SUPERSYMMETRIC QCD FLAVOUR CHANGING TOP QUARK DECAY

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Abstract

We present a detailed and complete calculation of the gluino and scalar quarks contribution to the flavour-changing top quark decay into a charm quark and a photon, gluon, or a $Z^0$ boson within the minimal supersymmetric standard model including flavour changing gluino-quarks-scalar quarks couplings in the right-handed sector. We compare the results with the ones presented in an earlier paper where we considered flavour changing couplings only in the left-handed sector. We show that these new couplings have important consequences leading to a large enhancement when the mixing of the scalar partners of the left- and right-handed top quark is included. Furthermore CP violation in the flavour changing top quark decay will occur when a SUSY phase is taken into account.
1 Introduction

Flavour changing top quark decay modes are a promising test ground for models beyond the standard model (SM). While in the SM the branching ratios of the decays $t \rightarrow c\gamma$, $cg$ and $cZ$ are far away from experimental reach [1]-[5], the authors of [5]-[6] showed that they are enhanced by several (3-4) orders of magnitude in Two-Higgs-doublet models (THDM's).

Nowadays, CDF [7]-[8] and D∅ [9] have begun to explore flavour changing top quark decays and interesting bounds have been reported [8]. A systematic examination of anomalous top quark quark interactions is actively pursued [10]-[15].

Within supersymmetry, the decays $t \rightarrow cV$ were first considered in [16] and the authors obtained the same enhancement as in the THDM’s. However as we have pointed out in a recent paper [17], in their calculation of the QCD corrections they had an inconsistency basically due to the lack of gauge invariance arising from the omission of the gluino-gluino-gluon coupling. They also did not include the non-negligible mixing of the scalar partners of the left and right handed quarks. In a very recent paper [18] the calculations were redone for the weak sector with charginos and neutralinos within the relevant loops including the mixing of the scalar quarks, where it was shown that supersymmetric contributions to $t \rightarrow cV$ can be up to 5 orders of magnitude larger than their SM counterparts.

In our previous paper [17] the results of the gluino and scalar quarks contribution to the flavour-changing top quark decay into a charm quark and a photon, a gluon or a $Z^0$ boson within the minimal supersymmetric standard model (MSSM) were presented. We included the mixing of the scalar partners of the left- and right-handed top quark and showed that it has several effects, the most important of which is to greatly enhance the $cZ$ decay mode for large values of the soft SUSY-breaking scalar mass $m_S$ and to give rise to a GIM-like suppression in the $c\gamma$ mode for certain combinations of parameters.

However the analysis of [17] considered flavour-changing strong interactions between the gluino, the quarks and their scalar partners only in the left-handed sector and kept the right-handed sector flavour-diagonal. This is a common assumption within the MSSM (see [17] and references therein) and might not be necessarily the case in any kind of extension of the MSSM, or more general assumptions within the MSSM.

The goal of this paper is to recalculate the flavour-changing top quark decays including flavour-changing couplings within the right-handed sector. We will assume maximal flavour-
changing in both sectors and analyse how the previous results will be changed. Furthermore
we show that flavour-changing couplings in the left- and right-handed sectors lead to a
CP violating term proportional to the gluino mass for the top quark decay modes under
consideration, which will be investigated in a further paper [19].

The Feynman diagrams and the couplings leading to the decay modes \( t \to c\gamma, \ cZ \) and \( cg \)
as well as the mass matrix of the scalar top quark are given in [17]. The only difference will
be the flavour-changing gluino-scalar quarks-quarks coupling in Eq. (6) of [17], which
will be taken in the present paper in the most general way:

\[
\mathcal{L}_{FC} = -\sqrt{2} g_s T^\alpha \bar{y}_a [ K^\alpha_L (c_\Theta \tilde{q}_1 - s_\Theta \tilde{q}_2) P_L - K^\alpha_R (s_\Theta \tilde{q}_1 + c_\Theta \tilde{q}_2) P_R ] q + h.c. \quad (1)
\]

Here \( K^\alpha_{L,R} \) is the supersymmetric version of the Kobayashi–Maskawa matrix and \( \Theta \) is the
mixing angle of the scalar partners of the left- and right-handed quarks. \( \tilde{q}_1 \) and \( \tilde{q}_2 \) are the
mass eigenstates which are related to the current eigenstates \( \tilde{q}_L \) and \( \tilde{q}_R \) by

\[
\tilde{q}_1 = \cos \Theta \tilde{q}_L + \sin \Theta \tilde{q}_R, \quad \tilde{q}_2 = -\sin \Theta \tilde{q}_L + \cos \Theta \tilde{q}_R \quad (2)
\]

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After summation over all diagrams, we obtain the following effective \( tcV \) vertex, neglecting
the charm quark mass:

\[
M_{\mu \nu} = -i \frac{\alpha_s}{2\pi} \frac{1}{p_\mu} \left[ \gamma_\mu \left( P_L V^\alpha_{\mu \nu} + P_R V^\alpha_{\mu \nu} \right) + \frac{P_\mu}{m_{top}} \left( P_L T^{\alpha}_{\mu \nu} + P_R T^{\alpha}_{\mu \nu} \right) \right] u_{p_\nu} \quad (3)
\]

\[
V^\alpha_{\gamma L} = ee_\gamma C_2(F) \left\{ K^{* \alpha}_{a2L} K^{\alpha}_{a3L} [ c_\Theta (C_{11}^{\alpha} - C_{11}^{SE}) + s_\Theta (C_{11}^{22\alpha} - C_{22}^{2\alpha}) ] + K^{* \alpha}_{a2L} K^{\alpha}_{a3R} c_\Theta s_\Theta \frac{m_{\tilde{g}}}{m_{top}} \left[ C_{11}^{\alpha}_{SEG} - C_{11}^{\alpha}_{SEG} | m_{top} = 0 - C_{22}^{2\alpha}_{SEG} + C_{22}^{2\alpha}_{SEG} | m_{top} = 0 \right] \right\}
\]

\[
V^\alpha_{\gamma R} = ee_\gamma C_2(F) \left\{ K^{* \alpha}_{a2L} K^{\alpha}_{a3L} [ c_\Theta (C_{11}^{\alpha} - C_{11}^{SE}) + s_\Theta (C_{11}^{22\alpha} - C_{22}^{2\alpha}) ] + K^{* \alpha}_{a2R} K^{\alpha}_{a3R} c_\Theta s_\Theta \frac{m_{\tilde{g}}}{m_{top}} \left[ C_{11}^{\alpha}_{SEG} - C_{11}^{\alpha}_{SEG} | m_{top} = 0 - C_{22}^{2\alpha}_{SEG} + C_{22}^{2\alpha}_{SEG} | m_{top} = 0 \right] \right\}
\]

\[
T^\alpha_{\gamma L} = ee_\gamma C_2(F) \left\{ K^{* \alpha}_{a2L} K^{\alpha}_{a3L} [ s_\Theta C_{11}^{\alpha} + c_\Theta C_{22}^{\alpha} ] - K^{* \alpha}_{a2L} K^{\alpha}_{a3L} c_\Theta s_\Theta \left[ C_{11}^{\alpha}_{\tilde{g}_{top}} - C_{22}^{\alpha}_{\tilde{g}_{top}} \right] \right\}
\]

\[
T^\alpha_{\gamma R} = ee_\gamma C_2(F) \left\{ K^{* \alpha}_{a2L} K^{\alpha}_{a3L} [ s_\Theta C_{11}^{\alpha} + c_\Theta C_{22}^{\alpha} ] - K^{* \alpha}_{a2L} K^{\alpha}_{a3R} c_\Theta s_\Theta \left[ C_{11}^{\alpha}_{\tilde{g}_{top}} - C_{22}^{\alpha}_{\tilde{g}_{top}} \right] \right\}
\]
\[ V_\alpha^L = g_s T_\alpha \left\{ K^{*g}_{\alpha 2L} K^{g}_{\alpha 3L} \left\{ \left[ -\frac{1}{2} C_2(G) + C_2(F) \right][s^2_{\Theta_a} C^{11a}_\epsilon + s^2_{\Theta_a} C^{22a}_\epsilon] - C_2(F) s^2_{\Theta_a} C^{1\alpha}_{SE} + s^2_{\Theta_a} C^{2a}_{SE} \right] + \frac{1}{2} C_2(G) \left[ s^2_{\Theta_a} C^{1\alpha}_{\tilde{g}} + C^{2\alpha}_{\tilde{g}} + C^{1\alpha}_{\tilde{t}} + s^2_{\Theta_a} C^{22\alpha}_{\tilde{t}} \right] \right\} \]

\[ V_\alpha^R = g_s T_\alpha \left\{ K^{*g}_{\alpha 2R} K^{g}_{\alpha 3R} \left\{ \left[ -\frac{1}{2} C_2(G) + C_2(F) \right][s^2_{\Theta_a} C^{11a}_\epsilon + s^2_{\Theta_a} C^{22a}_\epsilon] - C_2(F) s^2_{\Theta_a} C^{1\alpha}_{SE} + s^2_{\Theta_a} C^{2a}_{SE} \right] + \frac{1}{2} C_2(G) \left[ s^2_{\Theta_a} C^{1\alpha}_{\tilde{t}} + s^2_{\Theta_a} C^{2\alpha}_{\tilde{t}} \right] \right\} \]

\[ T_\alpha^L = g_s T_\alpha \left\{ K^{*g}_{\alpha 2L} K^{g}_{\alpha 3L} \left\{ \left[ -\frac{1}{2} C_2(G) + C_2(F) \right][s^2_{\Theta_a} C^{11a}_\epsilon + s^2_{\Theta_a} C^{22a}_\epsilon] - C_2(F) s^2_{\Theta_a} C^{1\alpha}_{SE} + s^2_{\Theta_a} C^{2a}_{SE} \right] + \frac{1}{2} C_2(G) \left[ s^2_{\Theta_a} C^{1\alpha}_{\tilde{g}} + s^2_{\Theta_a} C^{2\alpha}_{\tilde{g}} \right] \right\} \]

\[ T_\alpha^R = g_s T_\alpha \left\{ K^{*g}_{\alpha 2R} K^{g}_{\alpha 3R} \left\{ \left[ -\frac{1}{2} C_2(G) + C_2(F) \right][s^2_{\Theta_a} C^{11a}_\epsilon + s^2_{\Theta_a} C^{22a}_\epsilon] - C_2(F) s^2_{\Theta_a} C^{1\alpha}_{SE} + s^2_{\Theta_a} C^{2a}_{SE} \right] + \frac{1}{2} C_2(G) \left[ s^2_{\Theta_a} C^{1\alpha}_{\tilde{t}} + s^2_{\Theta_a} C^{2\alpha}_{\tilde{t}} \right] \right\} \]

\[ V_{ZL} = \frac{\epsilon}{s_W c_W} C_2(F) \left\{ K^{*g}_{\alpha 2L} K^{g}_{\alpha 3L} \left\{ \left( T_3 L c^2_{\Theta_a} - e q s^2_W \right) c^2_{\Theta_a} C^{11a}_\epsilon + \left( T_3 L s^2_{\Theta_a} - e q s^2_W \right) s^2_{\Theta_a} C^{22a}_\epsilon \right\} + T_3 L c^2_{\Theta_a} s^2_{\Theta_a} C^{1\alpha}_{SE} + s^2_{\Theta_a} C^{2\alpha}_{SE} \right\} \]

\[ V_{ZR} = \frac{\epsilon}{s_W c_W} C_2(F) \left\{ K^{*g}_{\alpha 2R} K^{g}_{\alpha 3R} \left\{ \left( T_3 L c^2_{\Theta_a} - e q s^2_W \right) s^2_{\Theta_a} C^{11a}_\epsilon + \left( T_3 L s^2_{\Theta_a} - e q s^2_W \right) c^2_{\Theta_a} C^{22a}_\epsilon \right\} - T_3 L c^2_{\Theta_a} s^2_{\Theta_a} C^{1\alpha}_{SE} + \epsilon q s^2_W C^{2\alpha}_{SE} \right\} \]
\[ T_{ZL}^\alpha = \frac{e}{s_W c_W} C_2(F) \left\{ K_{\alpha 2L}^g K_{\alpha 3R}^\tilde{g} \left( (T_3 L c_{\theta_\alpha} - e_q s_W^2) s_{\theta_\alpha} c_{\theta_\alpha}^{11\alpha} + (T_3 L s_{\theta_\alpha} - e_q s_W^2) c_{\theta_\alpha}^{22\alpha} - T_3 L c_{\theta_\alpha}^{2}\alpha s_{\theta_\alpha} (C_{\theta_\alpha}^{12\alpha} + C_{\theta_\alpha}^{21\alpha}) \right) \right. \\
\left. - K_{\alpha 2L}^g K_{\alpha 3R}^\tilde{g} c_{\theta_\alpha} s_{\theta_\alpha} \left[ (T_3 L c_{\theta_\alpha} - e_q s_W^2) c_{\theta_\alpha}^{11\alpha} + (T_3 L s_{\theta_\alpha} - e_q s_W^2) c_{\theta_\alpha}^{22\alpha} - T_3 L [c_{\theta_\alpha}^{2}\alpha C_{\theta_\alpha}^{12\alpha} - s_{\theta_\alpha}^{2} C_{\theta_\alpha}^{21\alpha}] \right] \right\} \]

\[ T_{ZR}^\alpha = \frac{e}{s_W c_W} C_2(F) \left\{ K_{\alpha 2L}^g K_{\alpha 3L}^\tilde{g} \left( (T_3 L c_{\theta_\alpha} - e_q s_W^2) c_{\theta_\alpha}^{11\alpha} + (T_3 L s_{\theta_\alpha} - e_q s_W^2) s_{\theta_\alpha}^{22\alpha} + T_3 L c_{\theta_\alpha}^{2}\alpha s_{\theta_\alpha} (C_{\theta_\alpha}^{12\alpha} + C_{\theta_\alpha}^{21\alpha}) \right) \right. \\
\left. - K_{\alpha 2L}^g K_{\alpha 3R}^\tilde{g} c_{\theta_\alpha} s_{\theta_\alpha} \left[ (T_3 L c_{\theta_\alpha} - e_q s_W^2) c_{\theta_\alpha}^{11\alpha} + (T_3 L s_{\theta_\alpha} - e_q s_W^2) s_{\theta_\alpha}^{22\alpha} + T_3 L [s_{\theta_\alpha}^{2} C_{\theta_\alpha}^{12\alpha} - c_{\theta_\alpha}^{2}\alpha C_{\theta_\alpha}^{21\alpha}] \right] \right\} \]

\[ C_{\epsilon}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \left[ \frac{1}{\epsilon} - \gamma + \ln(4\pi\mu^2) - \ln(f_{kl}^\alpha) \right] \]

\[ C_{\text{top}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{m_{\text{top}}^{2} \alpha_1 (1 - \alpha_1 - \alpha_2)}{f_{kl}^\alpha} \]

\[ C_{\tilde{g}\text{top}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{m_{\tilde{g}} m_{\text{top}} (1 - \alpha_1 - \alpha_2)}{f_{kl}^\alpha} \]

\[ C_{SE}^{klo} = \int_0^1 d\alpha_1 \alpha_1 \left[ \frac{1}{\epsilon} - \gamma + \ln(4\pi\mu^2) - \ln(g_{k}^\alpha) \right] \]

\[ C_{\text{SEG}}^{klo} = \frac{1}{\alpha_1} C_{SE}^{klo} \]

\[ C_{\tilde{g}\tilde{g}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \left[ \frac{1}{\epsilon} - \gamma - 1 + \ln(4\pi\mu^2) - \ln(h_{k}^\alpha) \right] \]

\[ C_{\tilde{g}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{m_{\tilde{g}}^2}{h_{k}^\alpha} \]

\[ C_{q_{\alpha}^{2}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{q_{\alpha}^{2} \alpha_1 \alpha_2}{h_{k}^\alpha} \]

\[ C_{t_{\alpha}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{m_{\text{top}}^{2} \alpha_1 (1 - \alpha_1 - \alpha_2)}{h_{k}^\alpha} \]

\[ C_{\tilde{g}t_{\alpha}}^{klo} = \int_0^1 d\alpha_1 \int_0^{1-\alpha_1} d\alpha_2 \frac{m_{\tilde{g}} m_{\text{top}} (1 - \alpha_1 - \alpha_2)}{h_{k}^\alpha} \]

\[ f_{kl}^\alpha = m_{\tilde{g}}^2 - (m_{\tilde{g}}^2 - m_{\tilde{q}_k}^2) \alpha_1 - (m_{\tilde{g}}^2 - m_{\tilde{q}_l}^2) \alpha_2 - m_{\text{top}}^{2} \alpha_1 (1 - \alpha_1 - \alpha_2) - q_{\alpha}^{2} \alpha_1 \alpha_2 \]
\[ g_k^\alpha = m_g^2 - (m_q^2 - m_{\tilde{g}}^2)\alpha_1 - m_{\text{top}}^2\alpha_1(1 - \alpha_1) \]
\[ h_k^\alpha = m_{\tilde{q}_k}^2 - (m_{\tilde{q}_k}^2 - m_{\tilde{g}}^2)(\alpha_1 + \alpha_2) - m_{\text{top}}^2\alpha_1(1 - \alpha_1 - \alpha_2) - q^2\alpha_1\alpha_2 \]

where \( \epsilon = 2 - d/2 \), \( C_2(F) = 4/3 \) and \( C_2(G) = 3 \) for SU(3). If \( \alpha \neq \text{top} \) we have \( c_{\Theta_{\alpha}} = 1 \).

Using the spin-1 condition \( (q_\mu = (p_1 + p_2)_\mu = 0) \) we can write \( P_\mu = (p_1 + p_2)_\mu = 2p_1\mu \). \( K_{\alpha_i L,R}^\tilde{g} \) is the SUSY–Kobayashi–Maskawa matrix; which, as explained in [17], will be parameterized by a small number \( \epsilon \) (not to be confused with the \( \epsilon \) above) to be taken as \( \epsilon^2 = 1/4 \) [16, 20]. It is straightforward at this point to verify that all divergent terms cancel exactly in a nontrivial way, without making use of the GIM mechanism. The results of [17] are reproduced with \( K_{\alpha_{2,3} R}^\tilde{g} = 0 \), that is \( V_{V L,R} = 0 = T_{V L} \). Note that with \( K_{\alpha_{2,3} R}^\tilde{g} \neq 0 \) we obtain terms proportional to the gluino mass, which might become dominant for large gluino masses.

A further crucial test is also provided by the nature of the current. Using the following identity:

\[ \bar{u}_{p_2} \frac{P_\mu}{m_{\text{top}}} P_{L,R} u_{p_1} \equiv \bar{u}_{p_2} \left[ \gamma_\mu P_{R,L} + i\sigma_{\mu\nu} \frac{q^\mu}{m_{\text{top}}} P_{L,R} \right] u_{p_1} \]  

and after Feynman integration with:

\[ \left[ C_{\epsilon}^{ii\alpha} + C_{\text{top}}^{ii\alpha} - C_{SE}^{ii\alpha} \right] q^2 = 0 \]
\[ \left[ C_{\epsilon}^{ii\alpha} + C_{\text{top}}^{ii\alpha} - C_{\tilde{g}\alpha}^{ii\alpha} - C_{\tilde{g}}^{ii\alpha} \right] q^2 = 0 \]
\[ \left[ \frac{m_{\tilde{g}}}{m_{\text{top}}} \left( C_{SE}^{ii\alpha} - C_{SE}^{ii\alpha} \left| m_{\text{top}}^2 = 0 \right. \right) - C_{\tilde{g}\text{top}}^{ii\alpha} \right] q^2 = 0 \]
\[ \left[ C_{\tilde{g}\text{top}}^{ii\alpha} - C_{\tilde{g}\text{top}}^{ii\alpha} \right] q^2 = 0 \]  

we can show that the quantity in front of the \( \gamma^\mu \) term vanishes in the limit \( q^2 \to 0 \), as required by gauge invariance, that is \( V_{V L,R} = -T_{V R,L} \) for \( V = \gamma, g \). For \( V = Z \), that is \( q^2 = m_Z^2 \), the relations above do not hold anymore. We do the first Feynman integration by hand and the second one numerically\(^1\).

In a recent paper [22] one of us (H.K.) considered the gluino and neutralino contributions to the direct CP violating parameter \( \epsilon' \). The Feynman diagrams and calculations were

\[^1\text{We think that in the computer age it is not necessary to present the results in the form of the Passarino-Veltman functions, which would make the results only more difficult to read, but refer the interested reader to [21], where similar calculations have been done. See also [18, 23].}\]
similar. It is straightforward to show that eq. (3) reproduce the eq. (A.9) in \cite{22} by replacing $m_{top}$ with $m_s$ and putting the down quark there to zero.

We assumed that both couplings of the gluino to the left- and right-handed quarks and their superpartners are flavour non diagonal and to be of the same order, that is we take $K_{aR}^g = e^{-i\Phi_S} K_{ab}$ and $K_{aL}^g = e^{+i\Phi_S} K_{ab}$, where $\Phi_S$ is a supersymmetric CP violation phase \cite{22}.

In eq. (3) this phase only comes in when $K_{ab}^g$ is multiplied by $K_{ab}^g$ and, as can be seen, these terms are proportional to the gluino mass. However this SUSY CP violating phase is strongly bounded by the electric dipole moment of the neutron (EDMN) to be of the order of $10^{-2} - 10^{-3}$, if not the SUSY masses are heavier than several TEV's (see references given in \cite{22}). We are not interested here in the consequences of this phase leading to CP violating flavour changing top quark decay, which will be presented elsewhere \cite{19}. In the following we put $\Phi_S = 0$.

When summing over all scalar quarks within the loops, the scalar up quark contributions cancels because of the unitarity of $K_{ab}$, and with $K_{ab} = -K_{ba}$ the mass splitting of the scalar top quark and the scalar charm quark comes into account. This was taken to be $m_\xi = 0.9 m_t$ in \cite{16}, and therefore too small for a top quark mass of 174 GeV. If all scalar quark masses would be the same, the decay rate of $t \to cV$ would be identical to 0. As a final result we obtain:

$$\Gamma_S(t \to cV) = \frac{\alpha_s^2}{128\pi^3} m_{top} \left( 1 - \frac{m_\xi^2}{m_{top}^2} \right)^2 \varepsilon^2 \left[ (V_{VL}^2 + V_{VR}^2) \left( 2 + \frac{m_{top}^2}{m_V^2} \right) - 2(V_{VL} T_{VR} + V_{VR} T_{VL}) \left( 1 - \frac{m_{top}^2}{m_V^2} \right) - (T_{VL}^2 + T_{VR}^2) \left( 2 - \frac{m_V^2}{m_{top}^2} \right) - \right]$$

where $V_{VL,R} = V_{VL,R}^\dagger - V_{VL,R}^\dagger$ and $T_{VL,R} = T_{VL,R}^\dagger - T_{VL,R}^\dagger$. As explained above for $V = \gamma, g$ we have $V_{VL,R} = -T_{VL,R}$ and all terms containing $m_V^2$ are absent.

We define \cite{5}: $B(t \to cV) = \Gamma_S(t \to cV)/\Gamma_W(t \to bW^+)$ where

$$\Gamma_W(t \to bW^+) = \frac{\alpha}{16 \sin^2 \Theta_W} m_{top} \left( 1 - \frac{m_{W^+}^2}{m_{top}^2} \right)^2 \left( 2 + \frac{m_{top}^2}{m_{W^+}^2} \right)$$

Our input parameters are $m_{top} = 174$ GeV and the strong coupling constant $\alpha_s = 1.4675/\ln(m_{top}^2/\Lambda_{QCD}^2) = 0.107$ with $\Lambda_{QCD} = 0.18$ GeV \cite{3}.
3 Discussions

To compare the new results with flavour changing couplings in the right- and left-handed sector with the ones already presented in \[17\], where flavour changing couplings only in the left-handed sector was considered, we present the same plots as in \[17\]. The general discussion remains the same and we will only present the changes when flavour changing in the right-handed sector is included.

In Fig. 1 we present the branching ratio $B(t \rightarrow cZ)$ as a function of the scalar mass $m_S$ for a gluino mass of 100 GeV. We see that without mixing, the branching ratio decreases rapidly with increasing scalar mass and is hardly changed when flavour changing in the right-handed sector is included. However the mixing has a drastic effect. It enhances the branching ratio by up to 4 orders of magnitude for large $m_s$ and is enhanced by another factor of 5 when flavour changing occurs in both sectors.

In Fig. 2 we consider the same cases as in Fig. 1 but for $B(t \rightarrow cg)$. As before without mixing the results remain almost the same whether or not flavour changing in the right-handed sector is included. However when mixing is taken into account the results are changed drastically up to 7 orders of magnitude for large values of the scalar mass $m_S$ when flavour changing is considered in both sectors, compared with the case where flavour changing occurs only in the left-handed sector.

In Fig. 3 we consider the branching ratio $B(t \rightarrow c\gamma)$ As in the cases before, without mixing there is almost no difference between the results with flavour changing only in the left-handed sector or in both sectors. As in Fig. 2 the results are changed drastically, up to 6-7 orders of magnitude for large values of the scalar mass, when mixing is taken into account and flavour changing is considered in the left- and right-handed sector, compared with the case where flavour changing occurs only in left-handed sector.

A further important consequence is that the GIM-like supression where the contribution of the top quark exactly cancels the contribution from the c-quark is pushed to much smaller values of the scalar mass $m_S$. We have tried many different combinations of $\mu$ and $m_\tilde{g}$ and the cancellation is always pushed to smaller values of the scalar mass.
Figure 1: The ratio $\Gamma_S/\Gamma_W$ of the the top quark decay into a charm quark and $Z^0$ boson as a function of the scalar mass $m_S$. The gluino mass was taken to be 100 GeV. The solid line is the unphysical case with no mixing ($\mu = 0 = A_{\text{top}}$) and $\tan \beta = 10$, the dotted line the same case when flavour changing $g - q - \bar{q}$ in the right-handed sector is included. The other cases are with mixing ($A_{\text{top}} = m_S$). The dashed–dotted ones with $\mu = 500$ GeV and $\tan \beta = 10$. The shorter ones are with flavour changing in both sectors.
Figure 2: The same as Fig. 1 but for the decay of the top quark into a charm quark and a gluon.
Figure 3: The same as Fig. 1 but for the decay of the top quark into a charm quark and a photon.
4 Conclusions

In this paper we presented the supersymmetric QCD 1-loop correction to the flavour changing decay rate $t \to cV$. We included flavour changing $g - q - \tilde{q}$ couplings in the left- and right-handed sector, thus extending the previous analysis of [17], where flavour changing was only considered in the left-handed sector. We have shown that the results remain almost the same when mixing of the scalar top quark is neglected. This remains true for the $t \to cZ$ decay rate even when mixing is included. However the results are changed drastically, up to 7 orders of magnitude for the decay rates $t \to cg$ and $t \to c\gamma$ when mixing of the scalar top quark is included and flavour changing couplings are taken in both sectors. Furthermore in the $t \to c\gamma$ decay mode the GIM-like cancellation of the scalar top and charm quarks is pushed to much smaller values of the scalar mass $m_S$.

Note: While completing this work we have seen a paper by an Italian group [23], where the same processes were considered. Their statement is that the SUSY mixing angle between the second and the third generation ($K_{23} = \varepsilon$) has been over-estimated by at least one order of magnitude in our first paper [17]. There and in this present paper we took $\varepsilon$ as a free parameter and have taken it pretty large following the spirit of former papers. From eq.(6) it is obvious that the results are diminished drastically if smaller values are taken for $\varepsilon$. However the authors of [23] showed that relaxing the universality constraints on soft SUSY mass breaking terms of the off-diagonal squark masses between $\tilde{c}$ and $\tilde{t}$ reintroduces a large $\varepsilon$, that is a large mixing angle between $\tilde{c}$ and $\tilde{t}$.

They also find a difference in the result for the amplitude which can be traced back to the omission in [17] of the diagrams involving a helicity flip in the gluino line, which dominate the branching ratios when the gluino mass gets large. However in [17] we considered flavour changing only in the left-handed sector as is usually done in the MSSM and therefore no gluino helicity flip was possible, that is no term proportional to the gluino mass is introduced. In the present work, we also took into account flavour changing in the right-handed sector and as a consequence the mentioned effect occurs, which is expressed by the new terms proportional to the gluino mass in eq.(3).
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