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

A. Bartl[1,2], K. Hidaka[3], K. Hohenwarter-Sodek[1], T. Kernreiter[4], W. Majerotto[5] and W. Porod[6]

[1] *Faculty of Physics, Universität Wien, A-1090 Vienna, Austria*

[2] *AHEP Group, Instituto de Fisica Corpuscular - C.S.I.C., Universidad de Valencia, Edificio Institutos de Investigacion, Apt. 22085, E-46071 Valencia, Spain*

[3] *Department of Physics, Tokyo Gakugei University, Koganei, Tokyo 184-8501, Japan*

[4] *Departamento de Física and CFTP, Instituto Superior Técnico Av. Rovisco Pais 1, 1049-001 Lisboa, Portugal*

[5] *Institut für Hochenergiephysik der Österreichischen Akademie der Wissenschaften, A-1050 Vienna, Austria*

[6] *Institut für Theoretische Physik and Astrophysik, Universität Würzburg, D-97074 Würzburg, Germany*

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 $\tilde{g} \to c\ \bar{t}\ (\bar{c}\ t)\ \tilde{\chi}_1^0$ 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 very 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 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 6×6 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 $\tilde{t}_1 \to c\tilde{\chi}_1^0$ [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. 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 paper, 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 [1]:

$$\begin{aligned} \tilde{g} &\to \tilde{u}_i c \to ct\tilde{\chi}_j^0 \\ \tilde{g} &\to \tilde{u}_i t \to ct\tilde{\chi}_j^0. \end{aligned} \tag{1}$$

[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 anti-particle: qq' (with $q \neq q'$) means $q\bar{q'}$ and $\bar{q}q'$, and qq means $q\bar{q}$, e.g. $B(\tilde{g} \to ct\tilde{\chi}_1^0) \equiv B(\tilde{g} \to c\bar{t}\tilde{\chi}_1^0) + B(\tilde{g} \to \bar{c}t\tilde{\chi}_1^0)$.

We show that the QFV gluino decay branching ratio $\mathrm{B}(\tilde{g} \to ct\tilde{\chi}_1^0)$ 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

First 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 conventional super-CKM basis of the quark-flavour eigenstates $\tilde{u}_{0\gamma} = (\tilde{u}_L, \tilde{c}_L, \tilde{t}_L, \tilde{u}_R, \tilde{c}_R, \tilde{t}_R)$, $\gamma = 1, \ldots, 6$, is [10]

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

where the three 3×3 matrices read

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

$$(M^2_{\tilde{u}RR})_{\alpha\beta} = M^2_{U\alpha\beta} + \left[\frac{2}{3}\sin^2\theta_W \cos 2\beta\; m_Z^2 + m^2_{u_\alpha}\right]\delta_{\alpha\beta}\;, \qquad (4)$$

$$(M^2_{\tilde{u}RL})_{\alpha\beta} = (v_2/\sqrt{2})A_{U\beta\alpha} - m_{u_\alpha}\mu^*\cot\beta\;\delta_{\alpha\beta}\;. \qquad (5)$$

The indices $\alpha, \beta = 1, 2, 3$ characterize the quark flavours u, c, t, respectively. $M^2_{Q_u}$ and M^2_U 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 $M^2_{Q_u} = K \cdot M^2_Q \cdot K^\dagger$

[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].

due to the SU(2) symmetry, where M^2_Q is the hermitean soft-SUSY-breaking mass matrix for the left down-type squarks and K is the CKM matrix. Note also that $M^2_{Q_u} \simeq M^2_Q$ as $K \simeq 1$. A_U is the soft-SUSY-breaking trilinear coupling matrix of the up-type squarks: $\mathcal{L}_{\rm int} = -(A_{U\alpha\beta}\tilde{u}^\dagger_{R\beta}\tilde{u}_{L\alpha}H^0_2 + h.c.) + \cdots$. μ is the higgsino mass parameter. $v_{1,2}$ are the vacuum expectation values of the Higgs fields with $v_{1,2}/\sqrt{2} \equiv \langle H^0_{1,2}\rangle$, and $\tan\beta \equiv v_2/v_1$. m_{u_α} ($u_\alpha = u, c, t$) are the physical quark masses.
The physical mass eigenstates $\tilde{u}_i$, $i = 1, \ldots, 6$, are given by $\tilde{u}_i = R^{\tilde{u}}_{i\alpha}\tilde{u}_{0\alpha}$. The 6×6 mixing matrix $R^{\tilde{u}}$ and the mass eigenstates $\tilde{u}_i$ are obtained by an unitary transformation $R^{\tilde{u}} M^2_{\tilde{u}} R^{\tilde{u}\dagger} = {\rm diag}(m_{\tilde{u}_1}, \ldots, m_{\tilde{u}_6})$, where $m_{\tilde{u}_i} < m_{\tilde{u}_j}$ for $i < j$. Quark-flavour violation is induced by off–diagonal entries in the matrices $M^2_{Q_u}$, M^2_U and A_U, i.e. squark generation mixing terms. For instance, a non–zero A_{U32} (A_{U23}) gives rise to $\tilde{c}_R - \tilde{t}_L$ ($\tilde{t}_R - \tilde{c}_L$) mixing. Having in mind that $M^2_{Q_u} \simeq M^2_Q$, we define the QFV parameters $\delta^{uLL}_{\alpha\beta}$, $\delta^{uRR}_{\alpha\beta}$ and $\delta^{uRL}_{\alpha\beta}$ ($\alpha \neq \beta$) as follows [11]:

$$\delta^{uLL}_{\alpha\beta} \equiv M^2_{Q\alpha\beta}/\sqrt{M^2_{Q\alpha\alpha}M^2_{Q\beta\beta}}\ , \tag{6}$$

$$\delta^{uRR}_{\alpha\beta} \equiv M^2_{U\alpha\beta}/\sqrt{M^2_{U\alpha\alpha}M^2_{U\beta\beta}}\ , \tag{7}$$

$$\delta^{uRL}_{\alpha\beta} \equiv (v_2/\sqrt{2})A_{U\beta\alpha}/\sqrt{M^2_{U\alpha\alpha}M^2_{Q\beta\beta}}\ . \tag{8}$$

The down-type squark mass matrix can be analogously parametrized as the up-type squark mass matrix. Note that due to the SU(2) symmetry relation $M^2_{Q_u} = K \cdot M^2_Q \cdot K^\dagger \simeq M^2_Q$ the elements in the left–left block of the up-type squark mass matrix and the down-type squark mass matrix are not independent: one has $(M^2_{\tilde{u}LL})_{\alpha\beta} \simeq (M^2_{\tilde{d}LL})_{\alpha\beta}$ for $\alpha \neq \beta$. We do not introduce additional QFV terms (i.e. squark generation mixing terms) in the down-type squark mass matrix.

The properties of the charginos $\tilde{\chi}^\pm_i$ ($i = 1, 2$, $m_{\tilde{\chi}^\pm_1} < m_{\tilde{\chi}^\pm_2}$) and neutralinos $\tilde{\chi}^0_k$ ($k = 1, ..., 4$, $m_{\tilde{\chi}^0_1} < ... < m_{\tilde{\chi}^0_4}$) are determined by the parameters M_2, M_1, μ 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]:
$3.03 \times 10^{-4} < B(b \to s\ \gamma) < 4.01 \times 10^{-4}$ (95% CL) [12], $0.60 \times 10^{-6} < B(b \to s\ l^+ l^-) < 2.60 \times 10^{-6}$ with $l = e$ or μ (95% CL) [13], $B(B_s \to \mu^+\mu^-) < 4.8 \times 10^{-8}$ (90% CL) [12], $|R^{SUSY}_{B\tau\nu} - 1.77| < 1.27$ (95% CL) with $R^{SUSY}_{B\tau\nu} \equiv B^{SUSY}(B^-_u \to \tau^-\bar{\nu}_\tau)/B^{SM}(B^-_u \to \tau^-\bar{\nu}_\tau) \simeq (1 - (\frac{m_{B^+} \tan\beta}{m_{H^+}})^2)^2$ [14]. Moreover we impose the following condition on the SUSY prediction: $|\Delta M^{SUSY}_{B_s} - 17.77| < ((0.12 \times 1.96)^2 + 3.3^2)^{1/2}\ ps^{-1} = 3.31\ ps^{-1}$ (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 ρ parameter [17]: $\Delta\rho(SUSY) < 0.0012$.

(iii) The LEP limits on the SUSY particle masses [18]: $m_{\tilde{\chi}^\pm_1} > 103$ GeV, $m_{\tilde{\chi}^0_1} > 50$ GeV, $m_{\tilde{u}_1,\tilde{d}_1} > 100$ GeV, $m_{\tilde{u}_1,\tilde{d}_1} > m_{\tilde{\chi}^0_1}$, $m_{A^0} > 93$ GeV, $m_{h^0} > 110$ 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]: $m_{\tilde{g}} > 308$ GeV.

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

$$|A_{U\alpha\alpha}|^2 < 3\ Y^2_{U\alpha}\ (M^2_{Q_u\alpha\alpha} + M^2_{U\alpha\alpha} + m^2_2)\ , \tag{9}$$

$$|A_{D\alpha\alpha}|^2 < 3\ Y^2_{D\alpha}\ (M^2_{Q\alpha\alpha} + M^2_{D\alpha\alpha} + m^2_1)\ , \tag{10}$$

$$|A_{U\alpha\beta}|^2 < Y^2_{U\gamma}\ (M^2_{Q_u\alpha\alpha} + M^2_{U\beta\beta} + m^2_2)\ , \tag{11}$$

$$|A_{D\alpha\beta}|^2 < Y^2_{D\gamma}\ (M^2_{Q\alpha\alpha} + M^2_{D\beta\beta} + m^2_1)\ , \tag{12}$$

[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 $B(B_d \to \mu^+\mu^-)$, $B(b \to d\ l^+l^-)$, ΔM_{B_d} and ΔM_{D^0} as they practically do not constrain the 2nd and 3rd generation squark mixing which we are interested in here.

with ($\alpha \neq \beta; \gamma = \mathrm{Max}(\alpha, \beta); \alpha, \beta = 1, 2, 3$) and $m_1^2 = (m_{H^\pm}^2 + m_Z^2 \sin^2\theta_W) \sin^2\beta - \frac{1}{2}m_Z^2$, $m_2^2 = (m_{H^\pm}^2 + m_Z^2 \sin^2\theta_W) \cos^2\beta - \frac{1}{2}m_Z^2$. The Yukawa couplings of the up-type and down-type quarks are $Y_{U\alpha} = \sqrt{2}m_{u_\alpha}/v_2 = \frac{g}{\sqrt{2}}\frac{m_{u_\alpha}}{m_W \sin\beta}$ ($u_\alpha = u, c, t$) and $Y_{D\alpha} = \sqrt{2}m_{d_\alpha}/v_1 = \frac{g}{\sqrt{2}}\frac{m_{d_\alpha}}{m_W \cos\beta}$ ($d_\alpha = d, s, b$), with m_{u_α} and m_{d_α} 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 $m_W = 80.4$ GeV, $m_Z = 91.2$ GeV and the on-shell top-quark mass $m_t = 174.3$ GeV. We have found that our results shown in the following are fairly insensitive to m_t.

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_{Q23}^2, M_{D23}^2, A_{U23}, A_{D23} and A_{D32}; the constraints from $B(b \to s\gamma)$ and ΔM_{B_s} are especially important [22].

4 Quark flavour violating gluino decays

We study the effect of QFV due to the 2nd and 3rd generation squark mixing on the decays of gluinos which could be copiously produced at LHC. We focus on the QFV gluino decays

$$\tilde{g} \to \tilde{u}_i\, c \to c\, t\, \tilde{\chi}_1^0 \quad \text{and} \quad \tilde{g} \to \tilde{u}_i\, t \to c\, t\, \tilde{\chi}_1^0\,, \tag{13}$$

leading to the same final state $c\, t\, \tilde{\chi}_1^0$. We calculate the gluino and squark decay widths taking into account the following two–body decays:

$$\begin{aligned} \tilde{g} &\to \tilde{u}_i\, u_k,\ \tilde{d}_i\, d_k, \\ \tilde{u}_i &\to u_k\, \tilde{\chi}_n^0,\ d_k\, \tilde{\chi}_m^+,\ \tilde{d}_j\, W^+,\ \tilde{u}_j\, Z^0,\ \tilde{u}_j\, h^0, \end{aligned} \tag{14}$$

where $u_k = (u, c, t)$ and $d_k = (d, s, b)$. Note that the squark decays into the heavier Higgs bosons are kinematically forbidden in our scenarios studied below. The formulae for the two–body decays in (14) 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.

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

M_1	M_2	$m_{\tilde{g}}$	μ	$\tan\beta$	m_{A^0}
139	264	800	1000	10	800

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

$M^2_{U\alpha\beta}$	$\beta=1$	$\beta=2$	$\beta=3$
$\alpha=1$	$(820)^2$	0	0
$\alpha=2$	0	$(600)^2$	$(224)^2$
$\alpha=3$	0	$(224)^2$	$(580)^2$

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.

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

$\tilde{\chi}^0_1$	$\tilde{\chi}^0_2$	$\tilde{\chi}^0_3$	$\tilde{\chi}^0_4$	$\tilde{\chi}^\pm_1$	$\tilde{\chi}^\pm_2$
138	261	1003	1007	261	1007

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

We take $\tan\beta, m_{A^0}, M_1, M_2, m_{\tilde{g}}, \mu, M^2_{Q\alpha\beta}, M^2_{U\alpha\beta}, M^2_{D\alpha\beta}, A_{U\alpha\beta}$ 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 $M^2_{Q\alpha\beta}$, $M^2_{U\alpha\beta}$, $M^2_{D\alpha\beta}$, $A_{U\alpha\beta}$ and $A_{D\alpha\beta}$ with $\alpha \neq \beta$. Note that the so-called minimal flavour violation (MFV) corresponds to the case where all of these squark generation mixing terms are zero and the CKM mixing matrix is the only source of flavour violation (QFV). 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 $B(b \to s\gamma) = 3.57 \times 10^{-4}$, $B(b \to sl^+l^-) = 1.59 \times 10^{-6}$, $B(b \to s\nu\bar{\nu}) = 4.07 \times 10^{-5}$, $B(B_s \to \mu^+\mu^-) = 4.72 \times 10^{-9}$, $B(B_u^+ \to \tau^+\nu) = 7.85 \times 10^{-5}$, $\Delta M_{B_s} = 17.38\ ps^{-1}$ 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.

$R^{\tilde{u}}_{i\alpha}$	$\tilde{u}_L$	$\tilde{c}_L$	$\tilde{t}_L$	$\tilde{u}_R$	$\tilde{c}_R$	$\tilde{t}_R$
$\tilde{u}_1$	-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.

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 $B(\tilde{u}_1 \to c\tilde{\chi}^0_1) = 0.576$, $B(\tilde{u}_1 \to t\tilde{\chi}^0_1) = 0.401$, $B(\tilde{u}_2 \to c\tilde{\chi}^0_1) = 0.495$, $B(\tilde{u}_2 \to t\tilde{\chi}^0_1) = 0.469$. This leads to the following gluino decay branching ratios:

$$B(\tilde{g} \to ct\tilde{\chi}^0_1) = \sum_{i=1,2} \Big[B(\tilde{g} \to \tilde{u}_i c) B(\tilde{u}_i \to t\tilde{\chi}^0_1) + B(\tilde{g} \to \tilde{u}_i t) B(\tilde{u}_i \to c\tilde{\chi}^0_1) \Big] = 0.463, \quad (15)$$

$$B(\tilde{g} \to cc\tilde{\chi}^0_1) = \sum_{i=1,2} \Big[B(\tilde{g} \to \tilde{u}_i c) B(\tilde{u}_i \to c\tilde{\chi}^0_1) \Big] = 0.380, \quad (16)$$

$$B(\tilde{g} \to tt\tilde{\chi}^0_1) = \sum_{i=1,2} \Big[B(\tilde{g} \to \tilde{u}_i t) B(\tilde{u}_i \to t\tilde{\chi}^0_1) \Big] = 0.120. \quad (17)$$

Note that the QFV gluino decay branching ratio of Eq.(15) 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 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 $\tilde{\chi}^0_1(\simeq \tilde{B}^0)$ and practically do not couple to $\tilde{\chi}^0_2(\simeq \tilde{W}^0)$, $\tilde{\chi}^{\pm}_1(\simeq \tilde{W}^{\pm})$, and that $\tilde{\chi}^0_{3,4}$, $\tilde{\chi}^{\pm}_2$ 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.1 we show contours of $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ 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_{U22} and M^2_{U23} are fixed as in our reference scenario defined in Table 1. We see that the QFV decay branching ratio $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ quickly increases up to $\sim 50\%$ with incease of the effective $\tilde{c}_R - \tilde{t}_R$ mixing angle $\tan(2\theta^{eff}_{23}) \equiv 2M^2_{U23}/\Delta M^2_U$.

In Fig.2 we present contours of $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ 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. 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 $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ increases quickly with increase of the $\tilde{c}_R - \tilde{t}_R$ mixing parameter $|\delta^{uRR}_{23}|$ and can be very large in a significant part of the $\delta^{uLL}_{23} - \delta^{uRR}_{23}$ plane allowed by all of the conditions (i)-(v) including the $b \to s\gamma$ constraint.

Studying the branching ratios of the gluino and up-type squark two-body decays separately would allow for a better understanding of their contributions to the QFV gluino decay $\tilde{g} \to ct\tilde{\chi}^0_1$. In Fig.3 we show the δ^{uRR}_{23} dependences of the gluino and squark decay branching ratios, where all basic parameters other than M^2_{U23} are fixed as in the scenario specified in Table 1. We see that the QFV decay branching ratio $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ increases quickly with increase of $|\delta^{uRR}_{23}|$ for $|\delta^{uRR}_{23}| \lesssim 0.1$ and can be very large ($\sim 50\%$) in a wide range of δ^{uRR}_{23}. This behaviour can be explained as follows: The gluino decays into squarks other than $\tilde{u}_{1,2}$ are kinematically forbidden, and $\tilde{u}_1$ and $\tilde{u}_2$ become quickly a strong mixture of the flavour eigenstates $\tilde{c}_R$ and $\tilde{t}_R$ with increase of the $\tilde{c}_R$ - $\tilde{t}_R$ mixing term M^2_{U23} because of the small mass parameter difference $(M^2_{\tilde{u}RR})_{22} - (M^2_{\tilde{u}RR})_{33} = (599)^2 - (604)^2$ GeV^2 in this scenario (see Eq.(4)). This results in the quick increase of $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ with increase of $|\delta^{uRR}_{23}|$ in $|\delta^{uRR}_{23}| \lesssim 0.1$ and the very large $B(\tilde{g} \to ct\tilde{\chi}^0_1)$ in $0.1 \lesssim |\delta^{uRR}_{23}| \lesssim 1.0$ (see discussion just below Eq.(15)). Note that $\tilde{u}_1 = \tilde{c}_R$ and $\tilde{u}_2 = \tilde{t}_R$ for $\delta^{uRR}_{23} = 0$, which explains the behaviour of the gluino and squark two-body decay branching ratios around $\delta^{uRR}_{23} = 0$ in Fig.3. Notice also that $m_{\tilde{u}_1}$ ($m_{\tilde{u}_2}$) decreases (increases) with the increase of $|\delta^{uRR}_{23}|$, which explains the behaviour (including the various kinematical thresholds) of the gluino and squark two-body decay branching ratios with increasing $|\delta^{uRR}_{23}|$ for $|\delta^{uRR}_{23}| \gtrsim 0.1$. Moreover, as $\tilde{u}_2$ equals the flavour eigenstate $\tilde{u}_R$ for $|\delta^{uRR}_{23}| \gtrsim 0.9$, the branching ratios $B(\tilde{u}_2 \to c\tilde{\chi}^0_1)$ and $B(\tilde{u}_2 \to t\tilde{\chi}^0_1)$ vanish in this range.

In Fig.4 we show the δ_{23}^{uRL} dependences of the gluino decay branching ratios, where all basic parameters other than A_{U32} are fixed as in the scenario specified in 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 δ_{23}^{uRL}. We find that the QFV decay branching ratio $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ decreases (down to $\sim$ 30%) with increase of $|\delta_{23}^{uRL}|$ and the quark-generation violating (QGV) decay branching ratio $B(\tilde{g} \to cb\tilde{\chi}_1^\pm)$ increases (up to $\sim$ 20%) with the increase of $|\delta_{23}^{uRL}|$. This behaviour can be explained as follows: In this scenario the $\tilde{\chi}_1^\pm$ ($\simeq \tilde{W}^\pm$ (wino)) couples to $\tilde{q}_L$ and its coupling to $\tilde{q}_R$ is suppressed. On the other hand, $\tilde{\chi}_1^0$ ($\simeq \tilde{B}^0$ (bino)) couples much more strongly to $\tilde{c}_R$ and $\tilde{t}_R$ than to $\tilde{c}_L$ and $\tilde{t}_L$. Sizable δ_{23}^{uRL} (i.e. $\tilde{c}_R - \tilde{t}_L$ mixing parameter) induces a sizable $\tilde{t}_L$ component in $\tilde{u}_{1,2}(\sim \tilde{c}_R + \tilde{t}_R)$, which enhances the widths $\Gamma(\tilde{u}_{1,2} \to b\tilde{\chi}_1^\pm)$ and leads to a suppression of $B(\tilde{u}_{1,2} \to c\tilde{\chi}_1^0)$ and $B(\tilde{u}_{1,2} \to t\tilde{\chi}_1^0)$. As a result $B(\tilde{g} \to cb\tilde{\chi}_1^\pm) = \sum_{i=1,2} B(\tilde{g} \to \tilde{u}_i c)B(\tilde{u}_i \to b\tilde{\chi}_1^\pm)$ [5] is enhanced for sizable δ_{23}^{uRL} while $B(\tilde{g} \to ct\tilde{\chi}_1^0) = \sum_{i=1,2} [B(\tilde{g} \to \tilde{u}_i c)B(\tilde{u}_i \to t\tilde{\chi}_1^0) + B(\tilde{g} \to \tilde{u}_i t)B(\tilde{u}_i \to c\tilde{\chi}_1^0)]$ is suppressed.

As for the δ_{32}^{uRL} dependence plot of the gluino decay branching ratios, where all basic parameters other than A_{U23} are fixed as in the scenario specified in Table 1, we have obtained similar results (including the allowed range) to those for the δ_{23}^{uRL} dependence in Fig.4. We have found 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 allowed range $|\delta_{32}^{uRL}| \lesssim 0.3$, and that $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ decreases (down to $\sim$ 30%) with the increase of $|\delta_{32}^{uRL}|$ 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. This behaviour can be explained as in the case of the δ_{23}^{uRL} dependence: Sizable δ_{32}^{uRL} (i.e. $\tilde{c}_L - \tilde{t}_R$ mixing parameter) induces a sizable $\tilde{c}_L$ component in $\tilde{u}_{1,2}(\sim \tilde{c}_R + \tilde{t}_R)$, which enhances the widths $\Gamma(\tilde{u}_{1,2} \to s\tilde{\chi}_1^\pm)$ and suppresses $B(\tilde{u}_{1,2} \to c\tilde{\chi}_1^0)$ and $B(\tilde{u}_{1,2} \to t\tilde{\chi}_1^0)$. This means that $B(\tilde{g} \to st\tilde{\chi}_1^\pm) = \sum_{i=1,2} B(\tilde{g} \to \tilde{u}_i t)B(\tilde{u}_i \to s\tilde{\chi}_1^\pm)$ is enhanced for sizable δ_{32}^{uRL} whereas $B(\tilde{g} \to ct\tilde{\chi}_1^0) = \sum_{i=1,2} [B(\tilde{g} \to \tilde{u}_i c)B(\tilde{u}_i \to t\tilde{\chi}_1^0) + B(\tilde{g} \to \tilde{u}_i t)B(\tilde{u}_i \to c\tilde{\chi}_1^0)]$ is suppressed.

As for the δ_{23}^{uLL} dependence plot of the gluino decay branching ratios, where all basic parameters other than M_{Q23}^2 are fixed as in the scenario specified in Table 1, we have found that the QFV decay branching ratio $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ is insensitive to δ_{23}^{uLL} and can be quite large ($\sim$ 50%) in a sizable allowed range $0.03 \lesssim \delta_{23}^{uLL} \lesssim 0.12$ as can be seen in

[5] Note that gluino decays into a down-type squark, such as $B(\tilde{g} \to \tilde{d}_i b)$, are kinematically forbidden in this scenario and hence that such decays can not contribute to $B(\tilde{g} \to cb\tilde{\chi}_1^\pm)$.

Fig.2, where our reference scenario corresponds to $(\delta_{23}^{uLL}, \delta_{23}^{uRR}) = (0.068, 0.144)$. The reason for the insensitivity of $B(\tilde{g} \to ct\tilde{\chi}_1^0)$ to δ_{23}^{uLL} is as follows: the $\tilde{c}_L - \tilde{t}_L$ mixing parameter δ_{23}^{uLL} affects mainly the masses and mixing of the heavier up-type squarks $\tilde{u}_4$ and $\tilde{u}_5$, but mainly the on-shell $\tilde{u}_{1,2}$ mediate the QFV decay $\tilde{g} \to ct\tilde{\chi}_1^0$ in this scenario.

5 Invariant mass distributions

Here we study the invariant mass distributions (i.e. the differential decay branching ratios) $\mathrm{dBr}(\tilde{g} \to \tilde{u}_i u_j \to u_j u_k \tilde{\chi}_n^0)/\mathrm{d}M_{u_j u_k}$, 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 endpoinds of the distributions are given in terms of the masses of the involved particles by [24]

$$M_{u_j u_k}^{i(\min,\max)} = \Big\{ m_{u_j}^2 + m_{u_k}^2 + \frac{1}{2m_{\tilde{u}_i}^2} \Big[(m_{\tilde{g}}^2 - m_{u_j}^2 - m_{\tilde{u}_i}^2)(m_{\tilde{u}_i}^2 + m_{u_k}^2 - m_{\tilde{\chi}_n^0}^2) \mp \lambda^{\frac{1}{2}}(m_{\tilde{g}}^2, m_{u_j}^2, m_{\tilde{u}_i}^2)\ \lambda^{\frac{1}{2}}(m_{\tilde{u}_i}^2, m_{u_k}^2, m_{\tilde{\chi}_n^0}^2) \Big] \Big\}^{\frac{1}{2}}, \qquad (18)$$

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 $M_{u_j u_k}^{i(\min,\max)} \neq M_{u_k u_j}^{i(\min,\max)}$ 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:

$$\mathrm{dBr}(\tilde{g} \to u_j u_k \tilde{\chi}_n^0)/\mathrm{d}M_{u_j u_k} = \frac{1}{1+\delta_{jk}} \sum_i \Big[\mathrm{dBr}(\tilde{g} \to \tilde{u}_i u_j \to u_j u_k \tilde{\chi}_n^0)/\mathrm{d}M_{u_j u_k} + \mathrm{dBr}(\tilde{g} \to \tilde{u}_i u_k \to u_k u_j \tilde{\chi}_n^0)/\mathrm{d}M_{u_j u_k} \Big]. \qquad (19)$$

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

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.5 shows the invariant mass distributions of the top quark pairs for the QFV scenario, where one has B($\tilde{g} \to t\bar{t}\tilde{\chi}_1^0$) = 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($\tilde{g} \to ct\tilde{\chi}_1^0$) = 46.3%. Fig.5 shows the invariant mass distribution of ct. There are more edge structures due to the processes $\tilde{g} \to \tilde{u}_1 t \to tc\tilde{\chi}_1^0$ [with $M_{tc}^{1(\mathrm{min,max})} = (253, 526)$ GeV], $\tilde{g} \to \tilde{u}_1 c \to ct\tilde{\chi}_1^0$ [with $M_{ct}^{1(\mathrm{min,max})} = (254, 580)$ GeV], and $\tilde{g} \to \tilde{u}_2 c \to ct\tilde{\chi}_1^0$ [with $M_{ct}^{2(\mathrm{min,max})} = (219, 497)$ GeV]. Note that $\tilde{g} \to \tilde{u}_2 t$ 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.6 we show the two invariant mass distributions of tt and ct, where one has B($\tilde{g} \to tt\tilde{\chi}_1^0$) = 16.6%, and B($\tilde{g} \to ct\tilde{\chi}_1^0$) = 31.4%. Note that the QFV decay branching ratio B($\tilde{g} \to ct\tilde{\chi}_1^0$) 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 $\tilde{g} \to \tilde{u}_2 t$ is kinematically allowed but phase-space suppressed. Moreover, $\tilde{u}_2 \to t\tilde{\chi}_1^0$ is strongly suppressed because $\tilde{u}_2(\sim \tilde{t}_L + \tilde{c}_L)$ does not significantly couple to $\tilde{\chi}_1^0(\sim \tilde{B}^0$(Bino)) in this scenario. Hence, B($\tilde{g} \to \tilde{u}_2 t \to tt\tilde{\chi}_1^0$)(=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 $\tilde{g} \to \tilde{u}_1 t \to tc\tilde{\chi}_1^0$ [with $M_{tc}^{1(\mathrm{min,max})} = (601, 971)$ GeV], and $\tilde{g} \to \tilde{u}_1 c \to ct\tilde{\chi}_1^0$ [with $M_{ct}^{1(\mathrm{min,max})} = (183, 1022)$ GeV]. The decays $\tilde{g} \to \tilde{u}_2\ c/t$ are phase-space suppressed and the decays $\tilde{u}_2 \to\ c/t\ \tilde{\chi}_1^0$ are strongly suppressed in this scenario as is explained above. Hence, B($\tilde{g} \to \tilde{u}_2\ c/t \to ct\tilde{\chi}_1^0$)(=0.0004) is very small.

$M^2_{Q\alpha\beta}$	$\beta = 1$	$\beta = 2$	$\beta = 3$
$\alpha = 1$	$(1200)^2$	0	0
$\alpha = 2$	0	$(1200)^2$	$(500)^2$
$\alpha = 3$	0	$(500)^2$	$(1128)^2$

M_1	M_2	$m_{\tilde{g}}$	μ	$\tan\beta$	m_{A^0}
255	497	1300	756	5	800

$M^2_{D\alpha\beta}$	$\beta = 1$	$\beta = 2$	$\beta = 3$
$\alpha = 1$	$(1141)^2$	0	0
$\alpha = 2$	0	$(1141)^2$	0
$\alpha = 3$	0	0	$(1100)^2$

$M^2_{U\alpha\beta}$	$\beta = 1$	$\beta = 2$	$\beta = 3$
$\alpha = 1$	$(1149)^2$	0	0
$\alpha = 2$	0	$(1149)^2$	$(894)^2$
$\alpha = 3$	0	$(894)^2$	$(877)^2$

Table 4: The MSSM parameters in the QFV scenario with $m_{\tilde{g}} = 1300$ GeV. All of $A_{U\alpha\beta}$ and $A_{D\alpha\beta}$ are set to zero. All mass parameters are given in GeV.

$\tilde{u}_1$	$\tilde{u}_2$	$\tilde{u}_3$	$\tilde{u}_4$	$\tilde{u}_5$	$\tilde{u}_6$
466	1054	1149	1199	1275	1379

$\tilde{d}_1$	$\tilde{d}_2$	$\tilde{d}_3$	$\tilde{d}_4$	$\tilde{d}_5$	$\tilde{d}_6$
1046	1101	1141	1141	1201	1274

$\tilde{\chi}^0_1$	$\tilde{\chi}^0_2$	$\tilde{\chi}^0_3$	$\tilde{\chi}^0_4$	$\tilde{\chi}^\pm_1$	$\tilde{\chi}^\pm_2$
253	483	758	775	482	774

Table 5: Sparticles and corresponding masses (in GeV) in the scenario of Table 4.

The signature of the QFV decay $\tilde{g} \to$ c t $\tilde{\chi}^0_1$ at LHC would be 'charm-jet + top-quark + missing-energy'. Therefore charm-tagging would be very useful. Even if charm-tagging is not feasible, we could detect the signature of the QFV gluino decay since it yields a remarkable (detectable) signature of $\tilde{g} \to$ q t $\tilde{\chi}^0_1$ (q $\neq$ t). We have shown that QFV gluino decay branching ratios such as B($\tilde{g} \to$ c t $\tilde{\chi}^0_1$) can be very large despite the very strong experimental constraints from QFV processes. This shows that the QFV gluino decays can contribute significantly to signal event rates at LHC. Therefore 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 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.

$R^{\tilde{u}}_{i\alpha}$	$\tilde{u}_L$	$\tilde{c}_L$	$\tilde{t}_L$	$\tilde{u}_R$	$\tilde{c}_R$	$\tilde{t}_R$
$\tilde{u}_1$	-0.001	0.006	-0.021	0	0.587	-0.809
$\tilde{u}_2$	-0.137	0.621	-0.771	0	-0.024	0.006
$\tilde{u}_3$	0	0	0	-1.0	0	0
$\tilde{u}_4$	-0.976	-0.219	-0.003	0	0	0
$\tilde{u}_5$	0.171	-0.752	-0.636	0	-0.032	-0.012
$\tilde{u}_6$	0.003	-0.015	-0.033	0	0.808	0.588

Table 6: The up-type squark compositions in the flavour eigenstates, i.e. the mixing matrix $R^{\tilde{u}}_{i\alpha}$ for the scenario of Table 4.

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 $\tilde{\chi}^0_1$) 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 ΔM_{B_s} 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.

Acknowledgments

This work is supported by the "Fonds zur Förderung der wissenschaftlichen Forschung (FWF) of Austria, project No. P18959-N16. The authors acknowledge support from EU under the MRTN-CT-2006-035505 and MRTN-CT-2006-503369 network programs.

A. B. was supported by the Spanish grants SAB 2006-0072, FPA 2005-01269, FPA 2005-25348-E and FPA 2008-00319/FPA of the Ministero de Educacion y Ciencia. T. K. is supported by the Portuguese FCT through the projects POCI/FP/81919/2007 and CFTP-FCT UNIT 777, which are partially funded through POCTI (FEDER). W. P. is partially supported by the German Ministry of Education and Research (BMBF) project nr. 05HT6WWA.

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/0202259].

[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. D**72** (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. D**74** (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].

[25] M. Battaglia *et al.*, Eur. Phys. J. C **22** (2001) 535 [arXiv:hep-ph/0106204].

Figure Captions

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

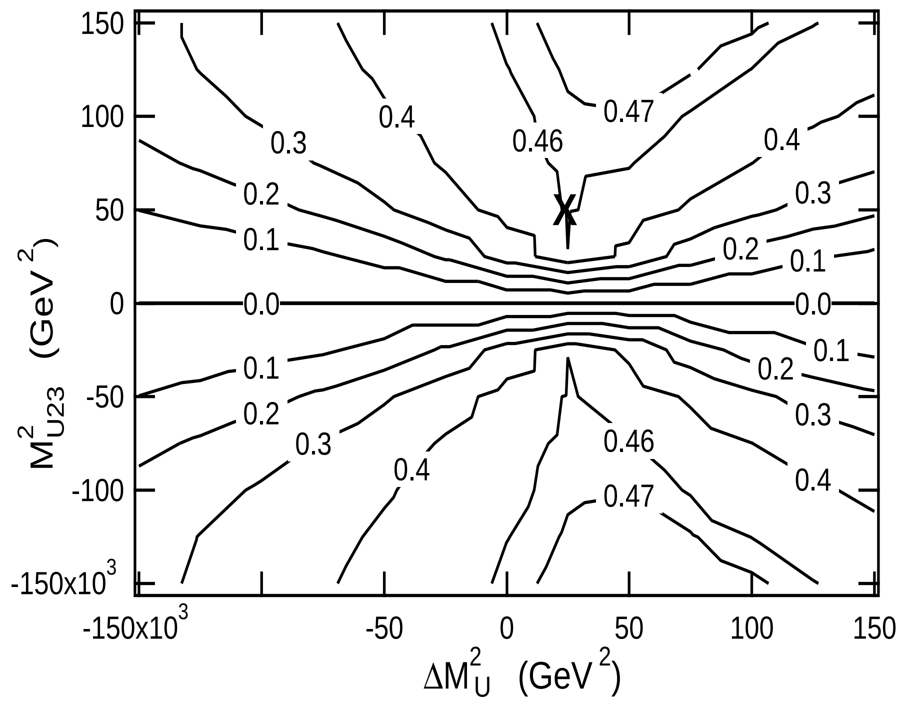
Figure 2: Contours of the QFV decay branching ratio $B(\tilde{g} \to c t \tilde{\chi}_1^0)$ (solid lines) in the $\delta_{23}^{uLL} - \delta_{23}^{uRR}$ plane where all of the conditions (i)-(v) except the $b \to s\gamma$ constraint are satisfied. Contours of $10^4 \times B(b \to s\gamma)$ (dashed lines) are also shown. The condition (i) requires $3.03 < 10^4 \times B(b \to s\,\gamma) < 4.01$. The point "x" of $(\delta_{23}^{uLL}, \delta_{23}^{uRR}) = (0.068, 0.144)$ corresponds to our reference scenario of Table 1.

Figure 3: δ_{23}^{uRR} 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 $\delta_{23}^{uRR} = 0.144$ corresponds to our reference scenario of Table 1. The shown range of δ_{23}^{uRR} is the whole range allowed by the conditions (i) to (v) given in the text; note that the range $|\delta_{23}^{uRR}| \gtrsim 1.0$ is excluded by the condition $m_{\tilde{u}_1} > m_{\tilde{\chi}_1^0}$ in (iii).

Figure 4: δ_{23}^{uRL} dependences of the branching ratios of the gluino cascade decays. The point "x" of $\delta_{23}^{uRL} = 0$ corresponds to our reference scenario of Table 1. The shown range of δ_{23}^{uRL} is the whole range allowed by the conditions (i) to (v) given in the text; note that the range $|\delta_{23}^{uRL}| \gtrsim 0.3$ is excluded by the condition (v).

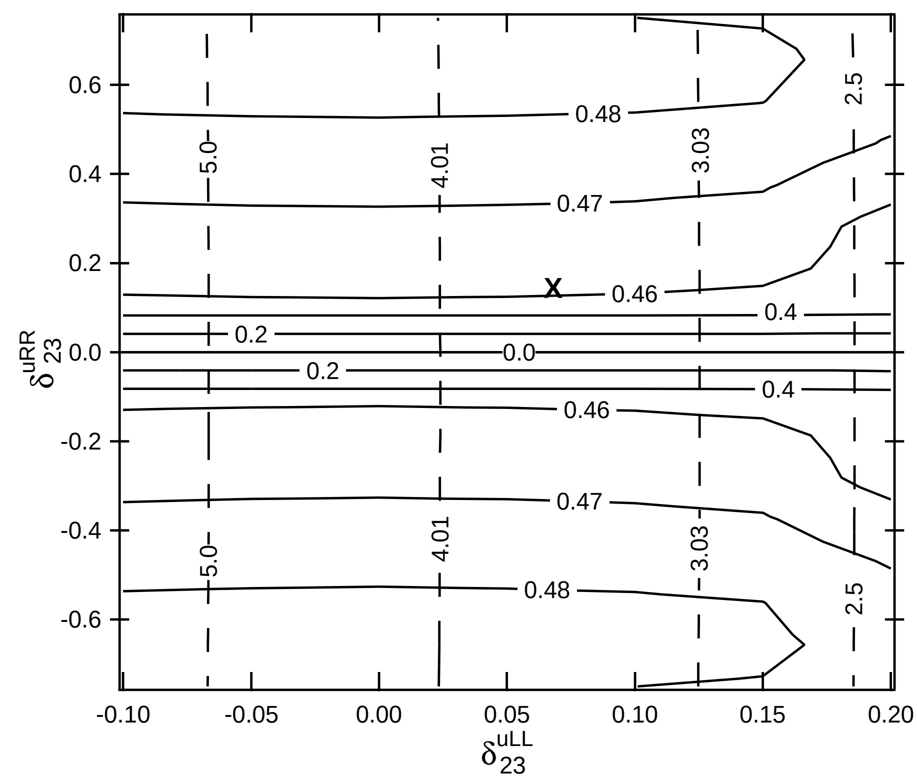
Figure 5: Invariant mass distributions of two up-type quarks from the decay $\tilde{g} \to u_j u_k \tilde{\chi}_1^0$ for the QFV scenario of Table 1.

Figure 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 4.



150
100
50
0
-50
-100
-150x10^3
-150x10^3
-50
0
50
100
150
M^2_{U23} (GeV^2)
ΔM^2_U (GeV^2)
0.3
0.4
0.46
0.47
0.4
0.3
0.2
0.1
0.2
0.1
0.0
0.0
0.1
0.2
0.1
0.2
0.3
0.3
0.4
0.46
0.47
0.4
X

Fig.1



0.6
0.4
0.2
0.0
-0.2
-0.4
-0.6
-0.10
-0.05
0.00
0.05
0.10
0.15
0.20
δ^{uRR}_{23}
δ^{uLL}_{23}
0.48
0.47
0.46
0.4
0.2
0.0
0.2
0.4
0.46
0.47
0.48
5.0
4.01
3.03
2.5
X

Fig.2

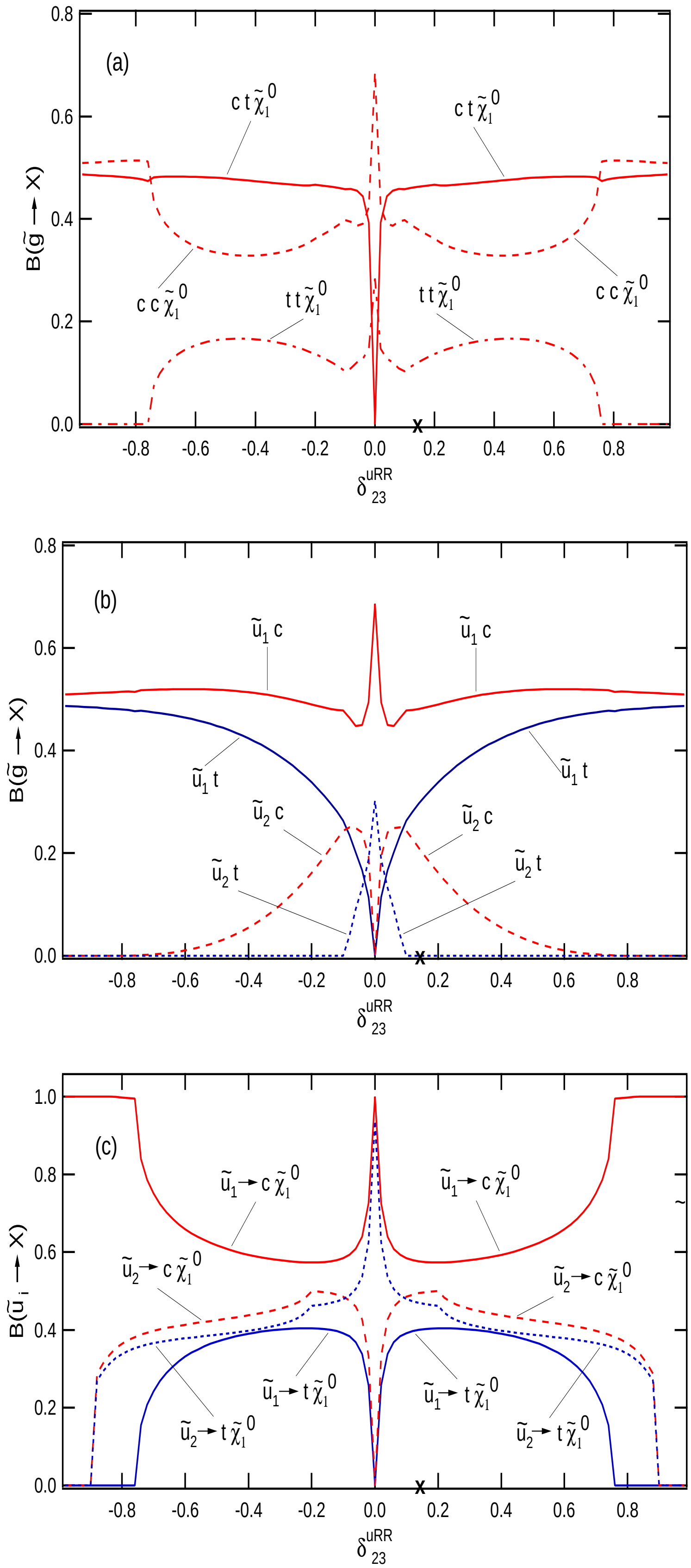


Fig.3

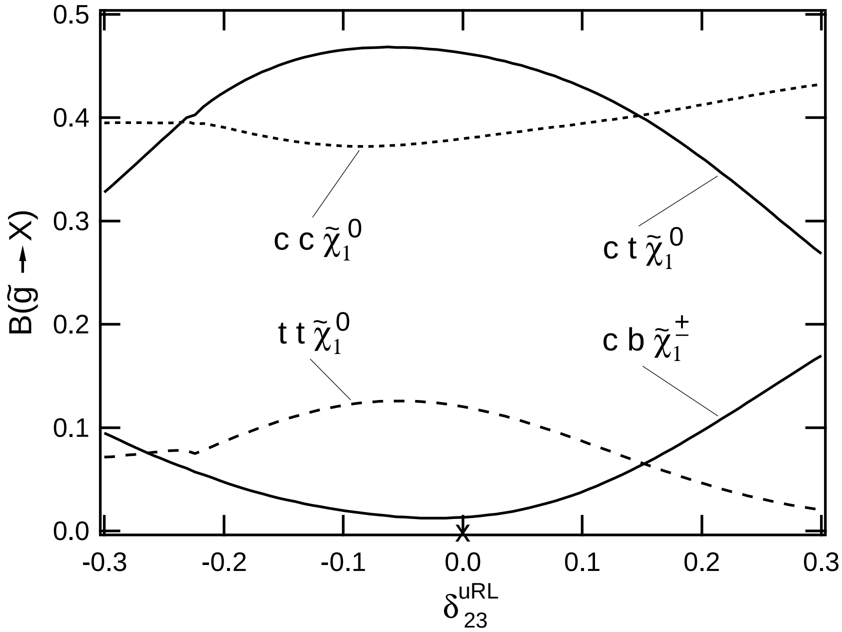
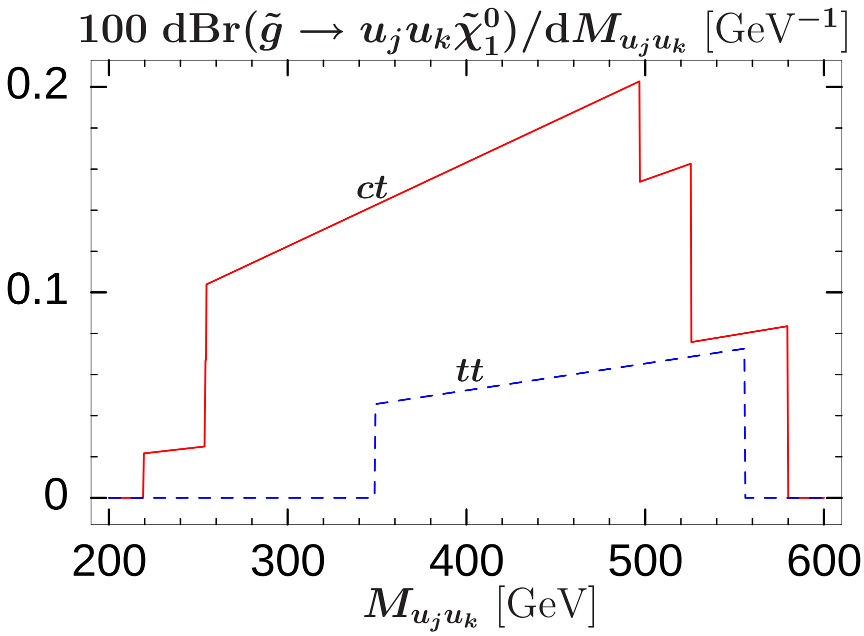
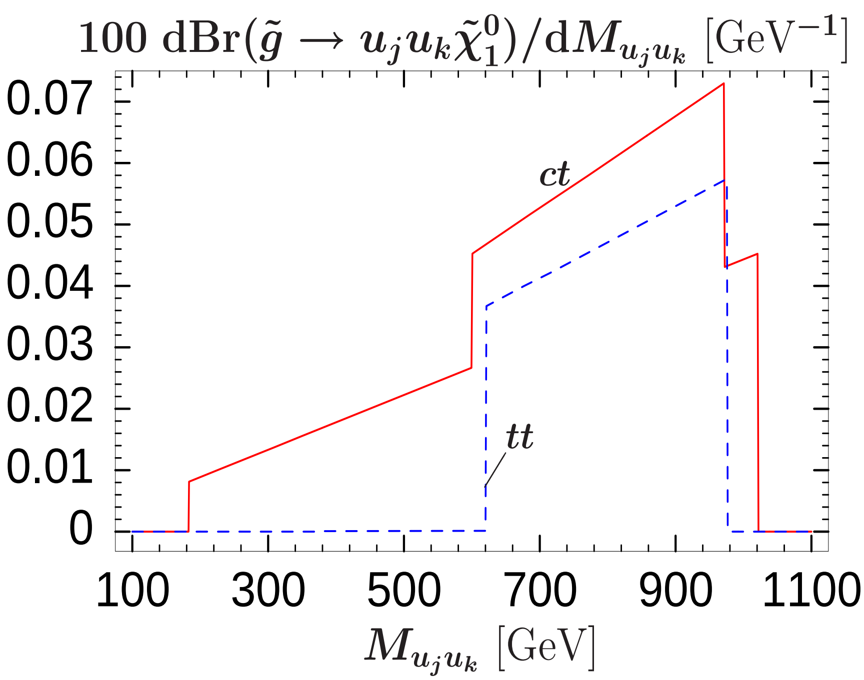

0.5
0.4
0.3
0.2
0.1
0.0
B(g̃ → X)
-0.3
-0.2
-0.1
0.0
0.1
0.2
0.3
δ_{23}^{uRL}
$c\,c\,\tilde{\chi}_1^0$
$c\,t\,\tilde{\chi}_1^0$
$t\,t\,\tilde{\chi}_1^0$
$c\,b\,\tilde{\chi}_1^{\pm}$

Fig.4



100 dBr($\tilde{g} \to u_j u_k \tilde{\chi}_1^0$)/d$M_{u_j u_k}$ [GeV^{-1}]
0.2
0.1
0
ct
tt
200
300
400
500
600
$M_{u_j u_k}$ [GeV]

Fig.5



100 dBr($\tilde{g} \to u_j u_k \tilde{\chi}_1^0$)/d$M_{u_j u_k}$ [GeV^{-1}]
0.07
0.06
0.05
0.04
0.03
0.02
0.01
0
ct
tt
100
300
500
700
900
1100
$M_{u_j u_k}$ [GeV]

Fig.6