Impact of squark generation mixing on the search for gluinos at LHC

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Abstract

We study gluino decays in the Minimal Supersymmetric Standard Model (MSSM) with squark generation mixing. We show that the effect of this mixing on the gluino decay branching ratios can be very large in a significant part of the MSSM parameter space despite the very strong experimental constraints on quark flavour violation (QFV) from B meson observables. Especially we find that under favourable conditions the branching ratio of the the QFV gluino decay $\tilde{g} \rightarrow c \bar{t} \chi_1^0$ can be as large as $\sim 50\%$. We also find that the squark generation mixing can result in a multiple-edge (3- or 4-edge) structure in the charm-top quark invariant mass distribution. The appearance of this remarkable structure provides an additional powerful test of supersymmetric QFV at LHC. These could have an important impact on the search for gluinos and the determination of the MSSM parameters at LHC.
1 Introduction

The search for supersymmetric (SUSY) particles will have a high priority at the Large Hadron Collider (LHC) at CERN. If weak scale SUSY is realized in nature, gluinos and squarks, the SUSY partners of gluons and quarks, will have high production rates for masses up to $O(1 \text{ TeV})$. The main decay modes of gluinos and squarks are usually assumed to be quark-flavour conserving (QFC). However, the squarks are not necessarily quark-flavour eigenstates and they are in general mixed by a $6 \times 6$ matrix. In this case quark-flavour violating (QFV) decays of gluinos and squarks could occur.

The effect of QFV in the squark sector on reactions at colliders has been studied only in a few publications. The pair production of quarks with different flavours at the LHC is studied in [1]. The QFV effect can also be probed in the top quark decay [2]. Moreover, QFV Higgs decays can have rates accessible at future colliders, see e.g. [3]. In all of these studies the external particles of the reactions are Standard Model (SM) particles (or SUSY Higgs bosons). This means that the effect of QFV in the squark sector is induced only by SUSY particle (sparticle) loops.

In sparticle reactions, on the other hand, the effect of QFV in the squark sector may be especially strong as they already occur at tree-level. The QFV decay $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ [4] and QFV gluino decays [5] were studied in the scenario of minimal flavour violation (MFV), where the only source of QFV is the mixing due to the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Note that the decay $\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0$ is actually the standard Tevatron search mode for light top-squarks. In [6, 7] squark pair production and their decays at LHC have been analyzed including also the effect of the squark generation mixing.

In the present Letter, we study the effect of mixing between the second and third squark generations in its most general form. More precisely, we study the influence of the mixing of charm squark and top squark on the gluino and squark decays. In particular, we calculate the branching ratios of the following gluino decays into two quarks plus neutralino via up-type squark decay (see Fig.1) $^1$

$$\tilde{g} \rightarrow \tilde{u}_i c \rightarrow c t \tilde{\chi}_1^0 \quad \text{and} \quad \tilde{g} \rightarrow \tilde{u}_i t \rightarrow c t \tilde{\chi}_1^0. \quad (1)$$

$^1$As we always sum over the particles and antiparticles of the (s)quarks, we do not indicate if it is a particle or its antiparticle: $qq'$ (with $q \neq q'$) means $q\bar{q}'$ and $\bar{q}q'$, and $qq$ means $q\bar{q}$, e.g. $B(\tilde{g} \rightarrow ct\tilde{\chi}_1^0) = B(\tilde{g} \rightarrow c\bar{t}\tilde{\chi}_1^0) + B(\tilde{g} \rightarrow \bar{c}t\tilde{\chi}_1^0)$. 

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We show that the QFV gluino decay branching ratio \( B(\tilde{g} \rightarrow \tilde{c}_L \tilde{\chi}^0_1) \) can be very large (up to \( \sim 50\% \)) due to the squark generation mixing in a significant part of the MSSM parameter space despite the very strong experimental constraints from B factories, Tevatron and LEP \(^2\). We also study the effect of the squark generation mixing on the invariant mass distributions of the two quarks from the gluino decay at LHC. We show that it can result in novel multiple-edge structures in the distributions \(^3\). These effects could have an important impact on the search for gluinos and the MSSM parameter determination at LHC.

## 2 Squark mixing with flavour violation

Here we summarize the MSSM parameters in our analysis. The most general up-type squark mass matrix including left-right mixing as well as quark-flavour mixing in the super-CKM basis of \( \tilde{u}_0 = (\tilde{u}_L, \tilde{c}_L, \tilde{t}_L, \tilde{u}_R, \tilde{c}_R, \tilde{t}_R) \), \( \gamma = 1, \ldots, 6 \), is [10]

\[
M^2_{\tilde{u}} = \begin{pmatrix}
M^2_{\tilde{u}_{LL}} & (M^2_{\tilde{u}_{RL}})^\dagger \\
M^2_{\tilde{u}_{RL}} & M^2_{\tilde{u}_{RR}}
\end{pmatrix},
\]  

where the three \( 3 \times 3 \) matrices read

\[
(M^2_{\tilde{u}_{LL}})_{\alpha\beta} = M^2_{\tilde{u}_{LL}} + \left[ \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right] \cos 2\beta m^2_Z + m^2_{u_{\alpha\alpha}} \delta_{\alpha\beta},
\]

\[
(M^2_{\tilde{u}_{RR}})_{\alpha\beta} = M^2_{\tilde{u}_{RR}} + \left[ \frac{2}{3} \sin^2 \theta_W \cos 2\beta m^2_Z + m^2_{u_{\alpha\alpha}} \right] \delta_{\alpha\beta},
\]

\[
(M^2_{\tilde{u}_{RL}})_{\alpha\beta} = \left( v_2/\sqrt{2} \right) A_{U_{\beta\alpha}} - m_{u_{\alpha\alpha}} \mu^* \cot \beta \delta_{\alpha\beta}.
\]

The indices \( \alpha, \beta = 1, 2, 3 \) characterize the quark flavours \( u, c, t \), respectively. \( M^2_{\tilde{u}_u} \) and \( M^2_{\tilde{u}_c} \) are the Hermitean soft-SUSY-breaking mass matrices for the left and right up-type squarks, respectively. Note that in the super-CKM basis one has \( M^2_{\tilde{u}_u} = K \cdot M^2_{\tilde{Q}_u} \cdot K^\dagger \) due to the SU(2) symmetry, where \( M^2_{\tilde{Q}_u} \) is the Hermitean soft-SUSY-breaking mass matrix for the left up-type squark.

\(^2\)This is in analogy to the case of lepton flavour violating (LFV) sneutrino decays due to slepton generation mixing [8].

\(^3\)This is in analogy to the case of LFV neutralino decays due to slepton generation mixing [9].
for the left down-type squarks and $K$ is the CKM matrix. Note also that $M_{Q_u}^2 \simeq M_{Q}^2$ as $K \simeq 1$. $A_U$ is the soft-SUSY-breaking trilinear coupling matrix of the up-type squarks: $L_{int} = -(A_{U\alpha\beta}\bar{u}_{R\beta}u_{L\alpha}H_2^0 + h.c.) + \cdots$. $\mu$ is the higgsino mass parameter. $v_{1,2}$ are the vacuum expectation values of the Higgs fields with $v_{1,2}/\sqrt{2} \equiv \langle H_{1,2} \rangle$, and $\tan \beta \equiv v_2/v_1$. $m_{u_\alpha}$ ($u_\alpha = u, c, t$) are the physical quark masses.

The physical mass eigenstates $\tilde{u}_i$, $i = 1, \ldots, 6$, are given by $\tilde{u}_i = R_{\alpha \beta} \tilde{u}_0 \alpha$. The mixing matrix $R_{\tilde{u}}$ and the mass eigenvalues are obtained by a unitary transformation $R_{\tilde{u}} M_{\tilde{u}}^2 R_{\tilde{u}}^\dagger = \text{diag}(m_{\tilde{u}_1}, \ldots, m_{\tilde{u}_6})$, where $m_{\tilde{u}_i} < m_{\tilde{u}_j}$ for $i < j$.

The down-type squark mass matrix can be analogously parametrized as the up-type squark mass matrix [10]. As $M_{Q_u}^2 \simeq M_{Q}^2$, one has $(M_{\tilde{d}_{LL}}^2)^{\alpha \beta} \simeq (M_{\tilde{d}_{LL}}^2)^{\alpha \beta}$ for $\alpha \neq \beta$. We do not introduce additional QFV terms in the down-type squark mass matrix.

The properties of the charginos $\tilde{\chi}_i^{\pm}$ ($i = 1, 2, m_{\tilde{\chi}_1^{\pm}} < m_{\tilde{\chi}_2^{\pm}}$) and neutralinos $\tilde{\chi}_k^0$ ($k = 1, ... , 4, m_{\tilde{\chi}_1^0} < ... < m_{\tilde{\chi}_4^0}$) are determined by the parameters $M_2, M_1, \mu$ and $\tan \beta$, where $M_2$ and $M_1$ are the SU(2) and U(1) gaugino masses, respectively. Assuming gaugino mass unification including the gluino mass $m_{\tilde{g}} = M_3$, we take $M_1 = (5/3)\tan^2 \theta_W M_2$.

## 3 Constraints

In our analysis, we impose the following conditions on the MSSM parameter space in order to respect experimental and theoretical constraints:

(i) Constraints from the B-physics experiments relevant mainly for the mixing between the second and third generations of squarks $^4$:

$^4$We do not consider the experimental constraints from $b \to sg$ and $b \to s\nu\bar{\nu}$ since they have large uncertainties. We do not include the constraints from the experimental data on $B(B_d \to \mu^+\mu^-)$,
\[ 3.03 \times 10^{-4} < B(b \to s \gamma) < 4.01 \times 10^{-4} \text{ (95\% CL)} \] [12], \[ 0.60 \times 10^{-6} < B(b \to s l^+l^-) < 2.60 \times 10^{-6} \text{ with } l = e \text{ or } \mu \text{ (95\% CL)} \] [13], \[ B(B_s \to \mu^+\mu^-) < 4.8 \times 10^{-8} \text{ (90\% CL)} \] [12], \[ |R_{Br}^{SUSY} - 1.77| < 1.27 \text{ (95\% CL)} \] with \[ R_{Br}^{SUSY} \equiv B^{SUSY}(B_u^- \to \tau^-\bar{\nu}_\tau)/B^{SM}(B_u^- \to \tau^-\bar{\nu}_\tau) \approx (1 - (m_{\mu^+} \tan \beta)^2)^2 \] [14]. Moreover we impose the following condition on the SUSY prediction: \[ |\Delta M_{B_s}^{SUSY} - 17.77| < ((0.12 \times 1.96)^2 + 3.3^2)^{1/2} \text{ ps}^{-1} = 3.31 \text{ ps}^{-1} \text{ (95\% CL)}, \] where we have combined the experimental error of 0.12 ps\(^{-1}\) (at 68\% CL) [15] quadratically with the theoretical uncertainty of 3.3 ps\(^{-1}\) (at 95\% CL) [16].

(ii) The experimental limit on SUSY contributions to the electroweak $\rho$ parameter [17]: \[ \Delta \rho(SUSY) < 0.0012. \]

(iii) The LEP limits on the SUSY particle masses [18]: \[ m_{\tilde{\chi}^+} > 103 \text{ GeV}, \quad m_{\tilde{\chi}^0} > 50 \text{ GeV}, \quad m_{\tilde{q}_1, \tilde{d}_1} > 100 \text{ GeV}, \quad m_{\tilde{q}_2, \tilde{d}_2} > m_{\tilde{\chi}_1}, \quad m_{A^0} > 93 \text{ GeV}, \quad m_{h^0} > 110 \text{ GeV}, \] where $A^0$ is the CP-odd Higgs boson and $h^0$ is the lighter CP-even Higgs boson.

(iv) The Tevatron limit on the gluino mass [19]: \[ m_{\tilde{g}} > 308 \text{ GeV}. \]

(v) The vacuum stability conditions for the trilinear coupling matrix [20]:

\[
\begin{align*}
|A_{U\alpha\alpha}|^2 &< 3 Y_{U\alpha}^2 \left( M_{Q_{u\alpha\alpha}}^2 + M_{U\alpha\alpha}^2 + m_2^2 \right), \\
|A_{D\alpha\alpha}|^2 &< 3 Y_{D\alpha}^2 \left( M_{Q_{d\alpha\alpha}}^2 + M_{D\alpha\alpha}^2 + m_1^2 \right), \\
|A_{U\alpha\beta}|^2 &< Y_{U\gamma}^2 \left( M_{Q_{u\alpha\alpha}}^2 + M_{U\beta\beta}^2 + m_2^2 \right), \\
|A_{D\alpha\beta}|^2 &< Y_{D\gamma}^2 \left( M_{Q_{d\alpha\alpha}}^2 + M_{D\beta\beta}^2 + m_1^2 \right),
\end{align*}
\]

with \( \alpha \neq \beta; \gamma = \text{Max}(\alpha, \beta); \alpha, \beta = 1, 2, 3 \) and \[ m_1^2 = (m_{H^\pm}^2 + m_Z^2 \sin^2 \theta_W) \sin^2 \beta - \frac{1}{2} m_Z^2, \quad m_2^2 = (m_{H^\pm}^2 + m_Z^2 \sin^2 \theta_W) \cos^2 \beta - \frac{1}{2} m_Z^2. \] The Yukawa couplings of the up-type and down-type quarks are \[ Y_{U\alpha} = \sqrt{2} m_{ua}/v_2 = \frac{q_{ua}}{\sqrt{2} m_w \sin \beta} \quad (u_\alpha = u, c, t) \] and \[ Y_{D\alpha} = \sqrt{2} m_{da}/v_1 = \frac{q_{da}}{\sqrt{2} m_w \cos \beta} \quad (d_\alpha = d, s, b), \] with \( m_{ua} \) and \( m_{da} \) being the running quark masses at the scale of \( m_Z \) and \( g \) the SU(2) gauge coupling. All soft-SUSY-breaking parameters are assumed to be given at the scale of \( m_Z \). As SM input we take \( m_W = 80.4 \text{ GeV}, \quad m_Z = 91.2 \text{ GeV} \) and the on-shell top-quark mass \( m_t = 174.3 \text{ GeV} \). We have found that our results shown in the following are fairly insensitive to \( m_t \).

\( B(b \to d l^+l^-), \Delta M_{B_s}, \text{ and } \Delta M_{D^0} \) as they practically do not constrain the 2nd and 3rd generation squark mixing which we are interested in here.
We calculate the observables in (i)-(iii) by using the public code SPheno v3.0 [21]. Condition (i) except for \( B(B_u^+ \to \tau^+ \nu) \) strongly constrains the 2nd and 3rd generation squark mixing parameters \( M_{Q_{23}}^2, M_{D_{23}}^2, A_{U_{23}}, A_{D_{23}} \) and \( A_{D_{32}} \); the constraints from \( B(b \to s\gamma) \) and \( \Delta M_{B_s} \) are especially important [22].

4 Quark flavour violating gluino decays

We study the effect of the 2nd and 3rd generation squark mixing on the gluino decays. We focus on the QFV gluino decays of Eq.(1) leading to the same final state \( c \, t \, \tilde{\chi}^0_1 \).

We calculate the gluino and squark decay widths taking into account the following two–body decays:

\[
\begin{align*}
\tilde{g} & \to \tilde{u}_i \, u_k, \ \tilde{d}_i \, d_k, \\
\tilde{u}_i & \to u_k \tilde{\chi}_n^0, \ d_k \tilde{\chi}_m^+, \ \tilde{d}_j \ W^+, \ \tilde{u}_j \ Z^0, \ \tilde{u}_j \ h^0,
\end{align*}
\]

where \( u_k = (u, c, t) \) and \( d_k = (d, s, b) \). The squark decays into the heavier Higgs bosons are kinematically forbidden in our scenarios studied below. The formulae for the two–body decays in (13) can be found in [6], except for the squark decays into the Higgs bosons for which we take the formulae of [23] modified appropriately with the squark mixing matrix in the general QFV case.

We take \( \tan \beta, m_{A^0}, M_1, M_2, m_{\tilde{g}}, \mu, M_{Q_{\alpha \beta}}, M_{U_{\alpha \beta}}^2, M_{D_{\alpha \beta}}^2, A_{U_{\alpha \beta}} \) and \( A_{D_{\alpha \beta}} \) as the basic MSSM parameters at the weak scale. We assume them to be real. The QFV parameters are the squark generation mixing terms \( M_{Q_{\alpha \beta}}^2, M_{U_{\alpha \beta}}^2, M_{D_{\alpha \beta}}^2, A_{U_{\alpha \beta}} \) and \( A_{D_{\alpha \beta}} \) with \( \alpha \neq \beta \). As a reference scenario, we take the scenario given in Table 1. This scenario is within the reach of LHC and satisfies the conditions (i)-(v). For the observables in (i) and (ii) we obtain \( B(b \to s\gamma) = 3.57 \times 10^{-4}, B(b \to s l^+ l^-) = 1.59 \times 10^{-6}, B(b \to s v \bar{v}) = 4.07 \times 10^{-5}, B(B_s \to \mu^+ \mu^-) = 4.72 \times 10^{-9}, B(B_u^+ \to \tau^+ \nu) = 7.85 \times 10^{-5}, \Delta M_{B_s} = 17.38 \, ps^{-1} \) and \( \Delta \rho(SUSY) = 1.50 \times 10^{-4} \). The resulting masses of squarks, neutralinos and charginos are given in Table 2. We show the up-type squark compositions in the flavour eigenstates in Table 3.

For the important branching ratios of the gluino and squark two–body decays we get \( B(\tilde{g} \to \tilde{u}_1 c) = 0.481, B(\tilde{g} \to \tilde{u}_1 t) = 0.300, B(\tilde{g} \to \tilde{u}_2 c) = 0.207, B(\tilde{g} \to \tilde{u}_2 t) = 0.0, \) and \( B(\tilde{u}_1 \to c \tilde{\chi}_1^0) = 0.576, B(\tilde{u}_1 \to t \tilde{\chi}_1^0) = 0.401, B(\tilde{u}_2 \to c \tilde{\chi}_1^0) = 0.495, B(\tilde{u}_2 \to t \tilde{\chi}_1^0) = 0.0. \)
Table 1: The MSSM parameters in our reference scenario with QFV. All of $A_{U\alpha\beta}$ and $A_{D\alpha\beta}$ are set to zero. All mass parameters are given in GeV.

| $M^2_{Q\alpha\beta}$ | $\beta = 1$ | $\beta = 2$ | $\beta = 3$ |
|----------------------|-------------|-------------|-------------|
| $\alpha = 1$        | (920)$^2$   | 0           | 0           |
| $\alpha = 2$        | 0           | (880)$^2$   | (224)$^2$   |
| $\alpha = 3$        | 0           | (224)$^2$   | (840)$^2$   |

| $M^2_{D\alpha\beta}$ | $\beta = 1$ | $\beta = 2$ | $\beta = 3$ |
|----------------------|-------------|-------------|-------------|
| $\alpha = 1$        | (830)$^2$   | 0           | 0           |
| $\alpha = 2$        | 0           | (820)$^2$   | 0           |
| $\alpha = 3$        | 0           | 0           | (810)$^2$   |

Table 2: Sparticles and corresponding masses (in GeV) in the scenario of Table 1.

| $\tilde{u}_1$ | $\tilde{u}_2$ | $\tilde{u}_3$ | $\tilde{u}_4$ | $\tilde{u}_5$ | $\tilde{u}_6$ | $\tilde{d}_1$ | $\tilde{d}_2$ | $\tilde{d}_3$ | $\tilde{d}_4$ | $\tilde{d}_5$ | $\tilde{d}_6$ |
|--------------|---------------|---------------|---------------|---------------|---------------|--------------|---------------|---------------|---------------|---------------|---------------|
| 558          | 642           | 819           | 837           | 897           | 918           | 800          | 820           | 830           | 835           | 897           | 922           |

| $\tilde{\chi}^0_1$ | $\tilde{\chi}^0_2$ | $\tilde{\chi}^0_3$ | $\tilde{\chi}^0_4$ | $\tilde{\chi}^0_1\tilde{\chi}^0_1$ | $\tilde{\chi}^0_2\tilde{\chi}^0_2$ |
|--------------------|--------------------|--------------------|--------------------|-----------------------------|-----------------------------|
| 138                | 261                | 1003               | 1007               | 261                         | 1007                        |

$t_{\tilde{\chi}^0_1} = 0.469$. This leads to the following gluino decay branching ratios:

$$B(\tilde{g} \to ct_{\tilde{\chi}^0_1}) = \sum_{i=1,2} \left[ B(\tilde{g} \to \tilde{u}_i c)B(\tilde{u}_i \to t_{\tilde{\chi}^0_1}) + B(\tilde{g} \to \tilde{u}_i t)B(\tilde{u}_i \to c_{\tilde{\chi}^0_1}) \right] = 0.463, \quad (14)$$
$$B(\tilde{g} \to cc_{\tilde{\chi}^0_1}) = \sum_{i=1,2} \left[ B(\tilde{g} \to \tilde{u}_i c)B(\tilde{u}_i \to c_{\tilde{\chi}^0_1}) \right] = 0.380, \quad (15)$$
$$B(\tilde{g} \to tt_{\tilde{\chi}^0_1}) = \sum_{i=1,2} \left[ B(\tilde{g} \to \tilde{u}_i t)B(\tilde{u}_i \to t_{\tilde{\chi}^0_1}) \right] = 0.120. \quad (16)$$

Note that the QFV gluino decay branching ratio of Eq. (14) is very large. The reason of this very large QFV gluino decay branching ratio is as follows: The gluino decays into squarks other than $\tilde{u}_{1,2}$ are kinematically forbidden, and $\tilde{u}_1$, $\tilde{u}_2$ are strong mixtures of the flavour eigenstates $\tilde{c}_R$ and $\tilde{t}_R$ due to the large $\tilde{c}_R - \tilde{t}_R$ mixing term $M^2_{U23} = (224 \text{ GeV})^2$ in this scenario. This results in the large branching ratios of $B(\tilde{g} \to \tilde{u}_i c), B(\tilde{g} \to \tilde{u}_i t)$ and $B(\tilde{u}_i \to c_{\tilde{\chi}^0_1}), B(\tilde{u}_i \to t_{\tilde{\chi}^0_1})$ with $i = 1, 2$, except for the branching ratio of the decay $\tilde{g} \to \tilde{u}_2 t$ which is kinematically forbidden. Note
that $\tilde{u}_{1,2}(\sim \tilde{c}_R + \tilde{t}_R)$ couple to $\tilde{\chi}_1^0(\sim \tilde{B}^0)$ and practically do not couple to $\tilde{\chi}_2^0(\sim \tilde{W}^0)$, $\tilde{\chi}_1^\pm$, $\tilde{\chi}_2^\pm$ are very heavy in this scenario. Here $\tilde{B}^0$ and $\tilde{W}^0, \pm$ are the U(1) and SU(2) gauginos, respectively.

We now study the basic MSSM parameter dependences of the QFV gluino and squark decay branching ratios for the reference scenario of Table 1. In Fig.2 we show contours of $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ in the $(\Delta M^2_U, M^2_{U23})$ plane with $\Delta M^2_U \equiv M^2_{U22} - M^2_{U33}$. All basic parameters other than $M^2_{Q23}$ and $M^2_{U23}$ are fixed as in our reference scenario defined in Table 1. We see that the QFV decay branching ratio $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ quickly increases up to $\sim 50\%$ with increase of the effective $\tilde{c}_R - \tilde{t}_R$ mixing angle $\tan(2\theta_{23}^{eff}) \equiv 2 M^2_{U23}/\Delta M^2_U$.

In Fig.3 we present contours of $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ in the $\delta_{23}^{uLL} - \delta_{23}^{uRR}$ plane where all of the conditions (i)-(v) except the $b \to s \gamma$ constraint are satisfied. For $b \to s \gamma$ we also show the corresponding branching ratio contours. All basic parameters other than $M^2_{Q23}$ and $M^2_{U23}$ are fixed as in our reference scenario defined in Table 1. We see that the QFV decay branching ratio $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ increases quickly with increase of the $\tilde{c}_R - \tilde{t}_R$ mixing parameter $|\delta_{23}^{uRR}|$ and can be very large (up to $\sim 50\%$) in a significant part of the $\delta_{23}^{uLL} - \delta_{23}^{uRR}$ plane allowed by all of the conditions (i)-(v) including the $b \to s \gamma$ constraint. $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ is insensitive to the $\tilde{c}_L - \tilde{t}_L$ mixing parameter $\delta_{23}^{uLL}$ and can be quite large ($\sim 50\%$) in a sizable allowed range $0.03 \lesssim \delta_{23}^{uLL} \lesssim 0.12$.

Studying the branching ratios of the gluino and up-type squark two-body decays separately allows for a better understanding of their contributions to the QFV gluino

| $R^\tilde{u}_{i\alpha}$ | $\tilde{u}_L$ | $\tilde{c}_L$ | $\tilde{t}_L$ | $\tilde{u}_R$ | $\tilde{c}_R$ | $\tilde{t}_R$ |
|------------------------|-------------|-------------|-------------|-------------|-------------|-------------|
| $\tilde{u}_1$          | -0.001     | 0.005       | -0.029      | 0           | 0.728       | -0.685      |
| $\tilde{u}_2$          | -0.002     | 0.008       | -0.040      | 0           | -0.686      | -0.727      |
| $\tilde{u}_3$          | 0          | 0           | 0           | 1.0         | 0           | 0           |
| $\tilde{u}_4$          | 0.128      | -0.583      | 0.801       | 0           | -0.007      | -0.045      |
| $\tilde{u}_5$          | -0.181     | 0.782       | 0.597       | 0           | -0.003      | -0.021      |
| $\tilde{u}_6$          | -0.975     | -0.221      | -0.005      | 0           | 0           | 0           |

Table 3: The up-type squark compositions in the flavour eigenstates, i.e. the mixing matrix $R^\tilde{u}_{i\alpha}$ for the scenario of Table 1.
decay $\tilde{g} \to ct\tilde{\chi}_1^0$. In Fig.4 we show the $\delta_{23}^{uRR}$ (i.e. $c_R - \tilde{t}_R$ mixing parameter) dependences of the gluino and squark decay branching ratios, where all basic parameters other than $M_{U_{23}}^2$ are fixed as in the scenario of Table 1. We see that $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ increases quickly with increase of $|\delta_{23}^{uRR}|$ for $|\delta_{23}^{uRR}| \lesssim 0.1$ and can be very large ($\sim 50\%$) in a wide range of $\delta_{23}^{uRR}$. This behaviour can be explained by an argument similar to that below Eq.(16). In Fig.4(b) [(c)] we see that $B(\tilde{g} \to \tilde{u}_i c)$ and $B(\tilde{g} \to \tilde{u}_it) [B(\tilde{u}_i \to c\tilde{\chi}_1^0)$ and $B(\tilde{u}_i \to t\tilde{\chi}_1^0)]$ with $i = 1,2$ are large in a wide range of $\delta_{23}^{uRR}$, except for $B(\tilde{g} \to \tilde{u}_2t)$ which is kinematically suppressed. This leads to the very large $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ in a wide range of $\delta_{23}^{uRR}$ (see Eq.(14)).

In Fig.5 we show the $\delta_{23}^{uRL}$ (i.e. $c_R - \tilde{t}_L$ mixing parameter) dependences of the gluino decay branching ratios, where all basic parameters other than $A_{U_{32}}$ are fixed as in the scenario of Table 1. We see that the QFV decay branching ratio $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ can be quite large ($\sim 30-50\%$) in a wide range of $\delta_{23}^{uRL}$. $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ decreases (down to $\sim 30\%$) and the QGV decay branching ratio $B(\tilde{g} \to cb\tilde{\chi}_1^0)$ increases (up to $\sim 20\%$) with increase of $|\delta_{23}^{uRL}|$. Sizable $\delta_{23}^{uRL}$ (i.e. $c_R - \tilde{t}_L$ mixing parameter) induces a sizable $\tilde{t}_L$ component in $\tilde{u}_{1,2}(\sim c_R + \tilde{t}_R)$, which enhances the widths $\Gamma(\tilde{u}_{1,2} \to b\tilde{\chi}_1^\pm(\sim \tilde{W}^\pm))$ and leads to a suppression of $B(\tilde{u}_{1,2} \to c\tilde{\chi}_1^0(\sim \tilde{B}^0))$ and $B(\tilde{u}_{1,2} \to t\tilde{\chi}_1^0)$. As a result $B(\tilde{g} \to cb\tilde{\chi}_1^\pm) = \sum_{i=1,2} B(\tilde{g} \to \tilde{u}_ic)B(\tilde{u}_i \to b\tilde{\chi}_1^\pm)$ is enhanced for sizable $\delta_{23}^{uRL}$ while $B(\tilde{g} \to ct\tilde{\chi}_1^0) = \sum_{i=1,2} [B(\tilde{g} \to \tilde{u}_ic)B(\tilde{u}_i \to t\tilde{\chi}_1^0) + B(\tilde{g} \to \tilde{u}_it)B(\tilde{u}_i \to c\tilde{\chi}_1^0)]$ is suppressed.

As for the $\delta_{32}^{uRL}$ (i.e. $c_L - \tilde{t}_R$ mixing parameter) dependence of the gluino decay branching ratios, we have obtained similar results to those for the $\delta_{23}^{uRL}$ dependence in Fig.5. We have found that $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ can be quite large ($\sim 30-50\%$) in a wide allowed range $|\delta_{32}^{uRL}| \lesssim 0.3$. $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ decreases (down to $\sim 30\%$) and the QGV decay branching ratio $B(\tilde{g} \to st\tilde{\chi}_1^\pm)$ increases (up to $\sim 5\%$) with the increase of $|\delta_{32}^{uRL}|$ while $B(\tilde{g} \to cb\tilde{\chi}_1^\pm)$ is small.

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5Note that gluino decays into a down-type squark, such as $B(\tilde{g} \to \tilde{d}_ib)$, are kinematically forbidden in this scenario and hence that such decays cannot contribute to $B(\tilde{g} \to cb\tilde{\chi}_1^\pm)$. 
5 Impact on collider signatures

Here we study the invariant mass distributions (i.e. the differential decay branching ratios) $d\text{Br}(\tilde{g} \to \tilde{u}_i u_j \to u_j u_k \tilde{\chi}_n^0)/dM_{u_j u_k}$ with $M_{u_j u_k}$ being the invariant mass of the two quark system $u_j u_k$ in the final state. The kinematical endpoints of the distributions are given in terms of the masses of the involved particles by [24]

$$M_{u_j u_k}^{\text{min}, \text{max}} = \left\{ m_{u_j}^2 + m_{u_k}^2 + \frac{1}{2m_{\tilde{u}_i}^2} \left[ (m_{\tilde{g}}^2 - m_{u_j}^2 - m_{\tilde{u}_i}^2)(m_{u_k}^2 + m_{u_k}^2 - m_{\tilde{\chi}_n}^2) + \lambda^2 \left( m_{\tilde{g}}^2, m_{u_j}^2, m_{\tilde{u}_i}^2 \right) \lambda^2 \left( m_{u_k}^2, m_{u_k}^2, m_{\tilde{\chi}_n}^2 \right) \right] \right\}^{1/2}.$$  

(17)

with $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + xz + yz)$, where $\tilde{u}_i$ is the intermediate squark, $u_j$ is from the primary decay (i.e. the two-body $\tilde{g}$ decay) and $u_k$ is from the secondary decay (i.e. the $\tilde{u}_i$ decay). Note that $M_{u_j u_k}^{\text{min}, \text{max}} \neq M_{u_k u_j}^{\text{min}, \text{max}}$ for $j \neq k$. We calculate the invariant mass distributions by summing over the intermediate up-type squarks giving rise to the same final state:

$$d\text{Br}(\tilde{g} \to u_j u_k \tilde{\chi}_n^0)/dM_{u_j u_k} = \frac{1}{1 + \delta_{jk}} \sum_i \left[ d\text{Br}(\tilde{g} \to \tilde{u}_i u_j \to u_j u_k \tilde{\chi}_n^0)/dM_{u_j u_k} \right.$$

$$+ d\text{Br}(\tilde{g} \to \tilde{u}_i u_k \to u_k u_j \tilde{\chi}_n^0)/dM_{u_j u_k} \left. \right].$$  

(18)

Note that the individual distribution $d\text{Br}(\tilde{g} \to \tilde{u}_i u_j \to u_j u_k \tilde{\chi}_n^0)/dM_{u_j u_k}$, is proportional to $M_{u_j u_k}^{\text{min}}$ and its allowed range is given by $\left[ M_{u_j u_k}^{\text{min}}, M_{u_j u_k}^{\text{max}} \right]$.

In the following we show how QFV due to the 2nd and 3rd generation mixing of the up-type squarks influences the invariant mass distributions. We discuss two scenarios, one with gluino mass $m_{\tilde{g}} = 800$ GeV and the other with $m_{\tilde{g}} = 1300$ GeV.

We start from the QFV scenario with $m_{\tilde{g}} = 800$ GeV given in Table 1. In this QFV scenario the squark mass eigenstates $\tilde{u}_1$ and $\tilde{u}_2$ are a strong mixture of the flavour eigenstates $\tilde{c}_R$ and $\tilde{t}_R$. First we consider the invariant mass distribution for a final state including two top quarks. Fig.6 shows the invariant mass distributions of the top quark pairs for the QFV scenario, where one has $B(\tilde{g} \to t\tilde{t}\tilde{\chi}_n^0) = 12.0\%$. Note that the invariant mass distribution of the two top quarks in the QFV scenario shows no additional edge structure. This is because only the lightest up-type squark, $\tilde{u}_1$, can
mediate this final state while the other squarks are too heavy.

Next we consider the invariant mass distribution for a final state including c and t quarks in the QFV scenario of Table 1, where one has \( B(\tilde{g} \rightarrow ct\tilde{\chi}^0_1) = 46.3\% \). Fig. 6 shows the invariant mass distribution of \( ct \). There are more edge structures due to the processes \( \tilde{g} \rightarrow \tilde{u}_1 t \rightarrow tc\tilde{\chi}^0_1 \) [with \( M^{(\text{min,max})}_{ct} = (253, 526) \) GeV], \( \tilde{g} \rightarrow \tilde{u}_1 c \rightarrow ct\tilde{\chi}^0_1 \) [with \( M^{(\text{min,max})}_{ct} = (254, 580) \) GeV], and \( \tilde{g} \rightarrow \tilde{u}_2 c \rightarrow ct\tilde{\chi}^0_1 \) [with \( M^{(\text{min,max})}_{ct} = (219, 497) \) GeV]. Note that \( \tilde{g} \rightarrow \tilde{u}_2 t \) is kinematically forbidden in this scenario. We see that the three remarkable endpoint-edges are fairly well separated.

Next we consider the invariant mass distribution of final state quarks for a QFV scenario with a heavier gluino \( (m_{\tilde{g}} = 1300 \) GeV) given in Table 4. This scenario is inspired by the mSUGRA scenario A of Ref. [25] and satisfies all of the conditions (i)-(v) in section 3. The resulting masses of squarks, neutralinos and charginos are given in Table 5. We show the corresponding up-type squark compositions in the flavour eigenstates in Table 6. In this scenario the squark mass eigenstate \( \tilde{u}_1 (\tilde{u}_2) \) is dominated by a strong mixture of the flavour eigenstates \( \tilde{t}_R \) and \( \tilde{c}_R \) (\( \tilde{t}_L \) and \( \tilde{c}_L \)). In Fig. 7 we show the two invariant mass distributions of \( tt \) and \( ct \), where one has \( B(\tilde{g} \rightarrow tt\tilde{\chi}^0_1) = 16.6\% \), and \( B(\tilde{g} \rightarrow ct\tilde{\chi}^0_1) = 31.4\% \). Note that the QFV decay branching ratio \( B(\tilde{g} \rightarrow ct\tilde{\chi}^0_1) \) is large.

The invariant mass distribution of two top quarks shows no additional edge structure for the same reason as in the scenario with \( m_{\tilde{g}} = 800 \) GeV discussed above. The decay \( \tilde{g} \rightarrow \tilde{u}_2 t \) is kinematically allowed but phase-space suppressed. Moreover, \( \tilde{u}_2 \rightarrow t\tilde{\chi}^0_1 \) is strongly suppressed because \( \tilde{u}_2 (\sim \tilde{t}_L + \tilde{c}_L) \) does not significantly couple to \( \tilde{\chi}^0_1 (\sim \tilde{B}^0(\text{Bino})) \) in this scenario. Hence, \( B(\tilde{g} \rightarrow \tilde{u}_2 t \rightarrow tt\tilde{\chi}^0_1) (=0.00035) \) is very small.

As for the invariant mass distribution of c and t quarks in the QFV scenario of Table 4, there are more edge structures due to the \( \tilde{u}_1 \)-mediated processes \( \tilde{g} \rightarrow \tilde{u}_1 t \rightarrow tc\tilde{\chi}^0_1 \) [with \( M^{(\text{min,max})}_{ct} = (601, 971) \) GeV], and \( \tilde{g} \rightarrow \tilde{u}_1 c \rightarrow ct\tilde{\chi}^0_1 \) [with \( M^{(\text{min,max})}_{ct} = (183, 1022) \) GeV]. The decays \( \tilde{g} \rightarrow \tilde{u}_2 c/t \) are phase-space suppressed and the decays \( \tilde{u}_2 \rightarrow c/t\tilde{\chi}^0_1 \) are strongly suppressed in this scenario as is explained above. Hence, \( B(\tilde{g} \rightarrow \tilde{u}_2 c/t \rightarrow ct\tilde{\chi}^0_1) (=0.0004) \) is very small.

Finally, we briefly discuss the measurability of the QFV decay \( \tilde{g} \rightarrow c t \tilde{\chi}^0_1 \) at LHC. It is important whether one can discriminate between the QFV decay \( \tilde{g} \rightarrow c t \tilde{\chi}^0_1 \) and the QFC decay \( \tilde{g} \rightarrow t t \tilde{\chi}^0_1 \). Therefore, it is necessary to identify the top quarks in the final states. This is possible by using the decay \( t \rightarrow bW \) with the W decaying into
two jets. For this purpose, a special method was proposed in [24], where it is assumed that the masses of the gluino and the $\tilde{\chi}_1^0$ are known from other measurements. The signature of the decay $\bar{g} \rightarrow c \bar{t} \tilde{\chi}_1^0$ would be 'charm-jet + top-quark + missing-energy'. Therefore, charm-tagging also would be very useful. If this is not possible, one should search for the decay $\bar{g} \rightarrow q \bar{t} \tilde{\chi}_1^0$ ($q \neq t$), i.e. for the signature 'jet + top-quark + missing-energy'. In the scenarios discussed, the most important SUSY background would be due to the QFC decay $\bar{g} \rightarrow t \tilde{\chi}_1^0$ and the pair production of the lightest up-type squarks, $pp \rightarrow \tilde{u}_1 + \tilde{u}_1 + X$, with $\tilde{u}_1 \rightarrow c \tilde{\chi}_1^0$ and $\tilde{u}_1 \rightarrow t \tilde{\chi}_1^0$. The most important SM background would be top-quark pair production. For the measurement of the endpoints in the multiple edge structure a good energy/momentum resolution of the detector would be necessary. In any case, one should take into account the possibility of significant contributions from QFV decays in the gluino search. Moreover one should also include the QFV squark parameters in the determination of the basic
Table 6: The up-type squark compositions in the flavour eigenstates, i.e. the mixing matrix $R_{\tilde{u}ia}$ for the scenario of Table 4.

SUSY parameters at LHC. It is clear that detailed Monte Carlo studies taking into account backgrounds and detector simulations would be necessary. Such studies are beyond the scope of the present article.

6 Conclusion

To conclude, we have studied gluino decays in the MSSM with squark mixing of the second and third generation, especially $\tilde{c}_{L/R} - \tilde{t}_{L/R}$ mixing. We have shown that QFV gluino decay branching ratios such as $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ can be very large due to the squark mixing in a significant part of the MSSM parameter space despite the very strong experimental constraints from B factories, Tevatron and LEP with those of $b \to s \gamma$ and $\Delta M_B$ being especially important.

We have also studied the effect of the squark generation mixing on the invariant mass distributions of the two quarks from the gluino decay at LHC. We have found that it can result in novel and characteristic edge structures in the distributions. In particular, multiple-edge (3- or 4-edge) structures can appear in the charm-top quark mass distribution. The appearance of these remarkable structures would provide an additional powerful test of supersymmetric QFV at LHC.

These could have an important impact on the search for gluinos and the MSSM parameter determination at LHC.
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References

[1] see e.g. J. J. Liu et al., Nucl. Phys. B 705 (2005) 3 [arXiv:hep-ph/0404099]; S. Bejar, J. Guasch and J. Sola, JHEP 0510 (2005) 113 [arXiv:hep-ph/0508043]; G. Eilam, M. Frank and I. Turan, Phys. Rev. D 74 (2006) 035012 [arXiv:hep-ph/0601253]; J. J. Cao et al., Phys. Rev. D 75 (2007) 075021 [arXiv:hep-ph/0702264]; D. Lopez-Val, J. Guasch and J. Sola, JHEP 0712 (2007) 054 [arXiv:0710.0587 [hep-ph]]; S. Bejar et al., Phys. Lett. B 668 (2008) 364 [arXiv:0805.0973 [hep-ph]].

[2] see e.g. C. S. Li, R. J. Oakes and J. M. Yang, Phys. Rev. D 49 (1994) 293 [Erratum-ibid. D 56 (1997) 3156]; G. Couture, C. Hamzaoui and H. Konig, Phys. Rev. D 52 (1995) 1713 [arXiv:hep-ph/9410230]; J. M. Yang and C. S. Li, Phys. Rev. D 49 (1994) 3412 [Erratum-ibid. D 51 (1995) 3974]; J. L. Lopez, D. V. Nanopoulos and R. Rangarajan, Phys. Rev. D 56 (1997) 3100 [arXiv:hep-ph/9702350]; J. Guasch and J. Sola, Nucl. Phys. B 562 (1999) 3 [arXiv:hep-ph/9906268]; J. L. Diaz-Cruz, H. J. He and C. P. Yuan, Phys. Lett. B 530 (2002) 179 [arXiv:hep-ph/0103178]; J. Cao et al., Phys. Rev. D 74 (2006) 031701 [arXiv:hep-ph/0604163].

[3] A. M. Curiel, M. J. Herrero and D. Temes, Phys. Rev. D 67 (2003) 075008 [arXiv:hep-ph/0210335]; D. A. Demir, Phys. Lett. B 571 (2003) 193 [arXiv:hep-
ph/0303249]; A. M. Curiel et al., Phys. Rev. D 69 (2004) 075009 [arXiv:hep-ph/0312135]; S. Bejar et al., JHEP 0408 (2004) 018 [arXiv:hep-ph/0402188]; T. Hahn et al., arXiv:hep-ph/0512315.

[4] K. i. Hikasa and M. Kobayashi, Phys. Rev. D 36 (1987) 724; T. Han et al., Phys. Rev. D 70 (2004) 055001 [arXiv:hep-ph/0312129]; E. Lunghi, W. Porod and O. Vives, Phys. Rev. D 74 (2006) 075003 [arXiv:hep-ph/0605177].

[5] W. Porod, JHEP 0205 (2002) 030 [arXiv:hep-ph/020259].

[6] G. Bozzi et al., Nucl. Phys. B 787 (2007) 1 [arXiv:0704.1826 [hep-ph]].

[7] B. Fuks, B. Herrmann and M. Klasen, Nucl. Phys. B 810 (2009) 266 [arXiv:0808.1104 [hep-ph]]; G. D. Kribs, A. Martin and T. S. Roy, arXiv:0901.4105 [hep-ph].

[8] A. Bartl, K. Hidaka, K. Hohenwarter-Sodek, T. Kernreiter, W. Majerotto, W. Porod, Phys. Lett. B 660 (2008) 228 [arXiv:0709.1157 [hep-ph]].

[9] A. Bartl, K. Hidaka, K. Hohenwarter-Sodek, T. Kernreiter, W. Majerotto, W. Porod, Eur. Phys. J. C 46 (2006) 783 [arXiv:hep-ph/0510074].

[10] B. Allanach et al., Comput. Phys. Commun. 180 (2009) 8 [arXiv:0801.0045 [hep-ph]].

[11] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B 477 (1996) 321 [arXiv:hep-ph/9604387].

[12] P. Chang, plenary talk at The 34th International Conference on High Energy Physics (ICHEP 2008), Philadelphia, 30 July to 5 August 2008; E. Barberio et al. [HFAG Collaboration], arXiv:0808.1297 [hep-ex].

[13] M. Iwasaki et al. [Belle Collaboration], Phys. Rev. D72 (2005) 092005 [arXiv:hep-ex/0503044]; B. Aubert et al. [BABAR Collaboration], Phys. Rev. Lett. 93 (2004) 081802 [arXiv:hep-ex/0404006]; see also E. Lunghi, W. Porod and O. Vives, Phys. Rev. D74 (2006) 075003 [arXiv:hep-ph/0605177].

[14] K. Hara, talk at The 34th International Conference on High Energy Physics (ICHEP 2008), Philadelphia, 30 July to 5 August 2008 [arXiv:0810.3301 [hep-ex]].
[15] A. Abulencia et al. [CDF Collaboration], Phys. Rev. Lett. 97 (2006) 242003; see also [12].

[16] M. Carena et al., Phys. Rev. D 74 (2006) 015009 [arXiv:hep-ph/0603106]; see also P. Ball and R. Fleischer, Eur. Phys. J. C 48 (2006) 413 [arXiv:hep-ph/0604249].

[17] G. Altarelli, R. Barbieri, F. Caravaglios, Int. J. Mod. Phys. A13 (1998) 1031.

[18] G. Sguazzoni, Proceedings of the 31st International Conference on High Energy Physics, Amsterdam, The Netherlands, 25 - 31 July 2002, Eds. S. Bentvelson, P. de Jong, J. Koch, E. Laenen, p. 709 [arXiv:hep-ex/0210022]; P. Lutz, the same Proceedings, p. 735; S. Schael et al., ALEPH Collaboration and DELPHI Collaboration and L3 Collaboration and OPAL Collaborations and LEP Working Group for Higgs Boson Searches, Eur. Phys. J. C 47 (2006) 547.

[19] K. Peters, talk at The 34th International Conference on High Energy Physics (ICHEP 2008), Philadelphia, 30 July to 5 August 2008.

[20] J. A. Casas and S. Dimopoulos, Phys. Lett. B 387 (1996) 107 [arXiv:hep-ph/9606237].

[21] The code SPheno v3.0 can be obtained from http://theorie.physik.uni-wuerzburg.de/~porod/SPheno.html; W. Porod, Comput. Phys. Commun. 153 (2003) 275 [arXiv:hep-ph/0301101].

[22] T. Hurth and W. Porod, arXiv:0904.4574 [hep-ph].

[23] A. Bartl, S. Hesselbach, K. Hidaka, T. Kernreiter and W. Porod, Phys. Rev. D 70 (2004) 035003 [arXiv:hep-ph/0311338].

[24] J. Hisano, K. Kawagoe, R. Kitano and M. M. Nojiri, Phys. Rev. D 66 (2002) 115004 [arXiv:hep-ph/0204078]; J. Hisano, K. Kawagoe and M. M. Nojiri, Phys. Rev. D 68 (2003) 035007 [arXiv:hep-ph/0304214].

[25] M. Battaglia et al., Eur. Phys. J. C 22 (2001) 535 [arXiv:hep-ph/0106204].
Figure 1: Feynman diagrams for $\tilde{g} \to \tilde{u}_i c \to c t \tilde{\chi}_1^0$ (left) and $\tilde{g} \to \tilde{u}_i t \to c t \tilde{\chi}_1^0$ (right).

Figure 2: Contours of the QFV decay branching ratio $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ in the $(\Delta M^2_U, M^2_{U23})$ plane where all of the conditions (i)-(v) are satisfied. The point ”x” of $(\Delta M^2_U, M^2_{U23}) = (2.36 \times 10^4, 5 \times 10^4)$ GeV$^2$ corresponds to our reference scenario of Table 1.

Figure 3: Contours of the QFV decay branching ratio $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ (solid lines) in the $\delta^{uLL}_{23} - \delta^{uRR}_{23}$ plane where all of the conditions (i)-(v) except the $b \to s \gamma$ constraint are satisfied. Contours of $10^4 \times B(b \to s \gamma)$ (dashed lines) are also shown. The condition (i) requires $3.03 < 10^4 \times B(b \to s \gamma) < 4.01$. The point ”x” of $(\delta^{uLL}_{23}, \delta^{uRR}_{23}) = (0.068, 0.144)$ corresponds to our reference scenario of Table 1.

Figure 4: $\delta^{uRR}_{23}$ dependences of the branching ratios of (a) the gluino cascade decays, (b) the gluino two-body decays and (c) the up-type squark two-body decays. The point ”x” of $\delta^{uRR}_{23} = 0.144$ corresponds to our reference scenario of Table 1. The shown range of $\delta^{uRR}_{23}$ is the whole range allowed by the conditions (i) to (v) given in the text; note that the range $|\delta^{uRR}_{23}| > 1.0$ is excluded by the condition $m_{\tilde{u}_1} > m_{\tilde{\chi}_1^0}$ in (iii).

Figure 5: $\delta^{uRL}_{23}$ dependences of the branching ratios of the gluino cascade decays. The point ”x” of $\delta^{uRL}_{23} = 0$ corresponds to our reference scenario of Table 1. The shown range of $\delta^{uRL}_{23}$ is the whole range allowed by the conditions (i) to (v) given in the text; note that the range $|\delta^{uRL}_{23}| > 0.3$ is excluded by the condition (v).

Figure 6: Invariant mass distributions of two up-type quarks from the decay $\tilde{g} \to u_j u_k \tilde{\chi}_1^0$ for the QFV scenario of Table 1.

Figure 7: Invariant mass distributions of two up-type quarks from the decay $\tilde{g} \to u_j u_k \tilde{\chi}_1^0$ for the QFV scenario of Table 4.
Fig. 1
Fig. 2

Fig. 3
Fig. 4
Fig. 5
100 d\text{Br}(\tilde{g} \rightarrow u_j u_k \tilde{\chi}_1^0)/dM_{u_j u_k} [\text{GeV}^{-1}] \\

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig6}
\caption{Fig. 6}
\end{figure}

100 d\text{Br}(\tilde{g} \rightarrow u_j u_k \tilde{\chi}_1^0)/dM_{u_j u_k} [\text{GeV}^{-1}] \\

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig7}
\caption{Fig. 7}
\end{figure}