^3 T^{al}}{8\pi^2}\sum_{k=1}^{6}\bigg[C_2[F](\Gamma_{UR}^{k3}\Gamma_{UR}^{*k2}C_{\tilde{g}\rm top}^{[kk]} - \Gamma_{UL}^{k3}\Gamma_{UR}^{*k2}C_{\rm top}^{\tilde{g}[kk]}) + \frac{C_2[G]}{2}\bigg(\Gamma_{UR}^{k3}\Gamma_{UR}^{*k2}\bigg(C_{\tilde{g}t}^{[k]} - C_{\tilde{g}\rm top}^{[kk]} - C_{\tilde{g}}^{[k]}\frac{m_t}{m_{\tilde{g}}}\bigg) \\
&\quad + \Gamma_{UL}^{k3}\Gamma_{UR}^{*k2}(C_{t}^{\tilde{g}[k]} + C_{\rm top}^{\tilde{g}[kk]})\bigg)\bigg].
\end{aligned} \tag{4.3}$$

2. Chargino contribution

$$
\begin{aligned}
F_{gL}^{\chi^+} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G_{DL}^{*[a,k,2]} - H_{DR}^{*[a,k,2]})\Big[(G_{DL}^{[a,k,3]} - H_{DR}^{[a,k,3]})(C_{\epsilon}^{c\chi^+[akk]} - C_{\mathrm{SE}}^{\chi^+[ak]}) \\
&\quad - (G_{DR}^{[a,k,3]} - H_{DL}^{[a,k,3]})(C_{\mathrm{SEG}}^{\chi^+[ak]} - C_{\mathrm{SEG}}^{\chi^+[ak]}(m_t^2=0))\frac{m_{\chi^+}^a}{m_t}\Big], \\
F_{gR}^{\chi^+} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G_{DR}^{*[a,k,2]} - H_{DL}^{*[a,k,2]})\Big[(G_{DR}^{[a,k,3]} - H_{DL}^{[a,k,3]})(C_{\epsilon}^{c\chi^+[akk]} - C_{\mathrm{SE}}^{\chi^+[ak]}) \\
&\quad - (G_{DL}^{[a,k,3]} - H_{DR}^{[a,k,3]})(C_{\mathrm{SEG}}^{\chi^+[ak]} - C_{\mathrm{SEG}}^{\chi^+[ak]}(m_t^2=0))\frac{m_{\chi^+}^a}{m_t}\Big], \\
T_{gL}^{\chi^+} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G_{DR}^{*[a,k,2]} - H_{DL}^{*[a,k,2]})[(G_{DL}^{[a,k,3]} - H_{DR}^{[a,k,3]})C_{\chi^+\mathrm{top}}^{[akk]} + (G_{DR}^{[a,k,3]} - H_{DL}^{[a,k,3]})C_{\mathrm{top}}^{\chi^+[akk]}], \\
T_{gR}^{\chi^+} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G_{DL}^{*[a,k,2]} - H_{DR}^{*[a,k,2]})[(G_{DR}^{[a,k,3]} - H_{DL}^{[a,k,3]})C_{\chi^+\mathrm{top}}^{[akk]} + (G_{DL}^{[a,k,3]} - H_{DR}^{[a,k,3]})C_{\mathrm{top}}^{\chi^+[akk]}].
\end{aligned}
\tag{4.4}
$$

3. Neutralino contribution

$$
\begin{aligned}
F_{gL}^{\chi^0} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]})\Big[(\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})(C_{\epsilon}^{c\chi^0[nkk]} - C_{\mathrm{SE}}^{\chi^0[nk]}) \\
&\quad + (\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})(C_{\mathrm{SEG}}^{\chi^0[nk]} - C_{\mathrm{SEG}}^{\chi^0[nk]}(m_t^2=0))\frac{m_{\chi^0}^n}{m_t}\Big], \\
F_{gR}^{\chi^0} &= -\frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]})\Big[(\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})(C_{\epsilon}^{c\chi^0[nkk]} - C_{\mathrm{SE}}^{\chi^0[nk]}) \\
&\quad + (\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})(C_{\mathrm{SEG}}^{\chi^0[nk]} - C_{\mathrm{SEG}}^{\chi^0[nk]}(m_t^2=0))\frac{m_{\chi^0}^n}{m_t}\Big], \\
T_{gL}^{\chi^0} &= \frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]})[(\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})C_{\chi^0\mathrm{top}}^{[nkk]} - (\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})C_{\mathrm{top}}^{\chi^0[nkk]}], \\
T_{gR}^{\chi^0} &= \frac{g_s g^2 T^{a\prime}}{16\pi^2}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]})[(\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})C_{\chi^0\mathrm{top}}^{[nkk]} - (\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})C_{\mathrm{top}}^{\chi^0[nkk]}].
\end{aligned}
\tag{4.5}
$$

B. $t \to c\gamma$ decay

1. Gluino contribution

$$
\begin{aligned}
F^{\tilde g}_{\gamma L} &= \frac{g_s^2 g Q_u C_2[F]\sin\theta_{\rm W}}{8\pi^2}\sum_{k=1}^{6}\bigg[\Gamma^{k3}_{UL}\Gamma^{*k2}_{UR}(C^{c\tilde g[kk]}_{\epsilon} - C^{\tilde g[k]}_{\rm SE}) + \Gamma^{k3}_{UR}\Gamma^{*k2}_{UR}(C^{\tilde g[k]}_{\rm SEG} - C^{\tilde g[k]}_{\rm SEG}(m_t^2 = 0))\frac{m_{\tilde g}}{m_t}\bigg],\\
F^{\tilde g}_{\gamma R} &= \frac{g_s^2 g Q_u C_2[F]\sin\theta_{\rm W}}{8\pi^2}\sum_{k=1}^{6}\bigg[\Gamma^{k3}_{UR}\Gamma^{*k2}_{UL}(C^{c\tilde g[kk]}_{\epsilon} - C^{\tilde g[k]}_{\rm SE}) + \Gamma^{k3}_{UL}\Gamma^{*k2}_{UL}(C^{\tilde g[k]}_{\rm SEG} - C^{\tilde g[k]}_{\rm SEG}(m_t^2 = 0))\frac{m_{\tilde g}}{m_t}\bigg],\\
T^{\tilde g}_{\gamma L} &= \frac{g_s^2 g Q_u C_2[F]\sin\theta_{\rm W}}{8\pi^2}\sum_{k=1}^{6}(\Gamma^{k3}_{UR}\Gamma^{*k2}_{UL}C^{\tilde g[kk]}_{\rm top} - \Gamma^{k3}_{UL}\Gamma^{*k2}_{UL}C^{[kk]}_{\tilde g\rm top}),\\
T^{\tilde g}_{\gamma R} &= \frac{g_s^2 g Q_u C_2[F]\sin\theta_{\rm W}}{8\pi^2}\sum_{k=1}^{6}(\Gamma^{k3}_{UL}\Gamma^{*k2}_{UR}C^{\tilde g[kk]}_{\rm top} - \Gamma^{k3}_{UR}\Gamma^{*k2}_{UR}C^{[kk]}_{\tilde g\rm top}).
\end{aligned}
\tag{4.6}
$$

2. Chargino contribution

$$
\begin{aligned}
F^{\chi^+}_{\gamma L} &= -\frac{g^3\sin\theta_{\rm W}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DL} - H^{*[a,k,2]}_{DR})\bigg[(G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR})(Q_d C^{c\chi^+[akk]}_{\epsilon} - Q_u C^{\chi^+[ak]}_{\rm SE} + C^{d\chi^+[aak]}_{\epsilon}\\
&\quad + C^{[aak]}_{\chi^+} + C^{\chi^+[aak]}_{k^2} + C^{\chi^+[aak]}_{t}) + (G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})\bigg(C^{[aak]}_{\chi^+}\frac{m_t}{m^a_{\chi^+}} - Q_u(C^{\chi^+[ak]}_{\rm SEG} - C^{\chi^+[ak]}_{\rm SEG}(m_t^2 = 0))\frac{m^a_{\chi^+}}{m_t}\bigg)\bigg],\\
F^{\chi^+}_{\gamma R} &= -\frac{g^3\sin\theta_{\rm W}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DR} - H^{*[a,k,2]}_{DL})\bigg[(G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})(Q_d C^{c\chi^+[akk]}_{\epsilon} - Q_u C^{\chi^+[ak]}_{\rm SE} + C^{d\chi^+[aak]}_{\epsilon}\\
&\quad + C^{[aak]}_{\chi^+} + C^{\chi^+[aak]}_{k^2} + C^{\chi^+[aak]}_{t}) + (G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR})\bigg(C^{[aak]}_{\chi^+}\frac{m_t}{m^a_{\chi^+}} - Q_u(C^{\chi^+[ak]}_{\rm SEG} - C^{\chi^+[ak]}_{\rm SEG}(m_t^2 = 0))\frac{m^a_{\chi^+}}{m_t}\bigg)\bigg],\\
T^{\chi^+}_{\gamma L} &= -\frac{g^3\sin\theta_{\rm W}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DR} - H^{*[a,k,2]}_{DL})\bigg[(G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})(Q_d C^{\chi^+[akk]}_{\rm top} + C^{\chi^+[aak]}_{t}) + (G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR})\\
&\quad \times\bigg(C^{[aak]}_{\chi^+ t} + Q_d C^{[akk]}_{\chi^+\rm top} - C^{[aak]}_{\chi^+}\frac{m_t}{m^a_{\chi^+}}\bigg)\bigg],\\
T^{\chi^+}_{\gamma R} &= -\frac{g^3\sin\theta_{\rm W}}{16\pi^2}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DL} - H^{*[a,k,2]}_{DR})\bigg[(G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR})(Q_d C^{\chi^+[akk]}_{\rm top} + C^{\chi^+[aak]}_{t}) + (G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})\\
&\quad \times\bigg(C^{[aak]}_{\chi^+ t} + Q_d C^{[akk]}_{\chi^+\rm top} - C^{[aak]}_{\chi^+}\frac{m_t}{m^a_{\chi^+}}\bigg)\bigg].
\end{aligned}
\tag{4.7}
$$

3. Neutralino contribution

$$
\begin{aligned}
F_{\gamma L}^{\chi^0} &= -\frac{g^3 Q_u \sin\theta_{\mathrm{W}}}{16\pi^2} \sum_{k=1}^{6} \sum_{n=1}^{9} (\sqrt{2} G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]}) \bigg[(\sqrt{2} G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})(C_{\epsilon}^{c\chi^0[nkk]} - C_{\mathrm{SE}}^{\chi^0[nk]}) \\
&\quad + (\sqrt{2} G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})(C_{\mathrm{SEG}}^{\chi^0[nk]} - C_{\mathrm{SEG}}^{\chi^0[nk]}(m_t^2 = 0)) \frac{m_{\chi^0}^n}{m_t} \bigg], \\
F_{\gamma R}^{\chi^0} &= -\frac{g^3 Q_u \sin\theta_{\mathrm{W}}}{16\pi^2} \sum_{k=1}^{6} \sum_{n=1}^{9} (\sqrt{2} G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]}) \bigg[(\sqrt{2} G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})(C_{\epsilon}^{c\chi^0[nkk]} - C_{\mathrm{SE}}^{\chi^0[nk]}) \\
&\quad + (\sqrt{2} G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})(C_{\mathrm{SEG}}^{\chi^0[nk]} - C_{\mathrm{SEG}}^{\chi^0[nk]}(m_t^2 = 0)) \frac{m_{\chi^0}^n}{m_t} \bigg], \\
T_{\gamma L}^{\chi^0} &= \frac{g^3 Q_u \sin\theta_{\mathrm{W}}}{16\pi^2} \sum_{k=1}^{6} \sum_{n=1}^{9} (\sqrt{2} G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]}) [(\sqrt{2} G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]}) C_{\chi^0 \mathrm{top}}^{[nkk]} - (\sqrt{2} G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]}) C_{\mathrm{top}}^{\chi^0[nkk]}], \\
T_{\gamma R}^{\chi^0} &= \frac{g^3 Q_u \sin\theta_{\mathrm{W}}}{16\pi^2} \sum_{k=1}^{6} \sum_{n=1}^{9} (\sqrt{2} G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]}) [(\sqrt{2} G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]}) C_{\chi^0 \mathrm{top}}^{[nkk]} - (\sqrt{2} G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]}) C_{\mathrm{top}}^{\chi^0[nkk]}].
\end{aligned}
\tag{4.8}
$$

C. $t \to cZ$ decay

1. Gluino contribution

$$
\begin{aligned}
F_{ZL}^{\tilde{g}} &= -\frac{g_s^2 g C_2[F]}{8\pi^2 \cos\theta_{\mathrm{W}}} \sum_{k=1}^{6} \bigg[\bigg(\Gamma_{UL}^{k3} \Gamma_{UR}^{*k2} C_{\mathrm{SE}}^{\tilde{g}[k]} - \Gamma_{UR}^{k3} \Gamma_{UR}^{*k2} (C_{\mathrm{SEG}}^{\tilde{g}[k]} - C_{\mathrm{SEG}}^{\tilde{g}[k]}(m_t^2 = 0)) \frac{m_{\tilde{g}}}{m_t} \bigg) (T_{3u} - Q_u \sin^2\theta_{\mathrm{W}}) \\
&\quad - \sum_{h=1}^{6} \Gamma_{UL}^{h3} \Gamma_{UR}^{*k2} C_{\epsilon}^{c\tilde{g}[hk]} \bigg(T_{3u} \sum_{j=1}^{3} \Gamma_{UL}^{*hj} \Gamma_{UL}^{kj} - Q_u \delta^{hk} \sin^2\theta_{\mathrm{W}} \bigg) \bigg], \\
F_{ZR}^{\tilde{g}} &= \frac{g_s^2 g C_2[F]}{8\pi^2 \cos\theta_{\mathrm{W}}} \sum_{k=1}^{6} \bigg[\bigg(\Gamma_{UR}^{k3} \Gamma_{UL}^{*k2} C_{\mathrm{SE}}^{\tilde{g}[k]} - \Gamma_{UL}^{k3} \Gamma_{UL}^{*k2} (C_{\mathrm{SEG}}^{\tilde{g}[k]} - C_{\mathrm{SEG}}^{\tilde{g}[k]}(m_t^2 = 0)) \frac{m_{\tilde{g}}}{m_t} \bigg) Q_u \sin^2\theta_{\mathrm{W}} \\
&\quad + \sum_{h=1}^{6} \Gamma_{UR}^{h3} \Gamma_{UL}^{*k2} C_{\epsilon}^{c\tilde{g}[hk]} \bigg(T_{3u} \sum_{j=1}^{3} \Gamma_{UL}^{*hj} \Gamma_{UL}^{kj} - Q_u \delta^{hk} \sin^2\theta_{\mathrm{W}} \bigg) \bigg], \\
T_{ZL}^{\tilde{g}} &= -\frac{g_s^2 g C_2[F]}{8\pi^2 \cos\theta_{\mathrm{W}}} \sum_{k=1}^{6} \sum_{h=1}^{6} (\Gamma_{UL}^{h3} \Gamma_{UL}^{*k2} C_{\tilde{g}\mathrm{top}}^{[hk]} - \Gamma_{UR}^{h3} \Gamma_{UL}^{*k2} C_{\mathrm{top}}^{\tilde{g}[hk]}) \bigg(T_{3u} \sum_{j=1}^{3} \Gamma_{UL}^{*hj} \Gamma_{UL}^{kj} - Q_u \delta^{hk} \sin^2\theta_{\mathrm{W}} \bigg), \\
T_{ZR}^{\tilde{g}} &= -\frac{g_s^2 g C_2[F]}{8\pi^2 \cos\theta_{\mathrm{W}}} \sum_{k=1}^{6} \sum_{h=1}^{6} (\Gamma_{UR}^{h3} \Gamma_{UR}^{*k2} C_{\tilde{g}\mathrm{top}}^{[hk]} - \Gamma_{UL}^{h3} \Gamma_{UR}^{*k2} C_{\mathrm{top}}^{\tilde{g}[hk]}) \bigg(T_{3u} \sum_{j=1}^{3} \Gamma_{UL}^{*hj} \Gamma_{UL}^{kj} - Q_u \delta^{hk} \sin^2\theta_{\mathrm{W}} \bigg).
\end{aligned}
\tag{4.9}
$$

2. Chargino contribution

$$\begin{aligned}
F^{\chi^+}_{ZL} &= \frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DL} - H^{*[a,k,2]}_{DR})\Bigg[\Bigg((G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR})C^{\chi^+[ak]}_{\mathrm{SE}} + (G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})(C^{\chi^+[ak]}_{\mathrm{SEG}} \\
&\quad - C^{\chi^+[ak]}_{\mathrm{SEG}}(m_t^2 = 0))\frac{m^a_{\chi^+}}{m_t}\Bigg)(T_{3u} - Q_u\sin^2\theta_W) - \sum_{h=1}^{6}(G^{[a,h,3]}_{DL} - H^{[a,h,3]}_{DR})C^{c\chi^+[ahk]}_{\epsilon} \\
&\quad \times\Bigg(T_{3d}\sum_{j=1}^{3}\Gamma^{*hj}_{DL}\Gamma^{kj}_{DL} - Q_d\delta^{hk}\sin^2\theta_W\Bigg) + \sum_{b=1}^{5}\Bigg((G^{[b,k,3]}_{DL} - H^{[b,k,3]}_{DR})[(C^{d\chi^+[abk]}_{\epsilon} + C^{\chi^+[abk]}_{k^2} - C^{\chi^+[abk]}_{t})O^{ab}_R \\
&\quad + C^{[abk]}_{\chi^+}O^{ab}_L] + (G^{[b,k,3]}_{DR} - H^{[b,k,3]}_{DL})\Bigg(-C'^{[aak]}_{\chi^+t}O^{ab}_L + C'^{[bbk]}_{\chi^+t}O^{ab}_R + C^{[abk]}_{\chi^+}O^{ab}_L\frac{m_t}{m^b_{\chi^+}}\Bigg)\Bigg)\Bigg], \\
F^{\chi^+}_{ZR} &= -\frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DR} - H^{*[a,k,2]}_{DL})\Bigg[\Bigg((G^{[a,k,3]}_{DR} - H^{[a,k,3]}_{DL})C^{\chi^+[ak]}_{\mathrm{SE}} + (G^{[a,k,3]}_{DL} - H^{[a,k,3]}_{DR}) \\
&\quad \times(C^{\chi^+[ak]}_{\mathrm{SEG}} - C^{\chi^+[ak]}_{\mathrm{SEG}}(m_t^2 = 0))\frac{m^a_{\chi^+}}{m_t}\Bigg)Q_u\sin^2\theta_W + \sum_{h=1}^{6}(G^{[a,h,3]}_{DR} - H^{[a,h,3]}_{DL})C^{c\chi^+[ahk]}_{\epsilon} \\
&\quad \times\Bigg(T_{3d}\sum_{j=1}^{3}\Gamma^{*hj}_{DL}\Gamma^{kj}_{DL} - Q_d\delta^{hk}\sin^2\theta_W\Bigg) - \sum_{b=1}^{5}\Bigg((G^{[b,k,3]}_{DR} - H^{[b,k,3]}_{DL})[(C^{d\chi^+[abk]}_{\epsilon} + C^{\chi^+[abk]}_{k^2} - C^{\chi^+[abk]}_{t})O^{ab}_L \\
&\quad + C^{[abk]}_{\chi^+}O^{ab}_R] + (G^{[b,k,3]}_{DL} - H^{[b,k,3]}_{DR})\Bigg(-C'^{[aak]}_{\chi^+t}O^{ab}_R + C'^{[bbk]}_{\chi^+t}O^{ab}_L + C^{[abk]}_{\chi^+}O^{ab}_R\frac{m_t}{m^b_{\chi^+}}\Bigg)\Bigg)\Bigg], \\
T^{\chi^+}_{ZL} &= -\frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DR} - H^{*[a,k,2]}_{DL})\Bigg[\sum_{h=1}^{6}((G^{[a,h,3]}_{DL} - H^{[a,h,3]}_{DR})C^{[ahk]}_{\chi^+\mathrm{top}} + (G^{[a,h,3]}_{DR} - H^{[a,h,3]}_{DL})C^{\chi^+[ahk]}_{\mathrm{top}}) \\
&\quad \times\Bigg(T_{3d}\sum_{j=1}^{3}\Gamma^{*hj}_{DL}\Gamma^{kj}_{DL} - Q_d\delta^{hk}\sin^2\theta_W\Bigg) - \sum_{b=1}^{5}\Bigg((G^{[b,k,3]}_{DR} - H^{[b,k,3]}_{DL})O^{ab}_L + (G^{[b,k,3]}_{DL} - H^{[b,k,3]}_{DR})(-C'^{[bbk]}_{\chi^+t}O^{ab}_L \\
&\quad + \Bigg(C'^{[aak]}_{\chi^+t} + C^{[aak]}_{\chi^+t} - C^{[abk]}_{\chi^+}\frac{m_t}{m^b_{\chi^+}}\Bigg)O^{ab}_L\Bigg)\Bigg)\Bigg], \\
T^{\chi^+}_{ZR} &= -\frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{a=1}^{5}(G^{*[a,k,2]}_{DL} - H^{*[a,k,2]}_{DR})\Bigg[\sum_{h=1}^{6}((G^{[a,h,3]}_{DR} - H^{[a,h,3]}_{DL})C^{[ahk]}_{\chi^+\mathrm{top}} + (G^{[a,h,3]}_{DL} - H^{[a,h,3]}_{DR})C^{\chi^+[ahk]}_{\mathrm{top}}) \\
&\quad \times\Bigg(T_{3d}\sum_{j=1}^{3}\Gamma^{*hj}_{DL}\Gamma^{kj}_{DL} - Q_d\delta^{hk}\sin^2\theta_W\Bigg) - \sum_{b=1}^{5}\Bigg((G^{[b,k,3]}_{DL} - H^{[b,k,3]}_{DR})O^{ab}_R + (G^{[b,k,3]}_{DR} - H^{[b,k,3]}_{DL}) \\
&\quad \times(-C'^{[bbk]}_{\chi^+t}O^{ab}_R + \Bigg(C'^{[aak]}_{\chi^+t} + C^{[aak]}_{\chi^+t} - C^{[abk]}_{\chi^+}\frac{m_t}{m^b_{\chi^+}}\Bigg)O^{ab}_R\Bigg)\Bigg)\Bigg].
\end{aligned} \tag{4.10}$$

3. Neutralino contribution

$$
\begin{aligned}
F_{ZL}^{\chi^0} &= \frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]})\bigg[\bigg((\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})C_{\rm SE}^{\chi^0[nk]} - (\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})\\
&\quad \times (C_{\rm SEG}^{\chi^0[nk]} - C_{\rm SEG}^{\chi^0[nk]}(m_t^2=0))\frac{m_{\chi^0}^n}{m_t}\bigg)(T_{3u} - Q_u\sin^2\theta_W) - \sum_{h=1}^{6}(\sqrt{2}G_{UL}^{[n,h,3]} + H_{UR}^{[n,h,3]})C_\epsilon^{c\chi^0[nhk]}\\
&\quad \times\bigg(T_{3u}\sum_{j=1}^{3}\Gamma_{UL}^{*hj}\Gamma_{UL}^{kj} - Q_u\delta^{hk}\sin^2\theta_W\bigg) + \sum_{m=1}^{9}\bigg((\sqrt{2}G_{UL}^{[m,k,3]} + H_{UR}^{[m,k,3]})[(C_\epsilon^{d\chi^0[nmk]} + C_{k^2}^{\chi^0[nmk]} - C_t^{\chi^0[nmk]})O_R^{\prime nm}\\
&\quad + C_{\chi^0}^{[nmk]}O_L^{\prime nm}] + (\sqrt{2}G_{UR}^{[m,k,3]} - H_{UL}^{[m,k,3]})\bigg(C_{\chi^0 t}^{\prime[nnk]}O_L^{\prime nm} - C_{\chi^0 t}^{\prime[mmk]}O_R^{\prime nm} - C_{\chi^0}^{[nmk]}O_L^{\prime nm}\frac{m_t}{m_{\chi^0}^m}\bigg)\bigg)\bigg],\\
F_{ZR}^{\chi^0} &= -\frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]})\bigg[\bigg((\sqrt{2}G_{UR}^{[n,k,3]} - H_{UL}^{[n,k,3]})C_{\rm SE}^{\chi^0[nk]} - (\sqrt{2}G_{UL}^{[n,k,3]} + H_{UR}^{[n,k,3]})\\
&\quad (C_{\rm SEG}^{\chi^0[nk]} - C_{\rm SEG}^{\chi^0[nk]}(m_t^2=0))\frac{m_{\chi^0}^n}{m_t}\bigg)Q_u\sin^2\theta_W + \sum_{h=1}^{6}(\sqrt{2}G_{UR}^{[n,h,3]} - H_{UL}^{[n,h,3]})C_\epsilon^{c\chi^0[nhk]}\\
&\quad \times\bigg(T_{3u}\sum_{j=1}^{3}\Gamma_{UL}^{*hj}\Gamma_{UL}^{kj} - Q_u\delta^{hk}\sin^2\theta_W\bigg) - \sum_{m=1}^{9}\bigg((\sqrt{2}G_{UR}^{[m,k,3]} - H_{UL}^{[m,k,3]})[(C_\epsilon^{d\chi^0[nmk]} + C_{k^2}^{\chi^0[nmk]} - C_t^{\chi^0[nmk]})O_L^{\prime nm}\\
&\quad + C_{\chi^0}^{[nmk]}O_R^{\prime nm}] + (\sqrt{2}G_{UL}^{[m,k,3]} + H_{UR}^{[m,k,3]})\bigg(C_{\chi^0 t}^{\prime[nnk]}O_R^{\prime nm} - C_{\chi^0 t}^{\prime[mmk]}O_L^{\prime nm} - C_{\chi^0}^{[nmk]}O_R^{\prime nm}\frac{m_t}{m_{\chi^0}^m}\bigg)\bigg)\bigg],\\
T_{ZL}^{\chi^0} &= \frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UR}^{*[n,k,2]} - H_{UL}^{*[n,k,2]})\bigg[\sum_{h=1}^{6}((\sqrt{2}G_{UL}^{[n,h,3]} + H_{UR}^{[n,h,3]})C_{\chi^0{\rm top}}^{[nhk]}\\
&\quad - (\sqrt{2}G_{UR}^{[n,h,3]} - H_{UL}^{[n,h,3]})C_{\rm top}^{\chi^0[nhk]})\bigg(T_{3u}\sum_{j=1}^{3}\Gamma_{UL}^{*hj}\Gamma_{UL}^{kj} - Q_u\delta^{hk}\sin^2\theta_W\bigg) + \sum_{m=1}^{9}\bigg((\sqrt{2}G_{UR}^{[m,k,3]} - H_{UL}^{[m,k,3]})O_L^{\prime nm}\\
&\quad + (\sqrt{2}G_{UL}^{[m,k,3]} + H_{UR}^{[m,k,3]})\bigg(C_{\chi^0 t}^{\prime[mmk]}O_L^{\prime nm} - \bigg(C_{\chi^0 t}^{\prime[nnk]} + C_{\chi^0 t}^{[nnk]} - C_{\chi^0}^{[nmk]}\frac{m_t}{m_{\chi^0}^m}\bigg)O_R^{\prime nm}\bigg)\bigg)\bigg],\\
T_{ZR}^{\chi^0} &= \frac{g^3}{16\pi^2\cos\theta_W}\sum_{k=1}^{6}\sum_{n=1}^{9}(\sqrt{2}G_{UL}^{*[n,k,2]} + H_{UR}^{*[n,k,2]})\bigg[\sum_{h=1}^{6}((\sqrt{2}G_{UR}^{[n,h,3]} - H_{UL}^{[n,h,3]})C_{\chi^0{\rm top}}^{[nhk]} - (\sqrt{2}G_{UL}^{[n,h,3]} + H_{UR}^{[n,h,3]})C_{\rm top}^{\chi^0[nhk]})\\
&\quad \times\bigg(T_{3u}\sum_{j=1}^{3}\Gamma_{UL}^{*hj}\Gamma_{UL}^{kj} - Q_u\delta^{hk}\sin^2\theta_W\bigg) + \sum_{m=1}^{9}\bigg((\sqrt{2}G_{UL}^{[m,k,3]} + H_{UR}^{[m,k,3]})O_R^{\prime nm} + (\sqrt{2}G_{UR}^{[m,k,3]} - H_{UL}^{[m,k,3]})\\
&\quad \times\bigg(C_{\chi^0 t}^{\prime[mmk]}O_R^{\prime nm} - \bigg(C_{\chi^0 t}^{\prime[nnk]} + C_{\chi^0 t}^{[nnk]} - C_{\chi^0}^{[nmk]}\frac{m_t}{m_{\chi^0}^m}\bigg)O_L^{\prime nm}\bigg)\bigg)\bigg],
\end{aligned}
\tag{4.11}
$$

$$
\begin{aligned}
&C_{\epsilon}^{cw[ahk]} = \int_0^1 dx \int_0^{1-x} dy \left(\frac{1}{\epsilon} - \gamma - \log\frac{L^c_{ahk}(w)}{4\pi\Lambda^2}\right), \qquad C_{\rm top}^{w[ahk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_t^2 x(1-x-y)}{L^c_{ahk}(w)},\\
&C_{w{\rm top}}^{[ahk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_t m_w^a (1-x-y)}{L^c_{ahk}(w)}, \qquad C_{\rm SE}^{w[ak]} = \int_0^1 dx x \left(\frac{1}{\epsilon} - \gamma - \log\frac{L^{\rm self}_{ak}(w)}{4\pi\Lambda^2}\right),\\
&C_{\rm SEG}^{w[ak]} = \int_0^1 dx \left(\frac{1}{\epsilon} - \gamma - \log\frac{L^{\rm self}_{ak}(w)}{4\pi\Lambda^2}\right), \qquad C_{\epsilon}^{dw[abk]} = \int_0^1 dx \int_0^{1-x} dy \left(\frac{1}{\epsilon} - \gamma - 1 - \log\frac{L^d_{abk}(w)}{4\pi\Lambda^2}\right),\\
&C_{w}^{[abk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_w^a m_w^b}{L^d_{abk}(w)}, \qquad C_{k^2}^{w[abk]} = \int_0^1 dx \int_0^{1-x} dy \frac{k^2 xy}{L^d_{abk}(w)},\\
&C_{t}^{w[abk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_t^2 x(1-x-y)}{L^d_{abk}(w)}, \qquad C_{wt}^{[abk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_w^a m_t (1-x-y)}{L^d_{abk}(w)},\\
&C_{wt}^{\prime[abk]} = \int_0^1 dx \int_0^{1-x} dy \frac{m_w^a m_t x}{L^d_{abk}(w)}, \qquad w = \tilde{g}, \gamma, Z,
\end{aligned}
\tag{4.12}
$$

$$
\begin{aligned}
L^{\rm self}_{ak}(w) &= -m_t^2 x(1-x) + (m_w^a)^2 (1-x) + m^2_{\tilde{q}_k} x,\\
L^{c}_{ahk}(w) &= -m_t^2 x(1-x-y) + (m_w^a)^2 (1-x-y) + m^2_{\tilde{q}_h} x + m^2_{\tilde{q}_k} y - k^2 xy,\\
L^{d}_{abk}(w) &= -m_t^2 x(1-x-y) + (m_w^a)^2 y + (m_w^b)^2 x + m^2_{\tilde{q}_k}(1-x-y) - k^2 xy.
\end{aligned}
\tag{4.13}
$$

Here $\epsilon = 2 - d/2$ (d is the number of dimensions); $C_2[F] = 4/3 (C_2[G] = 3)$ is the quadratic Casimir operator of the fundamental (adjoint) representation of $SU(3)_C$ with $\sum_{a\prime}(T^{a\prime}T^{a\prime})_{ij} = C_2[F]\delta_{ij}$ [$T^{a\prime}$ are the $SU(2)$ generators in the fundamental representation with normalization $\mathrm{Tr}(T^{a\prime}T^{b\prime}) = \delta^{a\prime b\prime}/2$]. The indices a, b run over the intermediate chargino mass eigenstates from 1 to 5, while n, m run over the intermediate neutralino mass eigenstates from 1 to 9. The indices k, h are used to represent squarks mass eigenstates, which run from 1 to 6. The arbitrary parameter Λ is introduced in Eq. (4.12) to make the argument of the logarithm function dimensionless and our results are independent of Λ. For $w = \tilde{g}$, throughout the Eqs. (4.3), (4.4), (4.5), (4.6), (4.7), (4.8), (4.9), (4.10), (4.11), (4.12), and (4.13), we dropped the indices a, b or n, m carried by the C functions defined above with a generic w dependency, as gluinos are assumed to have identical masses.

There are two constraints on the form factors that one can check. The first one is the cancellations of infinities when we set $d \to 4 (\epsilon \to 0)$ limit and the second one is the gauge invariance requirement that the coefficient of the γ^μ term in the $t \to cg, c\gamma$ decays should vanish in the $k^2 \to 0$ limit. Obviously, this is not the case for the $t \to cZ$ mode.

Cancellation of divergencies can be realized in two ways in d dimension. Either one can adopt the on-shell renormalization scheme (see, for example, [25]), where there is no contribution from self-energy diagrams [Figs. 1(a) and 1(b)], and find counterterms which exactly cancel the $1/\epsilon$ terms, or include all possible diagrams contributing to the process and expect to have overall cancellation of $1/\epsilon$. We followed the latter. The explicit ϵ dependencies are included inside the C functions defined in Eq. (4.12). As mentioned before, this way of writing the form factors has an advantage over other formulations, such as Passarino-Veltman, since it makes the cancellation of $1/\epsilon$ terms more apparent. In the expressions given above, there are, for example, some combinations of C functions appearing repeatedly in which cancellation of $1/\epsilon$ terms requires no computation at all. Note that by using the Gordon identity one can decompose Eq. (4.2) into an alternative form which contains the Dirac structures γ^μ and $\sigma^{\mu\nu}k_\nu$, where $\sigma^{\mu\nu} = (i/2)[\gamma^\mu, \gamma^\nu]$. So there will be a γ^μ component coming from the $2p^\mu$ part, and, while we are checking the cancellation of divergencies, the combination $F^w_{VL(R)} + T^w_{VR(L)}$ should be considered. However, none of the $T^w_{VR(L)}$ form factors has $1/\epsilon$ dependency, so this point is going to

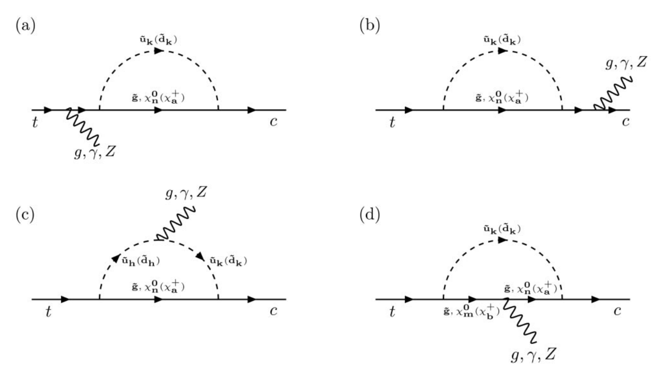

FIG. 1. The one-loop Feynman diagrams contributing to $t \to cg, \gamma, Z$ including gluino, chargino, and neutralino loops. For each decay mode, not all four types of diagrams contribute. In the $t \to cg$ case, diagram (d) for the chargino and the neutralino loops does not contribute. In the $t \to c\gamma$ case, diagram (d) for the gluino and the neutralino loops does not contribute. In the $t \to cZ$ case, diagram (d) for the gluino loop does not contribute.

be important only in the discussion of gauge invariance. Cancellation of $1/\epsilon$ happens in a nontrivial way but is relatively straightforward to check for the $t \to cg$ and $t \to c\gamma$ decay modes with respect to the case for the $t \to cZ$ decay, which requires use of some further properties of $\Gamma_{QL,QR}$,

$$\sum_{i=1}^{3}(\Gamma_{QL}^{*hi}\Gamma_{QL}^{ki} + \Gamma_{QR}^{*hi}\Gamma_{QR}^{ki}) = \delta^{hk}, \qquad \sum_{k=1}^{6}\Gamma_{QL,QR}^{*ki}\Gamma_{QL,QR}^{kj} = \delta^{ij}, \tag{4.14}$$

with unitarity properties of the (U^*, V) and N' matrices defined in Appendix A in the chargino and neutralino sectors, respectively.

Gauge invariance requires vanishing the γ^μ terms for massless gauge bosons. Thus, we should expect to have $F_{VL(R)}^w + T_{VR(L)}^w = 0$ for only $V = g, \gamma$ in each $w = \tilde{g}, \chi^+, \chi^0$ cases. Using the identities given in Ref. [8] for $w = \tilde{g}$ together with the similar ones for $w = \chi^+$ and χ^0, the vanishing of $F_{VL(R)}^w + T_{VR(L)}^w$ for each w is straightforward.

The branching ratio of the decay $t \to cV$ can be given, neglecting the mass of the charm quark

$$\begin{aligned}\mathrm{BR}(t \to cZ) &= \frac{m_t(1-\eta)^2}{32\pi\Gamma_t(t \to bW)}\sum_{w=\tilde{g},\chi^+,\chi^0}\Big[\Big(\eta - 2 + \frac{1}{\eta}\Big)(|T_{ZL}^w|^2 + |T_{ZR}^w|^2) + 2\Big(-1 + \frac{1}{\eta}\Big)\mathrm{Re}[T_{ZL}^w F_{ZR}^{*w} + T_{ZR}^{*w}F_{ZL}^w] \\ &\quad + \Big(2 + \frac{1}{\eta}\Big)(|F_{ZL}^w|^2 + |F_{ZR}^w|^2)\Big], \\ \mathrm{BR}(t \to cV) &= \frac{m_t}{16\pi\Gamma_t(t \to bW)}\sum_{w=\tilde{g},\chi^+,\chi^0}(|T_{VL}^w|^2 + |T_{VR}^w|^2), \qquad V = g, \gamma,\end{aligned} \tag{4.15}$$

where $\eta = m_Z^2/m_t^2$ and $\Gamma_t(t \to bW)$ is taken as the total decay width of the top quark.

V. NUMERICAL RESULTS

In this section, we present some numerical results for the branching ratio (BR) of the decays $t \to cV$ ($V = g, \gamma, Z$) including the contributions from gluino-up-squark, chargino-down-squark, and neutralino-up-squark loops within the context of both flavor-diagonal and -nondiagonal scenarios. We prefer to discuss each contribution separately, as this is missing from the literature. Within the framework we have chosen, there is a large set of parameters that need to be fixed. Throughout our numerical computations, the SM parameters are taken as $m_b = 4.5$ GeV, $m_t = 173.5$ GeV, $m_W = 80.425$ GeV, $m_Z = 91.187$ GeV, $\alpha_s(m_Z) = 0.1172$, and $\sin^2\theta_W = 0.2312$. The elements of the CKM matrix are taken as $K_{\rm CKM}^{ui} = (0.973, 0.220, 3.67 \times 10^{-3})$, $K_{\rm CKM}^{ci} = (-0.224, 0.986, 41.3 \times 10^{-3})$, $K_{\rm CKM}^{ti} = (0.01, -0.05, 0.997)$. The free SUSY parameters in the LR symmetric framework[3] are M_R, g_R, v_{Δ_R}, and v_{δ_R} in addition to the usual ones, M_L, M_V, $M_{\rm susy}$, A, μ, and $\tan\beta$ (we assume bilinear and trilinear scalar couplings in the soft symmetry breaking Lagrangian to be flavor-diagonal, $\mu_{ij} \equiv \mu$, $A_{ij} = A$). The mass of the gluino needs to be fixed as well. In addition to these, there are new flavor changing parameters, $(\delta_{U(D)}^{AB})_{ij}$, $i \neq j = 2, 3$, $A, B = L, R$ appearing in both the up-type and down-type squark mass-squared matrices defined in Eq. (3.4). These are relevant only in the flavor-nondiagonal case and we vary them in (0, 1) interval.

The 5×5 chargino and 9×9 neutralino mass matrices are diagonalized numerically and the following relations are assumed: $M_L = M_R$, $M_V = M_L/2$, $g_R = g_L$ and we set $M_L = 150$ GeV, $v_{\Delta_R} = v_{\delta_R} = 1$ TeV, $\mu = 200$ GeV, and $\tan\beta = 10$ throughout our analysis.[4] There are experimental lower bounds on the masses of charginos and neutralinos as well as gluinos and squarks [26]. For the values of the parameters given above, the lightest chargino is around 130 GeV and the lightest neutralino 90 GeV, which are consistent with the experimental lower bounds [26]. To determine the squark masses, we need to further fix $M_{\rm susy}$ and the coupling A. We set $A = M_{\rm susy}$ and vary them within (100, 1000) GeV in the flavor-diagonal case and set $M_{\rm susy} = 300, 400, 1000$ GeV in the flavor-nondiagonal case. For example, the lightest up-type (down-type) squark mass is around 250 (290) GeV for $M_{\rm susy} = 300$ GeV in the flavor-diagonal context. When the flavor changing parameters are turned on, the lightest up-type (down-type) squark is around 100 (200) GeV for intermediate values of $\delta_U^{\rm RR,LL} = \delta_U^{\rm RL,LR}$ ($\delta_D^{\rm RR,LL} = \delta_D^{\rm RL,LR}$). When the flavor violation comes only from the RR or the LL sector, we are getting slightly larger squark mass values. We now discuss in the following subsections the flavor-diagonal and -nondiagonal cases separately.

[3]The LRSUSY is assumed not embedded into some larger supersymmetric grand unified theories such as $SO(10)$, E_6. Otherwise, one can relate some of these parameters by using the relations among them at the unification scale with the use of renormalization group equations. Such a framework, however, would lead to more complicated particle spectra.

[4]We also considered large $\tan\beta$ values and discuss it at the end of Sec. V B.

A. Flavor-diagonal case

We discuss first the flavor-diagonal case and analyze the flavor changing effects based on the flavor-diagonal results. We assume that the SM contribution to the $t \to cV$ ($V = g, \gamma, Z$) decays is negligible with respect to the contributions from the LRSUSY as the total branching ratio is concerned. In some cases, chargino or neutralino loop contributions might be comparable with the SM values. However, we concentrate mainly on the relative contributions from each loop involving supersymmetric particles. In all our considerations, the gluino loop dominates both chargino and neutralino loops.

We first investigate the $t \to cg$ branching ratio as a function of M_{susy} within the range $(100, 1000)$ GeV for $\tan\beta = 10$ and $m_{\tilde{g}} = 300$ GeV. Figure 2 shows the dependency including the gluino, chargino, and neutralino contributions separately. Neutralino dominates the chargino contribution for the most part of the parameter space (except for large M_{susy} values, where they are roughly equal). There is a 2 orders of magnitude difference between them for $M_{\text{susy}} \sim 200$ GeV which gets smaller as M_{susy} gets larger. The gluino contribution is around 3 orders of magnitude larger than the neutralino contribution for $M_{\text{susy}} \sim 200$, and the difference gets even bigger, reaching 4×10^{-6} when M_{susy} is around 100 GeV. Overall, both the neutralino and the chargino contributions are practically negligible. In general, the neutralino contribution is bigger than the chargino contribution in the flavor-diagonal scenario for $t \to cV$ (with the notable exception of the $t \to cZ$ decay mode and the $t \to c\gamma$ decay for certain M_{susy} values), and this is consistent with the findings of Lopez *et al.* [9]. Note that this picture is reversed for the rare decays involving down-type quarks such as bottom-strange quark transitions (b decays). There the chargino contribution is always larger than the corresponding neutralino one. In b decays, the chargino couples with up-type squarks, while here it couples with down-type squarks, whereas the neutralino does the opposite. So the features seen here are consistent with the characteristics of the down-type rare decays [14].

In Fig. 3, we present the M_{susy} dependency of the branching ratio of the decay $t \to c\gamma$ for the same parameter values chosen above. The gluino contribution is 1 to 2 orders of magnitude smaller than the one in the $t \to cg$ decay and it reaches 2×10^{-7} at the maximum level. The neutralino contribution is 2 to 4 orders of magnitude smaller than the gluino. Even though, after $M_{\text{susy}} \sim 300$ GeV, the chargino contribution is larger than the neutralino, both are still negligible in the entire parameter space considered.

The final figure of this subsection, Fig. 4, shows the same dependency but for $t \to cZ$ decay. The general pattern is the same except that chargino dominates neutralino everywhere in the $(100, 1000)$ GeV range for M_{susy} (1 order of magnitude larger than the neutralino), and, furthermore, the chargino loop contribution is almost constant as M_{susy} changes. The gluino contribution takes maximum values of 7×10^{-7} but decreases sharply as M_{susy} gets bigger and even becomes comparable to the chargino contribution.

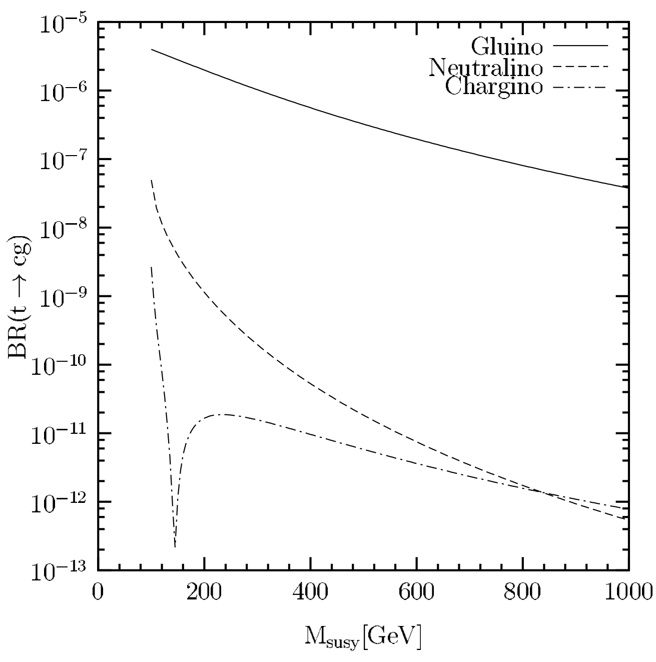

FIG. 2. The gluino, chargino, and neutralino contributions to the BR($t \to cg$) as a function of M_{susy} in the flavor-diagonal scenario for $m_{\tilde{g}} = 300$ GeV, $\tan\beta = 10$ together with the parameter values defined at the beginning of Sec. V. The solid, dashed, and dotted-dashed curves as coded represent gluino, neutralino, and chargino contributions, respectively.

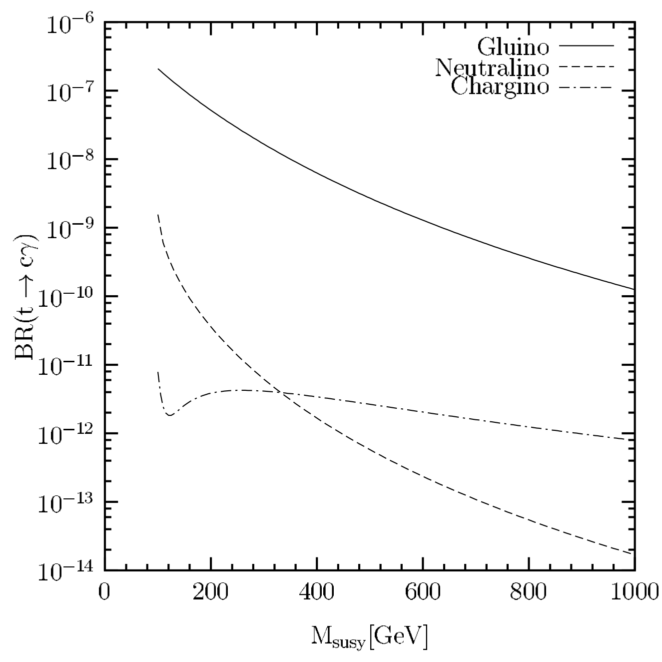

FIG. 3. The same as Fig. 2 but for the $t \to c\gamma$ decay mode.

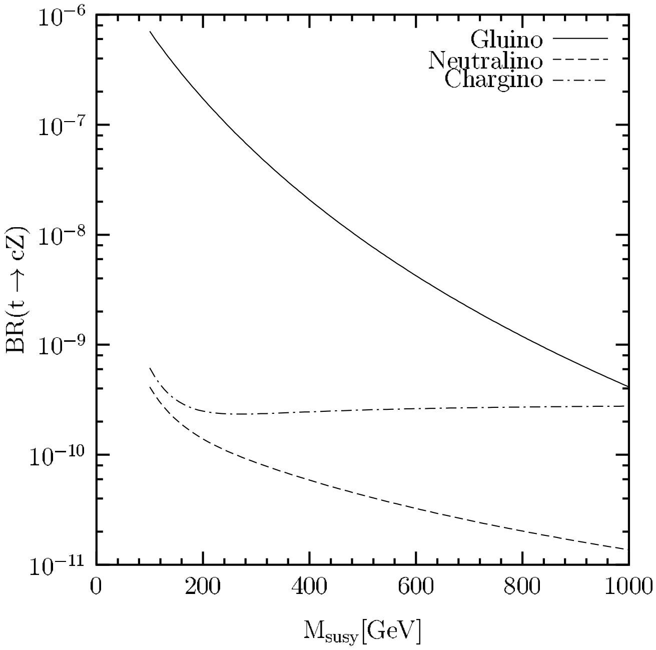

FIG. 4. The same as Fig. 2 but for the $t \to c\gamma$ decay mode.

When we consider the total branching ratios of the three decay modes, in the flavor-diagonal case, the largest one is $t \to cg$ as expected and the smallest is the $t \to c\gamma$ decay mode. For instance, we have $\mathrm{BR}(t \to cg) \sim 10^{-6}$, $\mathrm{BR}(t \to cZ) \sim 5 \times 10^{-8}$, $\mathrm{BR}(t \to c\gamma) \sim 10^{-8}$ for $m_{\tilde{g}} = 300$ GeV and the intermediate $\mathrm{M_{susy}} = 300$ GeV value. It is also possible to get 1 to 2 orders of magnitude larger values for some smaller gluino mass. In all three cases, both the chargino and the neutralino contributions are quite suppressed and negligible for practical purposes.

Before discussing the flavor-nondiagonal effects, we would like to comment on the current experimental limits on these decay modes. The best bound for $t \to c\gamma$ decay mode is from CDF [27], which looked for flavor changing top-quark interactions in $p\bar{p}$ at 1.8 TeV center of mass energy. The bound is $\mathrm{BR}(t \to c\gamma) \leq 3.2 \times 10^{-2}$(95%CL). The bound for the $t \to cZ$ decay channel is weaker, $\mathrm{BR}(t \to cZ) \leq 0.137$(95%CL) for $m_t = 174$ GeV, from the OPAL experiment [28] which searches for single top-quark production in the $e^+e^- \to \bar{t}c$ reaction at around 200 GeV center of mass energy.[5] For the top quark decaying into the charm quark and the gluon, there is no current experimental bound available since background problems make detection of such a channel difficult [30].

These bounds will hopefully be improved by LHC in the near future [31]. Both LHC and the future LC have an advantage over some other colliders for detecting rare top decays because of having better statistics capabilities with lower backgrounds. For example, the foreseen sensitivities to these channels can be as small as [30,31]

$$\begin{aligned} \mathrm{BR}(t \to cg) &\leq 7.4 \times 10^{-3}, \\ \mathrm{BR}(t \to c\gamma) &\leq 1.0 \times 10^{-4}, \\ \mathrm{BR}(t \to cZ) &\leq 1.1 \times 10^{-4}, \end{aligned} \tag{5.1}$$

so that consideration of the flavor-nondiagonal effects are essential, since the enhancement from such effects makes it possible to reach the experimentally accessible range for at least some of the modes considered here. Theoretical estimates for top flavor changing neutral couplings have been performed by using parton-level simulations [32]. Further references and a description of the processes involved can be found in Ref. [33].

B. Flavor-nondiagonal case

The remaining part of our study is devoted to discussing the effects of flavor changing mixings in both the up- and down-type squark mass matrices on the rare $t \to cV$ ($V = g, \gamma, Z$) decays. As motivated in Sec. III, we concentrate only on mixings between the second and the third generations and neglect any kind of mixing involving the first generation. Furthermore, unlike some previous studies where the so-called "mass insertion" method has been used, we follow the general "mass eigenstate" formalism. In addition to the parameters in the flavor-diagonal case, we have eight more, essentially unknown, parameters in each sector as given in Eq. (3.4) (that is, there are totally 16 parameters in both the up and the down sectors).

In order to reduce the set, we should conservatively make some further assumptions and consider some limiting cases. First of all, we set the flavor and chirality diagonal elements of the squark mass matrix to a common SUSY scale, $\mathrm{M_{susy}}$, which enables us to define dimensionless flavor changing parameters, $(\delta_Q^{AB})_{ij}$, $i \neq j = 2, 3$, $A, B = L, R$. In addition to that, we are going to restrict our consideration to a LRSUSY model with Hermitian or symmetric trilinear couplings so that the previously defined mixing parameters satisfy $(\delta_Q^{AB})_{ij} = (\delta_Q^{AB})_{ji}$, $i \neq j = 2, 3, A, B = L, R$. Furthermore, from the LR symmetry we expect $(\delta_Q^{\mathrm{LR}})_{ij} = (\delta_Q^{*\mathrm{RL}})_{ij}$ to hold, though we will not always assume it. Note that this is, in general, not necessarily true (MSSM is an example in this respect). Therefore, we are left with four additional parameters, $(\delta_U^{\mathrm{LL}})_{23}$, $(\delta_U^{\mathrm{RR}})_{23}$, $(\delta_U^{\mathrm{LR}})_{23}$, $(\delta_U^{\mathrm{RL}})_{23}$, in the up-squark sector and similarly for the down-squark sector.

We investigate five limiting cases: the case where the dominant mixing effects come from only the LL, or RR, or LL + LR, or RR + RL, or LR + RL blocks of the squark matrices. However, only the RR, RR + RL, and LR + RL cases are presented here[6] since the ones with LL are very

[5]There is a slightly weaker bound from the ALEPH experiment [29].

[6]We concentrate on mainly the RR sector since the model has an additional $SU(2)_R$ group, which makes the right-handed sector more interesting.

similar but not identical. This is basically because the flavor conserving LR mixing in the second generation is not equal to the one in the third generation. We do not consider limiting cases between δ's and flavor conserving but chirality changing entries of the matrix in Eq. (3.4). One reason is that we kept the trilinear soft terms set to $A = M_{susy}$, which makes such limiting cases very similar to the flavor-diagonal case which we have discussed in the previous section. So we proceed to investigate the five limiting cases in each decay.

1. $t \to cg$

In Fig. 5, we plot in a set of three graphs the gluino, chargino, and neutralino contributions to the BR($t \to cg$) as a function of $(\delta_U^{RR})_{23}$, for $M_{susy} = 300, 400, 1000$ GeV, assuming all other mixings to be zero; $m_{\tilde{g}} = 300$ GeV and $\tan\beta = 10$ are assumed. As seen from the graph on the left, the gluino contribution to the branching ratio depends strongly on $(\delta_U^{RR})_{23}$ and gets enhanced as $(\delta_U^{RR})_{23}$ takes larger values. When we compare with the flavor-diagonal case (Fig. 2), for $M_{susy} = 300$ GeV, there is at least a factor of 5 enhancement for intermediate values of $(\delta_U^{RR})_{23}$ and of 1 order of magnitude for larger values. As M_{susy} gets larger, the enhancement is 2 to 3 orders of magnitude for intermediate-large values of $(\delta_U^{RR})_{23}$. Two comments are in order. The first one is that the comparison with the flavor-diagonal case is not completely right. Here we assume all diagonal elements in the main diagonal identical (to M_{susy}), which leads to larger values for the branching ratio. That is, the enhancement would be slightly different if we relaxed this condition. The second one is that the full (0, 1) interval chosen for the flavor changing parameters is not always available since squark masses become unphysical when δ's exceed certain critical values, which also depend on M_{susy}.

In the same figure, for the same parameter values, we depict the chargino and the neutralino contributions to the BR($t \to cg$) in the second and the third graphs, respectively. The chargino contribution is not very sensitive to δ for $(\delta_D^{RR})_{23} \leq 0.8$ and the enhancement with respect to the flavor-diagonal case is 1 order of magnitude in that range but becomes 3 orders of magnitude after. When we compare it with the gluino contribution, it remains suppressed even under extreme conditions [i.e., $(\delta_U^{RR})_{23} = 0$, $(\delta_D^{RR})_{23} \sim 1$]. So the flavor mixing effects from the down-squark sector do not contribute to the total BR of the $t \to cg$ decay nearly as much as the ones from the up-squark sector. For the neutralino case, even though there is 3 to 4 orders of magnitude enhancement for intermediate or large $(\delta_D^{RR})_{23}$ values, it is at least 1 order of magnitude smaller than the gluino contribution. Neutralino contribution still dominates the chargino one unless $(\delta_U^{RR})_{23} = 0$, $(\delta_D^{RR})_{23} \sim 1$ occurs.

In Fig. 6, we show the same as Fig. 5 but this time additionally the mixing in the LR sector is turned on. For simplicity, we present the case $(\delta_{U(D)}^{RR})_{23} = (\delta_{U(D)}^{RL})_{23}$. With respect to the case where only the RR mixing is turned on, except the chargino contribution, the gluino and the neutralino contributions are suppressed about 1 order of magnitude. However, in the down sector, the chargino contribution depends strongly on $(\delta_D^{RR})_{23}$ and reaches 10^{-6} at maximum level. In the rest of the interval, it is still very suppressed in the total BR($t \to cg$).

The last case for the BR($t \to cg$), as a function of $(\delta_{U(D)}^{LR})_{23} = (\delta_{U(D)}^{RL})_{23}$, is depicted in Fig. 7. We have smaller contributions with respect to both the RR and RR + RL cases except for the neutralino contribution,

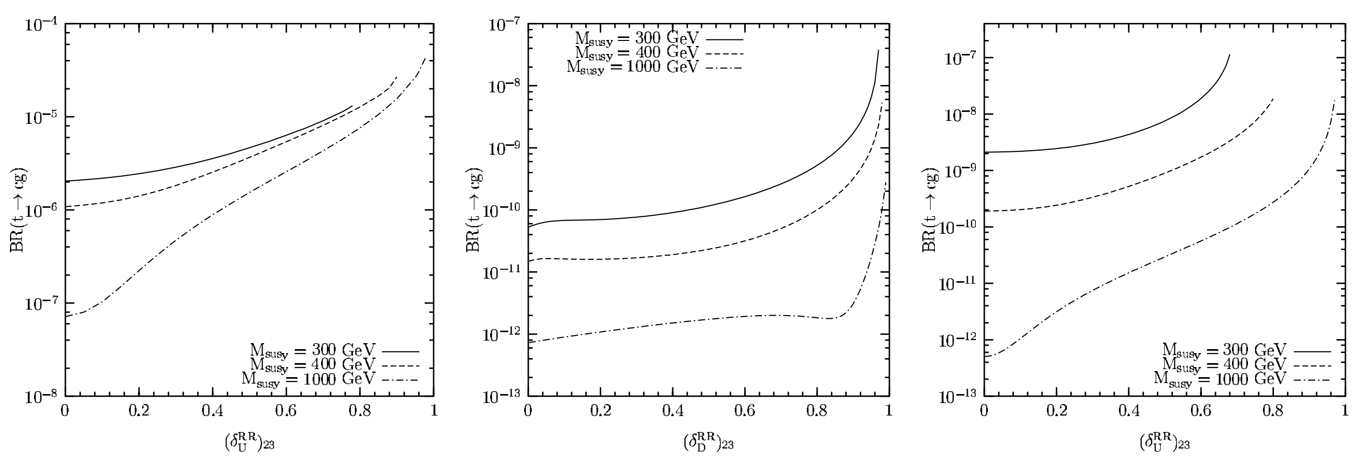

FIG. 5. The gluino, chargino, and neutralino contributions to the BR($t \to cg$) as a function of $(\delta_{U(D)}^{RR})_{23}$ for the case where the flavor changing effects come from the mixing in the RR sector only. The graph on the left (right) denotes the gluino (neutralino) contribution for the up RR mixing. The one in the middle is for the chargino as a function of $(\delta_D^{RR})_{23}$. Three representative M_{susy} values (300, 400, 1000) GeV are chosen with $m_{\tilde{g}} = 300$ GeV, $\tan\beta = 10$, $A = M_{susy}$, and $\mu = 200$ GeV.

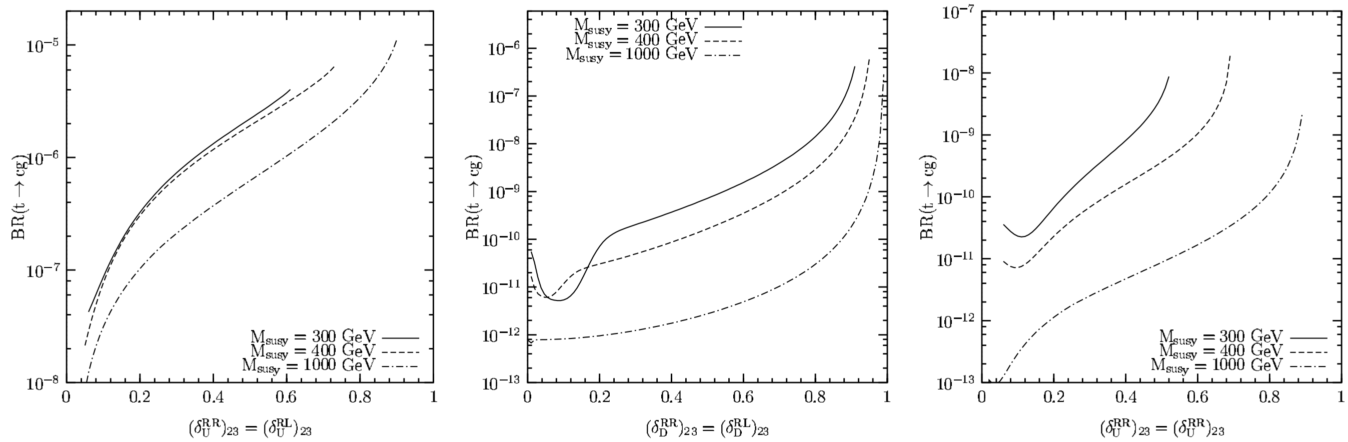

FIG. 6. The same as Fig. 5 but for the case where both $(\delta^{RR}_{U(D)})_{23}$ and $(\delta^{RL}_{U(D)})_{23}$ contribute, with the assumption $(\delta^{RR}_{U(D)})_{23} = (\delta^{RL}_{U(D)})_{23}$.

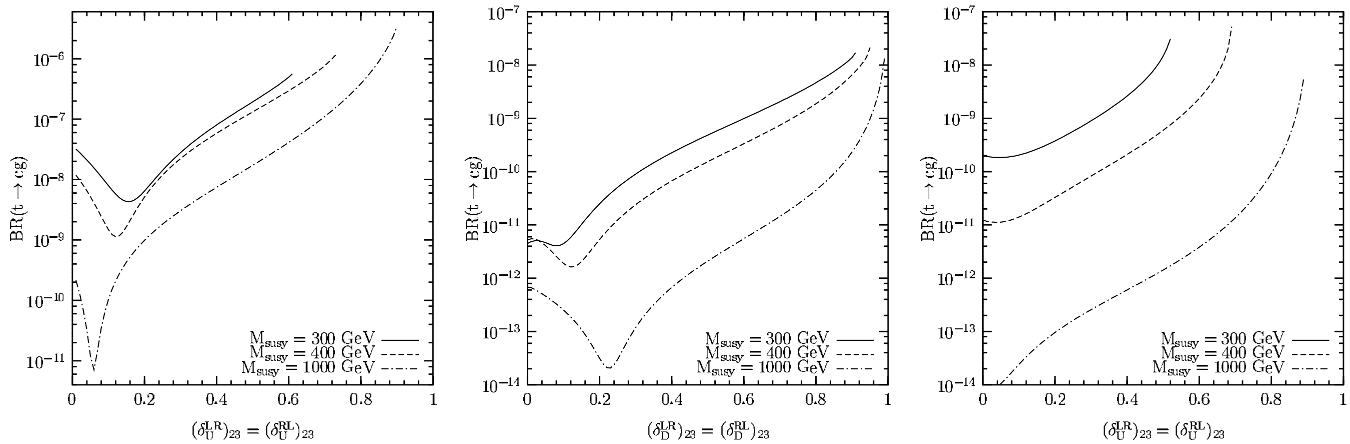

FIG. 7. The same as Fig. 5 but for the case where both $(\delta^{LR}_{U(D)})_{23}$ and $(\delta^{RL}_{U(D)})_{23}$ contribute, with the assumption $(\delta^{LR}_{U(D)})_{23} = (\delta^{RL}_{U(D)})_{23}$.

which reaches values obtained in the RR + RL case for upper values of $(\delta_U^{LR})_{23}$ at different M_{susy} values. So, overall, the gluino gives the largest contribution among three in all limiting cases considered above, and the total BR can get up to a few times 10^{-5} for $m_{\tilde{g}} = 300$ GeV.

2. $t \to c\gamma$

We analyze the $t \to c\gamma$ decay in the same order as we did the $t \to cg$ decay and we present only the RR and RR + RL cases. In Fig. 8, we show the dependence of the gluino, chargino, and neutralino contributions in the BR($t \to c\gamma$) to the parameter $(\delta^{RR}_{U(D)})_{23}$ at different M_{susy} values, 300, 400, 1000 GeV for $\tan\beta = 10$. The gluino contribution is now 1 order of magnitude bigger than in the flavor-diagonal case for $M_{susy} = 300$ GeV and $(\delta_U^{RR})_{23} \sim 0.5$, and it can go up to 3 orders larger for $M_{susy} = 1000$ GeV with maximum possible mixing in the up RR sector, or for smaller M_{susy} with slightly smaller mixings, when it reaches 10^{-6}. In the chargino case, in the down sector, the dependence of the BR to M_{susy} is very weak and enhancement with respect to the flavor-diagonal case varies from 2 to 3 times to about 2 orders of magnitude. Here the curve for $M_{susy} = 300$ GeV has a minimum which is peculiar only to this M_{susy} value and occurs due to some precise cancellations in the loop functions. For example, we do not have such behavior in the down LL mixing case which we have not discussed here explicitly. It is, however, possible to get bigger enhancements for smaller M_{susy} values. For the neutralino contribution, bigger enhancements ranging between 2 to 4 orders of magnitude occur, depending on the value of M_{susy} and the amount of mixing allowed in the up RR sector. Unlike the flavor-diagonal case, the neutralino contribution always dominates the chargino one through the entire interval scanned here, while it is still suppressed with respect to the gluino contribution.

The RR + RL case $[(\delta^{RR}_{U(D)})_{23} = (\delta^{RL}_{U(D)})_{23}]$ for the BR($t \to c\gamma$) is similar and shown in Fig. 9. The gluino contribution remains always 1 order smaller than in the RR mixing case. This is true for the neutralino as well, even

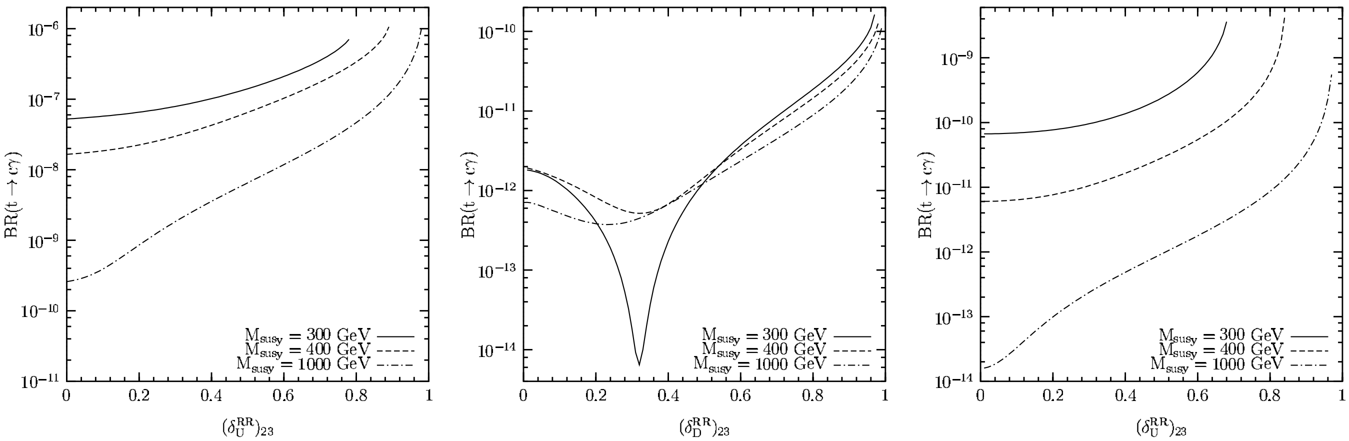

FIG. 8. The gluino, chargino, and neutralino contributions to the BR($t \to c\gamma$) as a function of $(\delta^{RR}_{U(D)})_{23}$ for the case where the flavor changing effects come from the mixing in the RR sector only. The graph on the left (right) denotes the gluino (neutralino) contribution for the up RR mixing. The one in the middle is for the chargino as a function of $(\delta^{RR}_{D})_{23}$. Three representative M$_{\rm susy}$ values (300, 400, 1000) GeV are chosen, with $m_{\tilde{g}} = 300$ GeV, $\tan\beta = 10$, $A = \mathrm{M_{susy}}$, and $\mu = 200$ GeV.

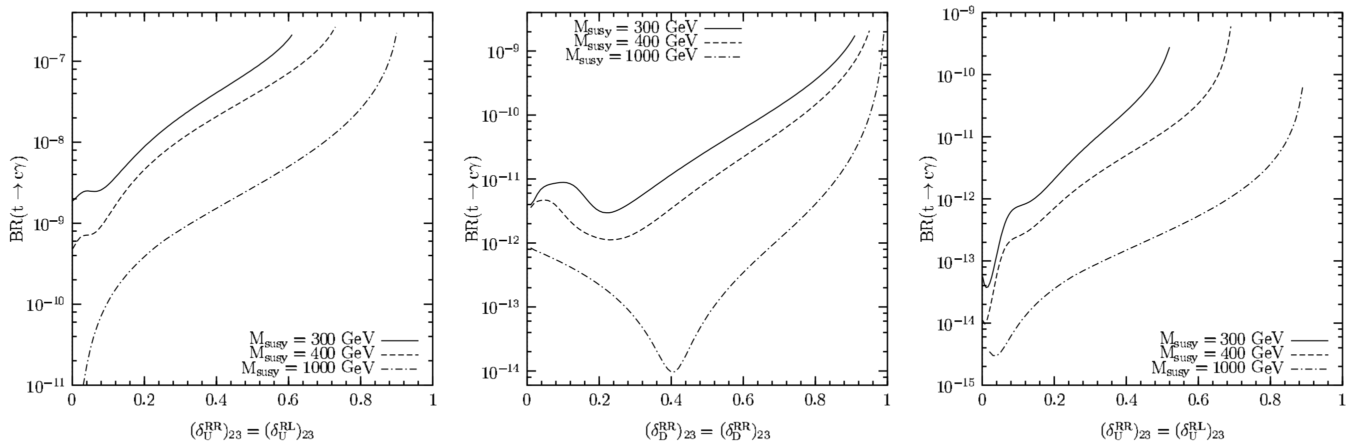

FIG. 9. The same as Fig. 8 but for the case where both $(\delta^{RR}_{U(D)})_{23}$ and $(\delta^{RL}_{U(D)})_{23}$ contribute, with the assumption $(\delta^{RR}_{U(D)})_{23} = (\delta^{RL}_{U(D)})_{23}$.

though there is a sharp dependency on the mixing parameters in the up sector. In this case, like in the flavor-diagonal case, the chargino curves can still cross the neutralino curves and become larger with larger down-type RR + RL mixing. If one considers the mixing in the up and the one in the down sector to be completely independent, one could end up with different conclusions.

3. $t \to cZ$

Finally, we present the final decay channel, $t \to cZ$. Figure 10 shows the branching ratio of $t \to cZ$ as a function of $(\delta^{RR}_{U(D)})_{23}$ when the rest of the flavor changing parameters are set to be zero. The gluino contribution becomes enhanced 1 to more than 4 orders of magnitude compared to the case where all flavor changing parameters are turned off. A 10^{-5} branching ratio seems to be reachable. We can discuss the neutralino contributions and directly compare with the gluino contributions since, unlike the chargino contribution, they vary with respect to the same mixing parameter. The neutralino contribution is more than 2 orders of magnitude enhanced over the flavor-diagonal case but still suppressed with respect to the gluino contribution. The effect in the chargino sector is tiny.

The RR + RL case, presented in Fig. 11, is very similar to the above case for both the gluino and the neutralino contributions but the mixing in the RR + RL down sector contributes to the chargino loop more effectively and increases its contribution to 1 order of magnitude with respect to the RR mixing case considered above. The chargino contribution, like in the flavor-diagonal case, is still dominant over the neutralino one.

Consideration of flavor changing effects in either sector is quite important for rare top decays $t \to cV$ ($V = g, \gamma, Z$)

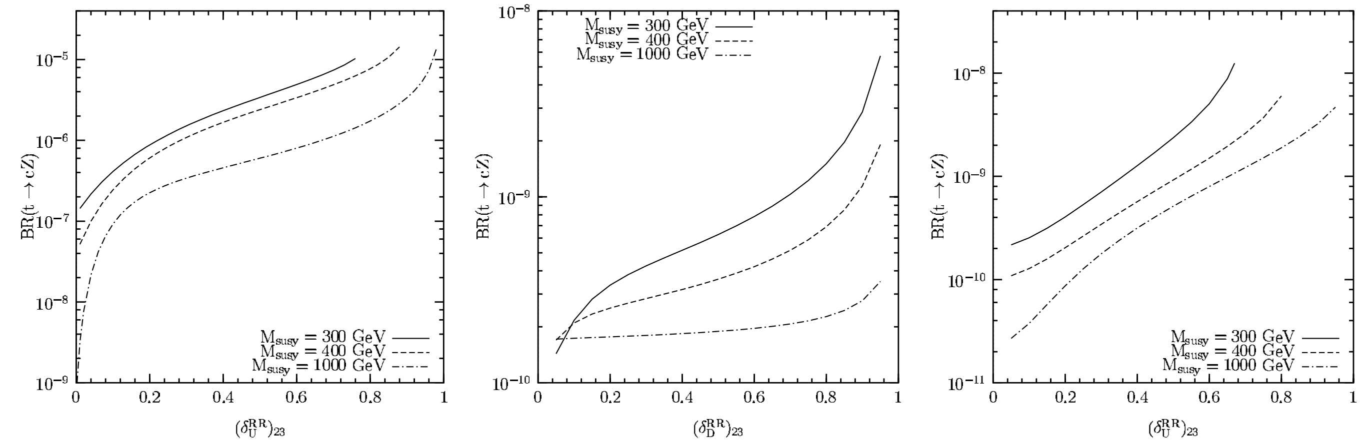

FIG. 10. The gluino, chargino, and neutralino contributions to the BR($t \to cZ$) as a function of $(\delta^{RR}_{U(D)})_{23}$ for the case where the flavor changing effects come from the mixing in the RR sector only. The graph on the left (right) denotes the gluino (neutralino) contribution for the up RR mixing. The one in the middle is for the chargino as a function of $(\delta^{RR}_{D})_{23}$. Three representative M$_{susy}$ values (300, 400, 1000) GeV are chosen with $m_{\tilde{g}} = 300$ GeV, $\tan\beta = 10$, $A = \mathrm{M}_{susy}$, and $\mu = 200$ GeV.

not only for the expected enhancements, but also for revealing the details of the SUSY flavor changing mechanism.

Before closing this section, we would like to comment on the large $\tan\beta$ effects in both flavor-diagonal and -nondiagonal scenarios. As can be seen from the Feynman rules for the gluino, the gluino loop for each decay is not sensitive to the $\tan\beta$ value. Basically, the branching ratios are slightly bigger for $\tan\beta$ values smaller than 10 and almost constant after 10. This indeed makes the total branching ratios of the decays insensitive to $\tan\beta$. Since we discuss contributions from each loop separately, it is worth mentioning the $\tan\beta$ dependencies of chargino and neutralino loops. We comment on specifically the $t \to c\gamma$ decay, but these comments apply to other decay modes as well. In the flavor-diagonal scenario, unlike the neutralino case, $\tan\beta$ dependency of the chargino contribution to $t \to c\gamma$ is quite strong. If one scans the BR(t → cγ) in $\tan\beta$ in the 0–50 interval, the chargino contribution can reach 2×10^{-10} (3×10^{-11}) for M$_{susy}$ = 300, 400 (1000) GeV, which is about 40 (30) times bigger than $\tan\beta = 10$ case and 3 orders of magnitude bigger with respect to very small $\tan\beta$ values. For the neutralino case, the branching ratio of $t \to c\gamma$ is almost constant as $\tan\beta$ changes for M$_{susy}$ = 300, 400 GeV. As M$_{susy}$ gets bigger, there is a slight increase in the branching ratio as $\tan\beta$ increases. So for M$_{susy}$ = 1000 GeV, we have 1 order of magnitude enhancement at maximum value of the branching ratio in the interval $\tan\beta \in (0, 50)$ and 5 times bigger branching ratios when we compare the case for $\tan\beta = 50$ with the one for $\tan\beta = 10$. Even though the chargino contribution becomes larger for large $\tan\beta$, it is still quite suppressed

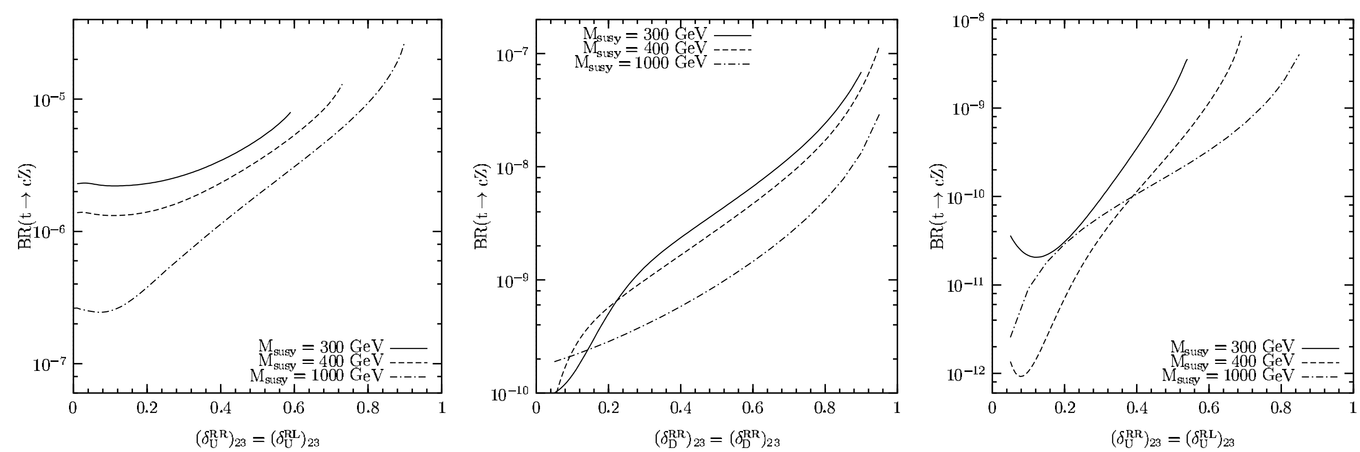

FIG. 11. The same as Fig. 10 but for the case where both $(\delta^{RR}_{U(D)})_{23}$ and $(\delta^{RL}_{U(D)})_{23}$ contribute, with the assumption $(\delta^{RR}_{U(D)})_{23} = (\delta^{RL}_{U(D)})_{23}$.

with respect to the gluino contribution. For example, for $\tan\beta = 50$, the chargino contribution is around 75; 40; 4 times smaller for $M_{susy} = 300; 400; 1000$ GeV. Where the chargino contribution was 1 to 2 orders of magnitude smaller than the neutralino contribution for small $\tan\beta$ values ($\tan\beta \leq 5$), it becomes more than 1 (2) order(s) bigger then the neutralino contribution for $M_{susy} = 300, 400$ (1000) GeV as $\tan\beta$ takes its maximum value considered here.

For the flavor-nondiagonal case, we consider again as an example the $t \to c\gamma$ decay with mixing only in the RR sector. The gluino contribution is very similar to the flavor-diagonal case. The chargino contribution is again most sensitive to $\tan\beta$ among the three contributions. The sensitivity in this case is slightly less, because we have assumed for simplicity the diagonal entries of the squark mass matrices identical to M_{susy}. For the flavor mixing parameter $(\delta_D^{RR})_{23} = 0.5$, the chargino contribution to the branching ratio becomes 35; 25; 10 times bigger for $M_{susy} = 300; 400; 1000$ GeV when we compare the cases $\tan\beta = 10$ and $\tan\beta = 50$. For the neutralino case, the discussion is very similar to the flavor-diagonal case with a reduced dependency on $\tan\beta$. In summary, when we add up all contributions the branching ratio is not sensitive to the $\tan\beta$ parameter because the gluino loop which dominates all contributions and it is not sensitive to $\tan\beta$. When we look at the individual contributions to the branching ratio, the chargino contribution has the strongest dependency on $\tan\beta$ in both flavor-diagonal and -nondiagonal scenarios. Thus, our overall analysis is valid as one varies $\tan\beta$ in the phenomenologically relevant interval.

VI. CONCLUSION

We have presented a complete analysis of the two-body flavor violating decays of the top quark in a fully left-right supersymmetric model. The model has a left-right scalar quark sector, as well as a right-handed gaugino in both neutral and charged sectors.

We have first evaluated the branching ratios in the flavor-diagonal case, which corresponds to the constrained LRSUSY, where explicit flavor violating terms do not exist in the squark mass matrix, and the only source of flavor violation is the CKM matrix. The enhancement due to the right-handed sector proves to be minimal, and top-quark flavor-changing neutral-current (FCNC) decay is somewhat below the detectable level at future collider experiments. We present the results for the gluino, chargino, and neutralino contributions to the decays $t \to cg$, $c\gamma$, and cZ explicitly and separately, unlike previous analyses. We confirm that branching ratios can reach 10^{-5} for the gluon and 10^{-6} for the photon and Z, in agreement with previous results.

We then analyzed the branching ratios for $t \to cV$ in the unconstrained LRSUSY, where flavor changing elements in the squark mass matrix are allowed to be arbitrarily large, though for the mixing between the second and the third generation only. This mixing is rather unconstrained in the top-charm squark sector and only very weakly constrained in the bottom-strange squark sector from b decays. We analyze again separately contributions from gluino, chargino, and neutralino to $t \to cg$, $c\gamma$, and cZ and look at results for the branching ratios when flavor violation is driven by the right-right, right-left, and left-right, or both, squark mixings. As expected, enhancements of the branching ratio could occur here, and we obtain, under the best circumstances, $\mathrm{BR}(t \to cg)$ close to 10^{-4}, while $\mathrm{BR}(t \to c\gamma)$ and $\mathrm{BR}(t \to cZ)$ can reach 10^{-6} and 10^{-5}, respectively. These enhancements are comparable to other studies in the unconstrained MSSM, so while t decays are important as a probe of SUSY FCNC couplings, they do not appear to differentiate between the MSSM and LRSUSY, unless there are clear indications that it is the right-handed scalar quark sector which is responsible for flavor violation. In this case, LRSUSY provides large enhancements, at least competitive with ones driven by the left-handed scalar quark mixings in MSSM.

Comparing t to b decays, it appears that there the chargino contribution is more evident and sometimes comparable to the gluino; while this is never the case for $t \to cV$, and, thus, here the effect of the right-handed gaugino is relatively obscured.

ACKNOWLEDGMENTS

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APPENDIX A: CHARGINO AND NEUTRALINO MIXINGS

In this appendix, we present the chargino and neutralino mass matrices and define their eigenvalues and eigenvectors.

1. Chargino mixing

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

$$\mathcal{L}_C = -\frac{1}{2}(\psi^{+T}, \psi^{-T})\begin{pmatrix} 0 & X^T \\ X & 0 \end{pmatrix}\begin{pmatrix} \psi^+ \\ \psi^- \end{pmatrix} + \mathrm{H.c.}, \tag{A1}$$

where $\psi^{+T} = (-i\lambda_L^+, -i\lambda_R^+, \tilde{\phi}_{1u}^+, \tilde{\phi}_{1d}^+, \tilde{\Delta}_R^+)$ and $\psi^{-T} = (-i\lambda_L^-, -i\lambda_R^-, \tilde{\phi}_{2u}^-, \tilde{\phi}_{2d}^-, \tilde{\delta}_R^-)$, and

$$X = \begin{pmatrix} M_L & 0 & g_L\kappa_u & 0 & 0 \\ 0 & M_R & g_R\kappa_u & 0 & \sqrt{2}g_R v_{\delta_R} \\ 0 & 0 & 0 & -\mu & 0 \\ g_L\kappa_d & g_R\kappa_d & -\mu & 0 & 0 \\ 0 & \sqrt{2}g_R v_{\Delta_R} & 0 & 0 & -\mu \end{pmatrix}, \tag{A2}$$

where we have taken, for simplification, $\mu_{ij} = \mu$. Note that although there are six charginos, $\tilde{\Delta}_L^-$ and $\tilde{\delta}_L^+$ decouple from the spectrum because we have chosen the VEV's of Δ_L and δ_L to be zero.

The chargino mass eigenstates χ_i are obtained by

$$\chi_i^+ = V_{ij}\psi_j^+, \qquad \chi_i^- = U_{ij}\psi_j^-, \qquad i, j = 1, \ldots 5, \tag{A3}$$

with V and U unitary matrices satisfying

$$U^* X V^{-1} = X_D. \tag{A4}$$

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

2. Neutralino mixing

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

$$\mathcal{L}_N = -\frac{1}{2}\psi^{0T} Z \psi^0 + \text{H.c.}, \tag{A5}$$

where $\psi^{0T} = (-i\lambda_L^0, -i\lambda_R^0, -i\lambda_V, \tilde{\phi}_{1u}^0, \tilde{\phi}_{2d}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0, \tilde{\phi}_{1d}^0, \tilde{\phi}_{2u}^0)$, and

$$Z = \begin{pmatrix} M_L & 0 & 0 & \frac{g_L\kappa_u}{\sqrt{2}} & -\frac{g_L\kappa_d}{\sqrt{2}} & 0 & 0 & 0 & 0 \\ 0 & M_R & 0 & \frac{g_R\kappa_u}{\sqrt{2}} & -\frac{g_R\kappa_d}{\sqrt{2}} & -\sqrt{2}g_R v_{\Delta_R} & -\sqrt{2}g_R v_{\delta_R} & 0 & 0 \\ 0 & 0 & M_V & 0 & 0 & 2\sqrt{2}g_V v_{\Delta_R} & 2\sqrt{2}g_V v_{\delta_R} & 0 & 0 \\ \frac{g_L\kappa_u}{\sqrt{2}} & \frac{g_R\kappa_u}{\sqrt{2}} & 0 & 0 & -\mu & 0 & 0 & 0 & 0 \\ -\frac{g_L\kappa_d}{\sqrt{2}} & -\frac{g_R\kappa_d}{\sqrt{2}} & 0 & -\mu & 0 & 0 & 0 & 0 & 0 \\ 0 & -\sqrt{2}g_R v_{\Delta_R} & 2\sqrt{2}g_V v_{\Delta_R} & 0 & 0 & 0 & -\mu & 0 & 0 \\ 0 & -\sqrt{2}g_R v_{\delta_R} & 2\sqrt{2}g_V v_{\delta_R} & 0 & 0 & -\mu & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu & 0 \end{pmatrix}. \tag{A6}$$

As in the case of charginos, $\tilde{\Delta}_L^0$ and $\tilde{\delta}_L^0$ decouple from the spectrum because we have chosen the VEV's of Δ_L and δ_L to be zero. The mass eigenstates are defined by

$$\chi_i^0 = N_{ij}\psi_j^0 \qquad (i, j = 1, 2, \ldots 9), \tag{A7}$$

where N is a unitary matrix chosen such that

$$NZN^T = Z_D, \tag{A8}$$

and Z_D is a diagonal matrix with non-negative entries.[7]

One can further switch to a basis involving the photino, left and right Z-ino states

$$\psi^{0\prime T} = (-i\lambda_{\tilde{\gamma}}, -i\lambda_{\tilde{Z}_L}, -i\lambda_{\tilde{Z}_R}, \tilde{\phi}_{1u}^0, \tilde{\phi}_{2d}^0, \tilde{\Delta}_R^0, \tilde{\delta}_R^0, \tilde{\phi}_{1d}^0, \tilde{\phi}_{2u}^0),$$

where the photino and left and right Z-ino states are defined as

$$\begin{pmatrix} \lambda_{\tilde{\gamma}} \\ \lambda_{\tilde{Z}_L} \\ \lambda_{\tilde{Z}_R} \end{pmatrix} = \begin{pmatrix} \sin\theta_{\rm W} & \sin\theta_{\rm W} & \sqrt{\cos2\theta_{\rm W}} \\ \cos\theta_{\rm W} & -\sin\theta_{\rm W}\tan\theta_{\rm W} & -\sqrt{\cos2\theta_{\rm W}}\tan\theta_{\rm W} \\ 0 & \sqrt{\cos2\theta_{\rm W}}\sec\theta_{\rm W} & -\tan\theta_{\rm W} \end{pmatrix} \begin{pmatrix} \lambda_L^0 \\ \lambda_R^0 \\ \lambda_V \end{pmatrix}. \tag{A9}$$

The mass matrix in the above basis becomes

[7]The positivity of the entries of Z_D can be achieved by multiplying some of the rows of N with i. This freedom comes from the fact that in order to determine N the square of Eq. (A8) needs to be considered. We have a similar situation for the chargino sector. For example, see [34] for details in the MSSM.

$$Z = \begin{pmatrix} m_{\tilde{\gamma}} & m_{\tilde{\gamma}\tilde{Z}_L} & m_{\tilde{\gamma}\tilde{Z}_R} & \sqrt{2}e\kappa_u & -\sqrt{2}e\kappa_d & \sqrt{2}ev_{\Delta_R} & \sqrt{2}ev_{\delta_R} & 0 & 0 \\ m_{\tilde{\gamma}\tilde{Z}_L} & m_{\tilde{Z}_L} & m_{\tilde{Z}_L\tilde{Z}_R} & A_u & -A_d & C_{\Delta_R} & C_{\delta_R} & 0 & 0 \\ m_{\tilde{\gamma}\tilde{Z}_R} & m_{\tilde{Z}_L\tilde{Z}_R} & m_{\tilde{Z}_R} & E_u & -E_d & B_{\Delta_R} & B_{\delta_R} & 0 & 0 \\ \sqrt{2}e\kappa_u & A_u & E_u & 0 & -\mu & 0 & 0 & 0 & 0 \\ -\sqrt{2}e\kappa_d & -A_d & -E_d & -\mu & 0 & 0 & 0 & 0 & 0 \\ \sqrt{2}ev_{\Delta_R} & C_{\Delta_R} & B_{\Delta_R} & 0 & 0 & 0 & -\mu & 0 & 0 \\ \sqrt{2}ev_{\delta_R} & C_{\delta_R} & B_{\delta_R} & 0 & 0 & -\mu & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & -\mu & 0 \end{pmatrix}, \tag{A10}$$

with the functions

$$\begin{gathered} m_{\tilde{\gamma}} = (M_L + M_R)\sin^2\theta_W + M_V\cos 2\theta_W, \qquad m_{\tilde{\gamma}Z_L} = M_L\cos\theta_W\sin\theta_W - (M_R\sin^2\theta_W + M_V\cos 2\theta_W)\tan\theta_W, \\ m_{\tilde{\gamma}Z_R} = (M_R - M_V)\sqrt{\cos 2\theta_W}\tan\theta_W, \qquad m_{\tilde{Z}_L} = M_L\cos^2\theta_W + M_R\sin^2\theta_W\tan^2\theta_W + M_V\cos 2\theta_W\tan^2\theta_W, \\ m_{\tilde{Z}_L\tilde{Z}_R} = -(M_R - M_V)\sqrt{\cos 2\theta_W}\tan^2\theta_W, \qquad m_{\tilde{Z}_R} = M_R(1 - \tan^2\theta_W) + M_V\tan^2\theta_W, \qquad A_d = \sqrt{2}e\kappa_d\cot 2\theta_W, \\ A_u = \sqrt{2}e\kappa_u\cot 2\theta_W, \qquad B_{\Delta_R} = \frac{-\sqrt{2}g_L v_{\Delta_R}}{\cos\theta_W\sqrt{\cos 2\theta_W}}, \qquad B_{\delta_R} = \frac{-\sqrt{2}g_L v_{\delta_R}}{\cos\theta_W\sqrt{\cos 2\theta_W}}, \qquad C_{\Delta_R} = -\sqrt{2}ev_{\Delta_R}\tan\theta_W, \\ C_{\delta_R} = -\sqrt{2}ev_{\delta_R}\tan\theta_W, \qquad E_d = \frac{g_L\sqrt{\cos 2\theta_W}}{\sqrt{2}\cos\theta_W}\kappa_d, \qquad E_u = \frac{g_L\sqrt{\cos 2\theta_W}}{\sqrt{2}\cos\theta_W}\kappa_u, \end{gathered} \tag{A11}$$

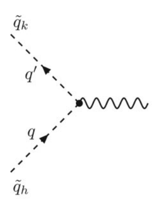
where we have used $g_L = g_R$ and $g_V = e/\sqrt{\cos 2\theta_W}$. Furthermore, the unitary matrix N can be expressed in the rotated basis as

$$\begin{pmatrix} N'_{j1} \\ N'_{j2} \\ N'_{j3} \end{pmatrix} = \begin{pmatrix} \sin\theta_W & \sin\theta_W & \sqrt{\cos 2\theta_W} \\ \cos\theta_W & -\sin\theta_W\tan\theta_W & -\sqrt{\cos 2\theta_W}\tan\theta_W \\ 0 & \sqrt{\cos 2\theta_W}\sec\theta_W & -\tan\theta_W \end{pmatrix} \begin{pmatrix} N_{j1} \\ N_{j2} \\ N_{j3} \end{pmatrix}, \qquad N'_{jk} = N_{jk}, \qquad j = 1, \ldots, 9, \qquad k = 4, \ldots, 9. \tag{A12}$$

APPENDIX B: FEYNMAN RULES

In this appendix, we give the Feynman rules involving SUSY particles. The ones for the SM particles can be found in many textbooks.

Gauge boson-squark-squark interaction

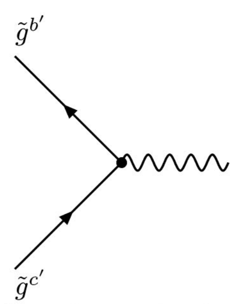


$g_\mu^{a'}$: $-ig_s T^{a'} (q + q')_\mu \delta_{hk}$

γ_μ : $-ieQ_q (q + q')_\mu \delta_{hk}$

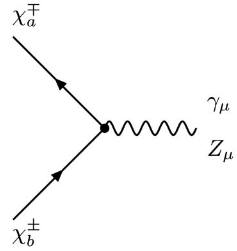
Z_μ : $-\dfrac{ig(q + q')_\mu}{\cos\theta_W}\left[T_{3Q}\sum_{i=1}^{3}\Gamma_{QL}^{*hi}\Gamma_{QL}^{ki} - Q_q\sin^2\theta_W\delta_{hk}\right]$

Gluon-gluino-gluino interaction



$g_\mu^{a'}$: $-ig_s f_{a'b'c'}\gamma_\mu$

Gauge boson-chargino-chargino interaction

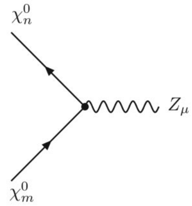


$: -ie\gamma_\mu$

$: \frac{ig}{\cos\theta_W}\gamma_\mu \left[O_L^{ab}P_L + O_R^{ab}P_R\right]$

$$O_L^{ab} = -V_{a1}V_{b1}^* - \tfrac{1}{2}V_{a3}V_{b3}^* - \tfrac{1}{2}V_{a4}V_{b4}^* + \delta_{ab}\sin^2\theta_W, \qquad O_R^{ab} = -U_{a1}^*U_{b1} - \tfrac{1}{2}U_{a3}^*U_{b3} - \tfrac{1}{2}U_{a4}^*U_{b4} + \delta_{ab}\sin^2\theta_W.$$

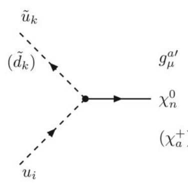
Z boson-neutralino-neutralino interaction



$: \frac{ig}{\cos\theta_W}\gamma_\mu \left[O_L^{\prime nm}P_L + O_R^{\prime nm}P_R\right]$

$$O_L^{\prime nm} = \tfrac{1}{2}\left(N'_{n4}N'^*_{m4} - N'_{n5}N'^*_{m5} + N'_{n6}N'^*_{m6} - N'_{n7}N'^*_{m7}\right)$$
$$O_R^{\prime nm} = -O_L^{\prime * nm}$$

$\tilde{g}$, χ^0, χ^+-quark-squark interaction



$: -i\sqrt{2}g_sT^{a'}\left(\Gamma_{UL}^{ki}P_L - \Gamma_{UR}^{ki}P_R\right)$

$: -ig\left[\left(\sqrt{2}G_{UL}^{[n,k,i]} + H_{UR}^{[n,k,i]}\right)P_L - \left(\sqrt{2}G_{UR}^{[n,k,i]} - H_{UL}^{[n,k,i]}\right)P_R\right]$

$: -ig\left[\left(G_{DL}^{[a,k,i]} - H_{DR}^{[a,k,i]}\right)P_L + \left(G_{DR}^{[a,k,i]} - H_{DL}^{[a,k,i]}\right)P_R\right]$

$$G_{UL}^{[n,k,i]} = \left[Q_u\sin\theta_W N'^*_{n1} + \frac{1}{\cos\theta_W}(T_{3u} - Q_u\sin^2\theta_W)N'^*_{n2} - \frac{Q_u+Q_d}{2}\frac{\sin\theta_W\tan\theta_W}{\sqrt{\cos 2\theta_W}}N'^*_{n3}\right]\Gamma_{UL}^{ki},$$

$$G_{UR}^{[n,k,i]} = \left[Q_u\sin\theta_W N'^*_{n1} - \frac{Q_u\sin^2\theta_W}{\cos\theta_W}N'^*_{n2} + \frac{\sqrt{\cos 2\theta_W}}{\cos\theta_W}\left(T_{3u} - \frac{Q_u+Q_d}{2}\frac{\sin^2\theta_W}{\sqrt{\cos 2\theta_W}}\right)N'^*_{n3}\right]\Gamma_{UR}^{ki},$$

$$H_{UL}^{[n,k,i]} = \frac{1}{\sqrt{2}m_W}\left[\frac{m_{u_i}}{\sin\beta}N'_{n5} + \frac{m_{d_i}}{\cos\beta}N'_{n7}\right]\Gamma_{UL}^{ki}, \qquad H_{UR}^{[n,k,i]} = \frac{1}{\sqrt{2}m_W}\left[\frac{m_{u_i}}{\sin\beta}N'^*_{n5} + \frac{m_{d_i}}{\cos\beta}N'^*_{n7}\right]\Gamma_{UR}^{ki},$$

$$G_{DL}^{[a,k,i]} = \sum_{j=1}^{3}K_{\text{CKM}}^{*ij}\Gamma_{DL}^{kj}U_{a1}^*, \qquad G_{DR}^{[a,k,i]} = \sum_{j=1}^{3}K_{\text{CKM}}^{*ij}\Gamma_{DR}^{kj}V_{a2}, \qquad H_{DL}^{[a,k,i]} = \frac{1}{\sqrt{2}m_W}\sum_{j=1}^{3}\left[\frac{m_{u_i}}{\sin\beta}V_{a4} + \frac{m_{d_j}}{\cos\beta}V_{a3}\right]K_{\text{CKM}}^{*ij}\Gamma_{DL}^{kj},$$

$$H_{DR}^{[a,k,i]} = \frac{1}{\sqrt{2}m_W}\sum_{j=1}^{3}\left[\frac{m_{u_i}}{\sin\beta}U_{a4}^* + \frac{m_{d_j}}{\cos\beta}U_{a3}^*\right]K_{\text{CKM}}^{*ij}\Gamma_{DR}^{kj}.$$

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