UAB-FT-430
KA-TP-1-1998
January, 1998
hep-ph/9802329

YUKAWA-COUPLING CORRECTIONS TO SCALAR QUARK DECAYS

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ABSTRACT

Heavy squark decays into top and charginos or neutralinos could be an unexpected source of top quarks at hadron colliders. A detailed treatment of these processes is necessary for a reliable calculation of both the top quark production cross-section and the standard top quark branching ratio in the MSSM. Along this line we compute the electroweak corrections to the sbottom decays $\tilde{b}_a \to \chi_i^- t$ within the Yukawa coupling approximation. The calculation of these higher order contributions requires renormalization of both the bottom squark mixing angle and of $\tan\beta$. This type of corrections gives the leading order electroweak effects at low and high $\tan\beta$.

A fundamental missing link between the highly successful Standard Model of the Elementary Particle Physics and the highly developed status of the experiments is the final elucidation of the nature of the spontaneous symmetry-breaking mechanism (SSB). Whether the understanding of the SSB will eventually point towards the existence of elementary spinless particles (Higgs bosons) or towards the finding of dynamically bound fermion-antifermion spinless states, is far from clear at present. Ultimately, experiment must decide. In the meanwhile, the Minimal Supersymmetric extension of the Standard Model (MSSM) remains immaculately consistent with all known high precision experiments at a level comparable to the SM [1]. This fact alone, if we bare in mind the vast amount of high precision data available at present both from low-energy and high-energy physics, should be considered at least as greatly notorious. To this fact we have to add another prominent feature, to wit: that the MSSM offers a starting point for a successful Grand Unified framework where a radiatively stable low-energy Higgs sector can survive. Putting things together it is well justified, we believe, to keep alive all efforts on all fronts trying to discover a supersymmetric particle. Undoubtedly, the next Tevatron run, and of course also the advent of the LHC, should offer us a gold-plated scenario for finding a hint of SUSY, if it is there at all.

Sparticles not much heavier than a few hundred GeV could be produced in significant numbers already at the Tevatron. For instance, selectron production was advocated in Ref.[2] to explain a purported non-SM event in the Collider Detector at Fermilab (CDF). Subsequently, in Ref.[3] it was argued that half of the top quarks at the Tevatron might come from gluino decays into top and stop, $\tilde{g} \rightarrow t\,\tilde{t}_1$. Similarly, we may envision the possibility that sbottom squarks are pair produced by the usual Drell-Yan mechanism and then decay into top quark and charginos: $\tilde{b}_a \rightarrow \chi_i^-\,t$. Indeed, this would be the leading two-body decay if gluinos are heavy enough that the strong decay mode $\tilde{b}_a \rightarrow \tilde{g}\,b$ is kinematically blocked up[1]. The observed cross-section is then equal to the Drell-Yan production cross-section convoluted over the parton distributions times the squared branching ratio. However, in the framework of the MSSM, we rather expect a generalization of this formula in the following way (schematically),

$$
\begin{aligned}
\sigma_{\rm obs.} \;=\; & \int dq\,d\bar{q}\;\sigma(q\,\bar{q} \rightarrow t\,\bar{t}) \times |BR(t \rightarrow W^+\,b)|^2 \\
+ \; & \int dq\,d\bar{q}\;\sigma(q\,\bar{q} \rightarrow \tilde{b}_a\,\bar{\tilde{b}}_a) \times |BR(\tilde{b}_a \rightarrow \chi_1^-\,t)|^2 \times |BR(t \rightarrow W^+\,b)|^2 + \ldots \qquad (1)
\end{aligned}
$$

We assume that gluinos are much heavier than squarks, so that their contribution to this cross-section through $q\,\bar{q} \rightarrow \tilde{g}\,\tilde{g}$ followed by $\tilde{g} \rightarrow t\,\tilde{t}_1$ is negligible. From eq.(1) we see that, if there are alternative (non-SM) sources of top quarks subsequently decaying into

[1] Squark decays have been discussed at the tree-level in several places of the literature. See e.g. Ref.[4] for some relatively recent references on the subject.

the SM final state, $W^+ b$, then one cannot rigorously place any stringent lower bound on $BR(t \to W^+ b)$ in the MSSM from the present FNAL data. Indeed, we could as well have non-SM top quark decay modes, such as e.g. $t \to \tilde{t}_a \chi^0_\alpha$ [3] and $t \to H^+ b$ [5, 6], that could serve, pictorially, as a "sinkhole" to compensate (at least in part) for the unseen source of extra top quarks produced at the Tevatron from sbottom pair production (Cf. eq.(1)). However, one usually assumes that $BR(t \to W^+ b) \geq 50\%$ in order to guarantee the (purportedly) standard top quark events observed at the Tevatron. Thus, from these considerations it is not excluded that the non-SM branching ratio of the top quark, $BR(t \to$"new"), could be as big as the SM one, i.e. $\sim 50\%$. We stress that at present one cannot exclude eq.(1) since the observed $t \to W^+ b$ final state involves missing energy, as it is also the case for the decays comprising supersymmetric particles. In particular, if $\tan\beta$ is large and there exists a relatively light chargino with a non-negligible higgsino component, the alternative mechanism suggested in eq.(1) could be a rather efficient non-SM source of top quarks that could compensate for the depletion in the SM branching ratio.

While the squark production cross-section has already received some attention in the literature at the level of NLO radiative corrections [7], an accurate treatment of the decay mechanisms is also very important to provide a solid basis for experimental analyses of the top quark production in the MSSM. Thus in this paper we consider the computation of the leading supersymmetric electroweak (SUSY-EW) quantum effects on $\tilde{b}_a \to \chi^-_i t$, namely the ones induced by potentially large Yukawa-couplings from the top and bottom quarks (normalized with respect to the $SU(2)_L$ gauge coupling):

$$\lambda_t \equiv \frac{h_t}{g} = \frac{m_t}{\sqrt{2}\, M_W \,\sin\beta} \quad , \quad \lambda_b \equiv \frac{h_b}{g} = \frac{m_b}{\sqrt{2}\, M_W \,\cos\beta}\,, \tag{2}$$

At large (≥ 20) or small (< 1) $\tan\beta$ [8] these effects could be competitive with the known QCD corrections [9]. Since in these conditions the full MSSM quantum effects can be rather large, their calculation is indispensable to account for the observed top quark production cross-section (1) in the MSSM or, alternatively, to better assess how much the determination of the SM branching ratio $BR(t \to W^+ b)$ is affected in the MSSM context after plugging in the experimental number on the LHS of eq.(1).

We address the calculation of the one-loop corrections to the partial width of $\tilde{b}_a \to \chi^-_i t$ within the context of the on-shell renormalization framework [10]. However, in SUSY extensions of the SM, there are additional ingredients concerning the renormalization program that must be fixed. Consider the MSSM interaction Lagrangian involving the top-sbottom-chargino vertex in the mass-eigenstate basis:

$$\mathcal{L}_{t\tilde{b}\chi} = -g\, \tilde{b}^*_a \bar{\chi}^+_i (A^{ai}_+ P_L + \epsilon_i\, A^{ai}_- P_R) t + \text{h.c.}\,, \tag{3}$$

where ϵ_i is the sign of the ith chargino eigenvalue $M_i\,(i=1,2$ with $|M_1|<|M_2|)$ in the real matrix representation, and the coupling matrices are denoted by[2]

$$A_+^{ai} = R_{1a} V_{i1} - \lambda_b R_{2a} V_{i2}\,, \qquad A_-^{ai} = -R_{1a}\lambda_t U_{i2}\,. \tag{4}$$

The explicit appearance of the Yukawa couplings (2) in the above Lagrangian requires both the introduction of top and bottom quark mass counterterms (in the on-shell scheme) and also a suitable prescription for $\tan\beta$ renormalization. We denote by $m_{\tilde{b}_a}\,(a=1,2)$, with $m_{\tilde{b}_1} < m_{\tilde{b}_2}$, the two sbottom mass eigenvalues. The sbottom mixing angle $\theta_{\tilde{b}}$ is defined by the transformation relating the weak-interaction $(\tilde{b}'_a = \tilde{b}_L, \tilde{b}_R)$ and the mass eigenstate $(\tilde{b}_a = \tilde{b}_1, \tilde{b}_2)$ squark bases:

$$\tilde{b}'_a = R_{ab}\,\tilde{b}_b\;;\quad R = \begin{pmatrix} \cos\theta_{\tilde{b}} & -\sin\theta_{\tilde{b}} \\ \sin\theta_{\tilde{b}} & \cos\theta_{\tilde{b}} \end{pmatrix}, \tag{5}$$

where R is the matrix appearing in eq.(4). From this change of basis the sbottom mass matrix,

$$\mathcal{M}_{\tilde{b}}^2 = \begin{pmatrix} M_{\tilde{b}_L}^2 + m_b^2 + \cos 2\beta(-\frac{1}{2}+\frac{1}{3}\,s_W^2)\,M_Z^2 & m_b\,(A_b - \mu\tan\beta) \\ m_b\,(A_b - \mu\tan\beta) & M_{\tilde{b}_R}^2 + m_b^2 - \frac{1}{3}\,\cos 2\beta\, s_W^2\, M_Z^2\,, \end{pmatrix}, \tag{6}$$

becomes diagonalized as follows: $R^\dagger\,\mathcal{M}_{\tilde{b}}^2\,R = \mathrm{diag}\left\{m_{\tilde{b}_2}^2, m_{\tilde{b}_1}^2\right\}$.

Formally, we start the renormalization procedure as usual [10] i.e. by introducing parameter and field renormalization constants:

$$\begin{aligned}
g &\to (1+\frac{\delta g}{g})\,g\,,\\
m_q &\to (1+\frac{\delta m_q}{m_q})\,m_q\,,\\
t &\to (1+\frac{1}{2}\delta Z_L^t P_L + \frac{1}{2}\delta Z_R^t P_R)\,t\,,\\
\chi_i^+ &\to (1+\frac{1}{2}\delta Z_L^i P_L + \frac{1}{2}\delta Z_R^i P_R)\,\chi_i^+\,,\\
\tilde{b}_a &\to (1+\frac{1}{2}\delta Z^a)\,\tilde{b}_a + \delta Z^{ab}\tilde{b}_b \quad (b\neq a)\,,\\
\lambda_{b,t} &\to (1+\frac{\delta\lambda_{b,t}}{\lambda_{b,t}})\,\lambda_{b,t}\,,\\
\tan\beta &\to (1+\frac{\delta\tan\beta}{\tan\beta})\,\tan\beta\,,\\
R &\to R+\delta R.
\end{aligned} \tag{7}$$

Notice that the sbottom wave-function renormalization involves the mixed scalar field renormalization constant $(b\neq a)$

$$\delta Z^{ab} = \frac{\Sigma^{ab}(m_{\tilde{b}_b}^2)}{m_{\tilde{b}_b}^2 - m_{\tilde{b}_a}^2}\,, \tag{8}$$

[2]See Refs.[6, 11] for full notation niceties.

where Σ^{ab} is the bare mixed self-energy connecting the physical states $\tilde{b}_a$ and $\tilde{b}_b$ ($a \neq b$). The renormalization of the Yukawa couplings (2) involves the mass counterterms and the counterterm to $\tan\beta$ through the relations

$$\begin{aligned}
\frac{\delta\lambda_b}{\lambda_b} &= \frac{\delta m_b}{m_b} - \frac{\delta M_W}{M_W} + \sin^2\beta\,\frac{\delta\tan\beta}{\tan\beta}\,,\\
\frac{\delta\lambda_t}{\lambda_t} &= \frac{\delta m_t}{m_t} - \frac{\delta M_W}{M_W} - \cos^2\beta\,\frac{\delta\tan\beta}{\tan\beta}\,.
\end{aligned} \tag{9}$$

Furthermore, as the last row of eq.(7) shows, we have introduced a counterterm δR for the rotation matrix R that diagonalizes the sbottom mass matrix. This counterterm can be replaced (see below) by a counterterm $\delta\theta_{\tilde{b}}$ for the sbottom mixing angle. Indeed, we shall proceed this way as we wish to take $\theta_{\tilde{b}}$ as an input in our renormalization program.

Substituting the relations (7) into the bare Lagrangian (3) we find:

$$\mathcal{L}_{\rm one-loop} = -g\,\tilde{b}_a^*\bar{\chi}_i^+\left[\left(A_+^{ai} + \delta C_+^{ai}\right)P_L + \epsilon_i\left(A_-^{ai} + \delta C_-^{ai}\right)P_R\right]t + {\rm h.c.}\,, \tag{10}$$

with

$$\begin{aligned}
\delta C_+^{ai} &= A_+^{ai}\left(\frac{\delta g}{g} + \frac{1}{2}\delta Z^a + \frac{1}{2}\delta Z_R^i + \frac{1}{2}\delta Z_L^t\right) + \delta A_+^{ai} + \delta Z^{ba}A_+^{bi}\,,\\
\delta C_-^{ai} &= A_-^{ai}\left(\frac{\delta g}{g} + \frac{1}{2}\delta Z^a + \frac{1}{2}\delta Z_L^i + \frac{1}{2}\delta Z_R^t\right) + \delta A_-^{ai} + \delta Z^{ba}A_-^{bi}\,.
\end{aligned} \tag{11}$$

The full structure of the four on-shell renormalized decay amplitudes for $\tilde{b}_a \to \chi_i^-\,t$ ($a = 1,2; i = 1,2$) follows from the previous Lagrangian after including the contributions from the (LH and RH) one-loop vertex form factors $F_{L,R}^{ai}$:

$$i\,T(\tilde{b}_a \to t\,\chi_i^-) = i\,g\,\bar{u}_t\left[\epsilon_i\left(A_-^{ai} + \Lambda_L^{ai}\right)P_L + \left(A_+^{ai} + \Lambda_R^{ai}\right)P_R\right]v_i \tag{12}$$

where

$$\Lambda_L^{ai} = \delta C_-^{ai} + \epsilon_i\,F_L^{ai}\,, \qquad \Lambda_R^{ai} = \delta C_+^{ai} + F_R^{ai}\,. \tag{13}$$

Let now Γ_0^{ai} be the tree-level partial width of the decay $\tilde{b}_a \to \chi_i^-\,t$ obtained from the Lagrangian (3):

$$\Gamma_0^{ai} = \frac{g^2}{16\,\pi\,m_{\tilde{b}_a}^3}\,\lambda(a,i,t)\left\{\left[(A_+^{ai})^2 + (A_-^{ai})^2\right]\left(m_{\tilde{b}_a}^2 - M_i^2 - m_t^2\right) - 4\,A_+^{ai}\,A_-^{ai}\,m_t\,\epsilon_i\,M_i\right\}\,, \tag{14}$$

with $\lambda(a,i,t) \equiv \lambda(m_{\tilde{b}_a}^2, M_i^2, m_t^2)$ the usual Källen function for the given arguments. The quantum correction to Γ_0^{ai} can be described in terms of the quantities $\delta^{ai} = (\Gamma^{ai} - \Gamma_0^{ai})/\Gamma_0^{ai}$, where Γ^{ai} is the corresponding one-loop corrected width. From the previous formulae, δ^{ai} can be worked out as follows:

$$\delta^{ai} = \frac{2\left(m_{\tilde{b}_a}^2 - M_i^2 - m_t^2\right)\left(A_-^{ai}\,\Lambda_L^{ai} + A_+^{ai}\,\Lambda_R^{ai}\right) - 4\,m_t\,\epsilon_i\,M_i\left(A_-^{ai}\Lambda_R^{ai} + A_+^{ai}\,\Lambda_L^{ai}\right)}{\left[(A_+^{ai})^2 + (A_-^{ai})^2\right]\left(m_{\tilde{b}_a}^2 - M_i^2 - m_t^2\right) - 4\,A_+^{ai}\,A_-^{ai}\,m_t\,\epsilon_i\,M_i}\,. \tag{15}$$

For the further procedure, we have to fix the renormalization conditions that determine the counterterms $\delta\tan\beta$ and $\delta\theta_{\tilde{b}}$ which enter the explicit structure of $\delta A_{\pm}^{ai}$ in eq.(11). Following Ref.[6] we use the decay process $H^+ \to \tau^+\nu_\tau$ to define the parameter $\tan\beta$ through

$$\Gamma(H^+ \to \tau^+\nu_\tau) = \frac{g^2 m_\tau^2\, M_H}{32\,\pi\, M_W^2}\,\tan^2\beta = \frac{G_F m_\tau^2\, M_H}{4\pi\sqrt{2}}\,\tan^2\beta\,(1-\Delta r^{MSSM})\,. \tag{16}$$

This determines the counterterm $\delta\tan\beta$ [6] as follows:

$$\frac{\delta\tan\beta}{\tan\beta} = \frac{1}{2}\left(\frac{\delta M_W^2}{M_W^2} - \frac{\delta g^2}{g^2}\right) - \frac{1}{2}\delta Z_{H^\pm} + \cot\beta\,\delta Z_{HW} + \Delta_\tau\,, \tag{17}$$

where Δ_τ stands for the full set of MSSM one-loop corrections to the partial width of $H^+ \to \tau^+\,\nu_\tau$. For fixing $\delta\theta_{\tilde{b}}$, we require that the renormalized sbottom mixing angle (that we use as an input data, Cf. eq.(20) below) does not feel a shift from the mixed sbottom bare self-energies Σ^{ab} between the physical states $\tilde{b}_a$ and $\tilde{b}_b$ $(a \neq b)$. This is similar to the prescription adopted in Refs.[9, 12], though it is not identical. In our formalism, the 3-point Green functions explicitly incorporate the mixed field renormalization constants δZ^{ab} $(a \neq b)$ and are therefore renormalized also in the $\theta_{\tilde{b}}$ parameter. The UV-divergent parts of the 3-point functions are cancelled against $\delta\theta_{\tilde{b}}$ by defining the latter as follows:

$$\delta\theta_{\tilde{b}} = \frac{1}{2}\,(\delta Z^{12} - \delta Z^{21}) = \frac{1}{2}\,\frac{\Sigma^{12}(m_{\tilde{b}_2}^2) + \Sigma^{12}(m_{\tilde{b}_1}^2)}{m_{\tilde{b}_2}^2 - m_{\tilde{b}_1}^2}\,. \tag{18}$$

Of course another equivalent choice could just be $\delta\theta_{\tilde{b}} = \delta Z^{12}$ (or $-\delta Z^{21}$), but eq.(18), is more symmetrical; the numerical differences among the finite parts of the two choices are negligible.

The analytical formulation developed thus far is well suited to tackle the general problem of the SUSY-EW corrections to squark decays. Since the dominant part is from the Yukawa sector we wish to pursue our calculation in the following within the Yukawa coupling approximation. This means that we are going to compute the leading electroweak effects of $\mathcal{O}(\lambda_t^2)$ *and* $\mathcal{O}(\lambda_b^2)$ that emerge for large values of the Yukawa couplings (2) when the remaining gauge contributions – of $\mathcal{O}(g^2)$– are subdominant. In practice we shall only explore the large $\tan\beta$ regime, typically $\tan\beta \geq 20$; the possibility $\tan\beta < 1$ is not so appealing from the theoretical point of view. Thus within our approximation we will include the correction Δ_τ in leading order $\mathcal{O}(\lambda_\tau^2)$ of the τ Yukawa-coupling, λ_τ. Notice furthermore that for $\lambda_b >> 1$ the tree-level decay rate, eq. (14), is maximized. Therefore, the large $\tan\beta$ range is expected to be the most relevant one for the decay under consideration.

In our approach, we set the $SU(2)_L$ gaugino mass parameter $M >> |\mu|, M_W$ in the chargino mass matrix [6], and therefore the chargino $\chi_1^\pm$ is mainly higgsino, whereas the chargino $\chi_2^\pm$ is mainly gaugino and does not contribute to our decays. It is only in this case that the Yukawa-coupling approximation makes sense. Thus, since $m_{\tilde{t}_1} > 80 - 90\,GeV$, in this approach the decay into stop and neutralino $t \to \tilde{t}_a\,\chi^0_\alpha$ is kinematically forbidden[3]. Therefore, in this approximation the relevant counterterms $\delta A^{ai}_\pm$ in eq.(11) boil down to

$$\begin{aligned} \delta A^{a2}_+ &= -\delta R_{2a}\lambda_b - R_{1a}\delta\lambda_b \\ \delta A^{a2}_- &= -\delta R_{1a}\,\lambda_t - R_{1a}\,\delta\lambda_t\,, \end{aligned} \tag{19}$$

with $\delta\lambda_b$ and $\delta\lambda_t$ as in eq.(9). The one-loop Feynman diagrams contributing to $\tilde{b}_a \to \chi_i^-\,t$ in our Yukawa-coupling approximation are depicted in Fig. 1. Specifically, in Fig. 1a we show the wave-function renormalization graphs from which the field renormalization constants δZ^a, δZ^i and $\delta Z^t_{L,R}$ are straightforwardly computed; and in Fig.1b we display the contributions going to the vertex form factors $F^{ai}_{L,H}$. We refrain from writing out cumbersome formulas in this letter [13]. Nevertheless we remark that the calculation is well-defined in the Yukawa coupling approximation. In this respect, we have checked the effective cancellation of divergences among the bare diagrams and the various counterterms. From the practical point of view we emphasize that the definition of $\tan\beta$ that we have used to fix $\delta\tan\beta$ (Cf. eq.(16)) is quite convenient as the Tevatron (and the LHC) collaborations have started to develop all the necessary techniques for τ-tagging and $H^\pm$ identification [14].

Let us now pass on to the numerical analysis. It is summarized in Figs. 2-5. The set of independent parameters that we use in the sbottom sector consists of the masses and the mixing angle

$$(m_{\tilde{b}_1}, m_{\tilde{b}_2}, \theta_{\tilde{b}})\,, \tag{20}$$

whereas for the stop sector we just have in addition

$$(m_{\tilde{t}_1}, \theta_{\tilde{t}}) \tag{21}$$

since by $SU(2)_L$ gauge invariance the value of the other stop mass $m_{\tilde{t}_2}$ is already determined. Similarly, the sbottom and stop trilinear terms A_b and A_t are fixed by the previous parameters as follows:

$$A_b = \mu\,\tan\beta + \frac{m^2_{\tilde{b}_2} - m^2_{\tilde{b}_1}}{2\,m_b}\,\sin 2\,\theta_{\tilde{b}}\,; \qquad A_t = \mu\,\cot\beta + \frac{m^2_{\tilde{t}_2} - m^2_{\tilde{t}_1}}{2\,m_t}\,\sin 2\,\theta_{\tilde{t}}\,. \tag{22}$$

We impose the approximate (necessary) condition

$$A_q^2 < 3\,(m^2_{\tilde{t}} + m^2_{\tilde{b}} + M^2_H + \mu^2)\,, \tag{23}$$

where $m_{\tilde{q}}$ is of the order of the average squark masses for $\tilde{q} = \tilde{t}, \tilde{b}$, to avoid colour-breaking minima in the MSSM Higgs potential [15].

[3] See, however, Ref.[13] for a more general discussion.

For the numerical analysis, we assume that the sbottom masses are not too heavy ($m_{\tilde{b}_1} \leq 350\,GeV$) and that the charged Higgs mass, M_H, is such that the decay $t \to H^+ b$ is available. In this way the quantum corrections that we shall evaluate could have an impact already for the Tevatron physics of top quark production. However, in Ref.[13] we include the corresponding study of the radiative corrections to sbottom decays also in the case of very heavy ("obese") sbottom and charged Higgs bosons ($m_{\tilde{q}} \geq 400\,GeV$, $M_H > m_t$) which should be relevant for LHC. Of course, $\tan\beta$ and the SUSY Higgs mixing parameter μ are also additional independent inputs for our calculation. In the relevant large $\tan\beta$ segment under consideration, namely

$$20 \lesssim \tan\beta \lesssim 40\,, \tag{24}$$

the bottom quark Yukawa coupling λ_b is comparable to the top quark Yukawa coupling, λ_t. Even though the extreme interval $40 < \tan\beta < 60$ can be tolerated by perturbation theory, we shall confine ourselves to the moderate range (24). This is necessary to preserve the condition (23) for the typical set of sparticle masses used in our analysis. In particular, notice that with our chosen set of independent inputs in eqs.(20)-(21) the A-parameters are not free –they are fixed by eq.(22). We point out that the colour stability requirement (23) could be satisfied independently of $\tan\beta$ if the A-parameters would be chosen directly as a part of the set of inputs and then taken sufficiently small. Nevertheless this possibility is not so convenient in our analysis where the sparticle masses are the natural inputs that we wish to control in order to make sure that sparticles can be produced and decay at the Tevatron as explained in connection to eq.(1).

In Figs. 2a and 2b we evaluate the branching ratio

$$\begin{aligned} BR_0(\tilde{b}_a \to \chi_1^- t) &= \frac{\Gamma_0(\tilde{b}_a \to \chi_1^- t)}{\Sigma_i \Gamma_i} \\ \Sigma_i \Gamma_i &= \Sigma_\alpha \Gamma_0(\tilde{b}_a \to \chi_\alpha^0 b) + \Gamma_0(\tilde{b}_a \to \chi_1^- t) \\ &+ \Gamma_0(\tilde{b}_a \to H^- \tilde{t}_1) + \Sigma_i\, \Gamma_0(\tilde{b}_a \to \Phi_i^0\, \tilde{b}_b)\,, \end{aligned} \tag{25}$$

(where $\Phi_i^0 = h^0, H^0, A^0$) as a function of $\tan\beta$ and $m_{\tilde{b}_1}$, respectively, in order to illustrate the relevant intervals for these parameters where to compute radiative corrections to our decay. In these figures we have also indicated the ranges of $\tan\beta$ and $m_{\tilde{b}_1}$ excluded by the condition (23) for the given values of the other parameters. The corresponding corrections δ^{ai} (15) are shown in Figs. 3a and 3b as a function of the lightest stop and sbottom masses, respectively. The allowed range for the sbottom and stop mixing angles is conditioned by the upper bound on the trilinear couplings and is obtained from eqs.(22) and (23). In the case of $\theta_{\tilde{b}}$, it is rather narrow, so that the physical sbottom masses basically coincide with the LH and RH electroweak eigenstates. In the physical $\theta_{\tilde{b}}$ range, the variation of the correction (15) is shown in Fig. 4a. On the other hand the permitted range for the stop mixing angle, $\theta_{\tilde{t}}$, is much larger and we have plotted the corrections within the allowed region in Fig. 4b. Notice that the sign of the quantum effects changes within the domain of variation of $\theta_{\tilde{t}}$. Finally, we display the evolution of the SUSY-EW effects as a function of $\tan\beta$ (Fig. 5a) and of μ (Fig. 5b) within the region of compatibility with the constraint (23). We remark that for $|\mu| > 120\,GeV$ and $\tan\beta > 20$ the corrections can be above 20% near the phase space limit of the lightest sbottom decay. However, this effect has nothing to do with the phase space exhaustion, which is described by the kinematic function $\lambda(a,i,t)$ on the RHS of the tree-level expression (14), but rather with the presence of the dynamical factor in brackets on that equation which also goes to the denominator of δ in eq.(15). That factor is fixed by the structure of the interaction Lagrangian of the sbottom decay into charginos and top; and, for the parameters in Fig. 5, it turns out to vanish near (actually past) the phase space limit in the case of the lightest sbottom $(\tilde{b}_1)$ decay. However, this is not so for the heaviest sbottom $(\tilde{b}_2)$ as it is patent in the same figure.

A few more words are in order to explain the origin of the leading electroweak effects. One could expect that they come from the well-known large $\tan\beta$ enhancement stemming from the chargino-stop corrections to the bottom mass[4] (Cf. Fig. 1a). Nonetheless this is only partially true, for in the present case the remaining contributions in Fig. 1 can be sizeable enough. To be more precise, in the region of the parameter space that we have dwelled upon the bottom mass contribution is seen to be dominant only for the lightest sbottom decay and for the lowest values of $\tan\beta$ in the range (24). This is indeed the case in Fig. 4b where $\tan\beta = 20$ and therefore the bottom mass effect modulates the electroweak correction in this process and δ becomes essentially an odd function of the stop mixing angle. On the other hand, from Fig. 5a it is obvious that the (approximate) linear behaviour on $\tan\beta$ expected from bottom mass renormalization becomes completely distorted by the rest of the contributions, especially in the high $\tan\beta$ end. In short, the final electroweak correction cannot be simply ascribed to a single renormalization source but to the full Yukawa-coupling combined yield.

In general the SUSY-EW corrections to $\Gamma(\tilde{b}_a \to \chi_i^- \, t)$ are smaller than the QCD corrections [9]. The reason why the electroweak corrections are smaller is in part due to the condition (23) restricting our analysis within the $\tan\beta$ interval (24). From Figs. 4 and 5a it is clear that outside this interval the SUSY-EW contributions could be much higher and with the same or opposite sign as the QCD effects depending on the choice of the sign of the mixing angles[5]. Moreover, since we have focused our analysis to sbottom

[4]See e.g. [6] and references therein.

[5]In the conditions of Ref.[9] the evolution of the QCD corrections with $\tan\beta$ would be very slow because the assumed sbottom mixing angle is zero. In the present case, however, $\theta_{\tilde{b}} \neq 0$ and the QCD corrections evolve significantly with $\tan\beta$ but become saturated [13].

masses accessible to Tevatron, again the theoretical bound (23) severely restricts the maximum value of the trilinear couplings and this prevents the electroweak corrections from being larger. This cannot be cured by assuming larger values of M_H and/or of μ due to our assumption that $t \to H^+ b$ is operative and because μ directly controls the value of the (higgsino-like) chargino final state in our decay, so that basically we have $|\mu| < m_{\tilde{b}_a} - M_H$. The restriction cannot be circumvented either if we assume larger values of $m_{\tilde{t}_a}$, for it has been shown that too heavy stops are incompatible with the CLEO data on $b \to s\gamma$ both at low and high $\tan\beta$ [16, 17]. We point out that the MSSM analysis of $b \to s\gamma$ also motivated the sign choice $A\mu < 0$ in our numerical calculation [17]. Admittedly, the situation with radiative B-decays is still under study and there are many sources of uncertainties that deserve further experimental consideration. Still, we have used this information to focus on a limited domain of the MSSM parameter space.

Finally, we mention that the treatment of the electroweak corrections to the closely related decay $\tilde{t}_a \to \chi^0_\alpha t$ is similar to the one presented here, but for lack of space it will be presented elsewhere [13]. It suffices to say that the branching ratio and the radiative corrections to that decay can be sizeable only for small values of $\tan\beta$, so that this process becomes optimized in a region of parameter space complementary to the one under consideration.

In summary, the MSSM corrections to squark decays into charginos and neutralinos can be significant and therefore must be included in any reliable analysis of top quark physics at the Tevatron within the MSSM. The main corrections stem from the strongly interacting sector of the theory (i.e. the one involving gluons and gluinos), but also non-negligible effects may appear from the electroweak sector (characterized by chargino-neutralino exchange) at large (or very small) values of $\tan\beta$. Failure of including these corrections in future studies of top quark physics at the Tevatron, both in the production and decay mechanisms, might seriously hamper the possibility of discovering clear-cut traces of SUSY physics from the identification of large non-SM quantum corrections in these processes. As already stated, we have mainly concentrated on the impact of these quantum signatures in the physics of the Tevatron, but important effects are also expected [13] for experiments aiming at the production and decay of "obese" squarks at the LHC. The latter type of squarks could be free of some of the restrictions that have been considered for the present calculation.

Acknowledgements:
Financial support of J.S. by the Spanish Ministerio de Educación y Ciencia and of W.H by the DAAD through the Programa de Acciones Integradas (Acción No. HA1996-0030) is gratefully acknowledged. The work of JG has been financed by a grant of the Comissionat per a Universitats i Recerca, Generalitat de Catalunya. J.S. and J.G. have also been

partially supported by CICYT under project No. AEN93-0474.

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Figure Captions

- **Fig.1** SUSY-EW Feynman diagrams, at one-loop order, correcting the partial width of $\tilde{b}_a \to \chi_1^- t$: **(a)** Self-energy diagrams; **(b)** Vertex contributions. There are other possible diagrams, but they do not contribute in the Yukawa-coupling approximation.

- **Fig.2 (a)** The branching ratio of $\tilde{b}_a \to \chi_1^- t$ as a function of $\tan\beta$ for the various decays $a = 1, 2$ with $m_{\tilde{b}_1} < m_{\tilde{b}_2}$; **(b)** As in (a), but as a function of $m_{\tilde{b}_1}$. The marked parts of the abscissa in both figures are excluded by the condition (23). The fixed parameters for (a) and (b) are given in the frame.

- **Fig.3 (a)** The SUSY-EW corrections (15) as a function of $m_{\tilde{t}_1}$; **(b)** As in (a), but as a function of $m_{\tilde{b}_1}$. Rest of inputs as in Fig.2.

- **Fig.4 (a)** Evolution of the SUSY-EW corrections as a function of the sbottom mixing angle, $\theta_{\tilde{b}}$, within its allowed range; **(b)** As in (a), but as a function of the stop mixing angle, $\theta_{\tilde{t}}$. Remaining inputs are as in Fig.2

- **Fig.5 (a)** The SUSY-EW correction as a function of $\tan\beta$; **(b)** As in (a), but as a function of μ. Rest of inputs and notation as in Fig.2.

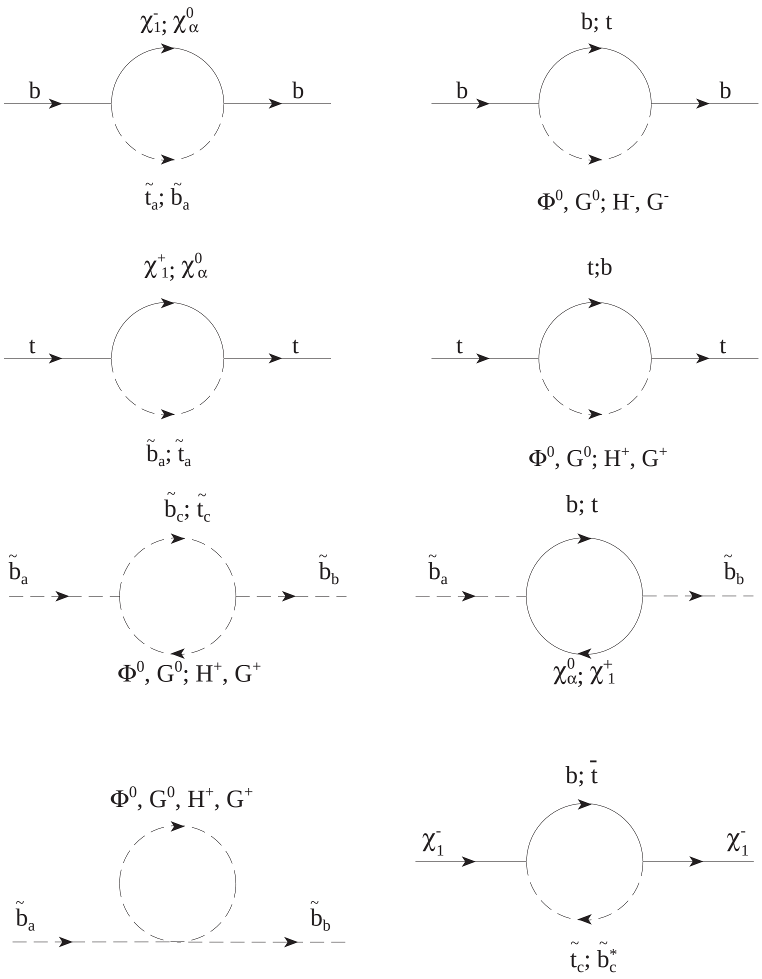

χ^-_1; χ^0_α
b
b
$\tilde{t}_a$; $\tilde{b}_a$
b; t
b
b
Φ^0, G^0; H^-, G^-
χ^+_1; χ^0_α
t
t
$\tilde{b}_a$; $\tilde{t}_a$
t;b
t
t
Φ^0, G^0; H^+, G^+
$\tilde{b}_c$; $\tilde{t}_c$
$\tilde{b}_a$
$\tilde{b}_b$
Φ^0, G^0; H^+, G^+
b; t
$\tilde{b}_a$
$\tilde{b}_b$
χ^0_α; χ^+_1
Φ^0, G^0, H^+, G^+
$\tilde{b}_a$
$\tilde{b}_b$
b; $\bar{t}$
χ^-_1
χ^-_1
$\tilde{t}_c$; $\tilde{b}^*_c$

Fig. 1 (a)

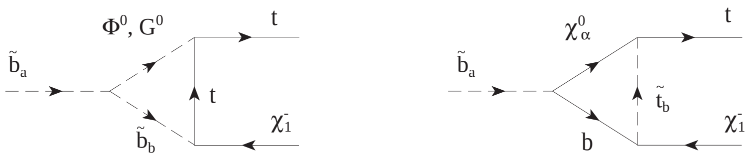
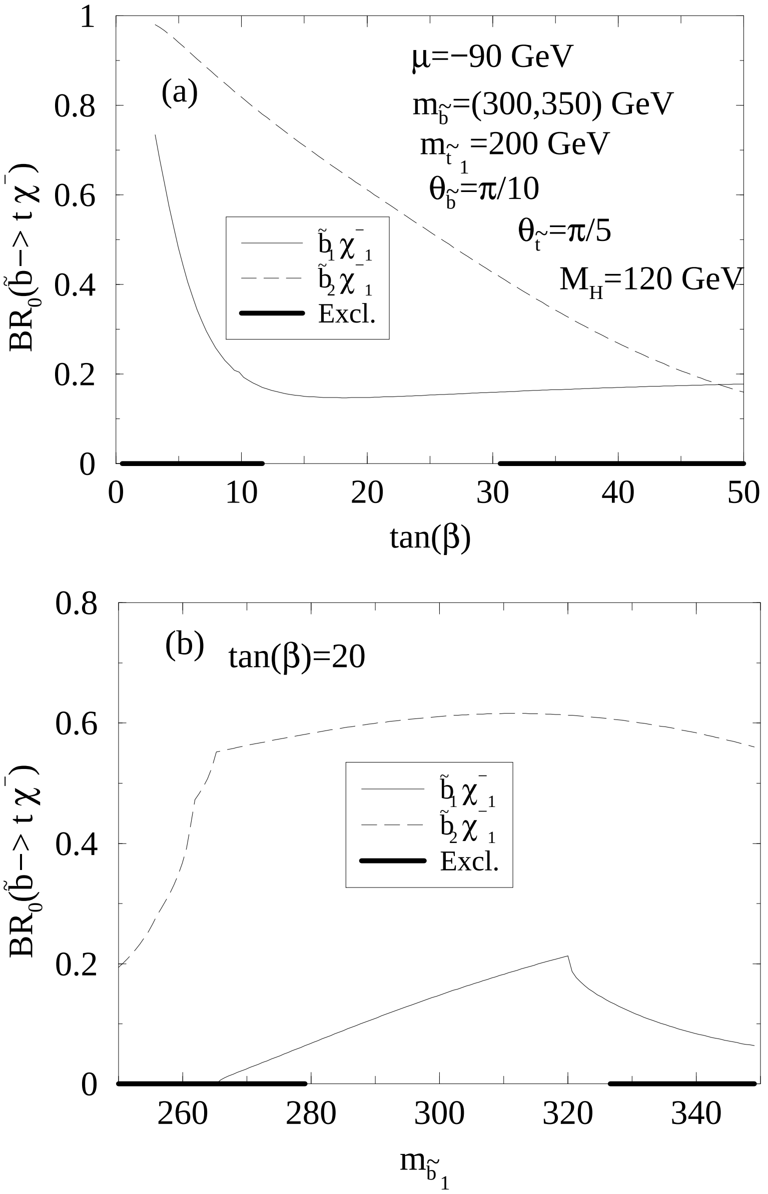

Φ^0, G^0
t
$\tilde{b}_a$
t
$\tilde{b}_b$
χ^-_1
χ^0_α
t
$\tilde{b}_a$
$\tilde{t}_b$
b
χ^-_1

Fig. 1 (b)



(a)
μ=−90 GeV
m_b̃=(300,350) GeV
m_t̃1 =200 GeV
θ_b̃=π/10
θ_t̃=π/5
M_H=120 GeV
b̃1 χ−1
b̃2 χ−1
Excl.
BR0(b̃−> t χ−)
tan(β)
(b)
tan(β)=20
b̃1 χ−1
b̃2 χ−1
Excl.
BR0(b̃−> t χ−)
m_b̃1

Fig. 2

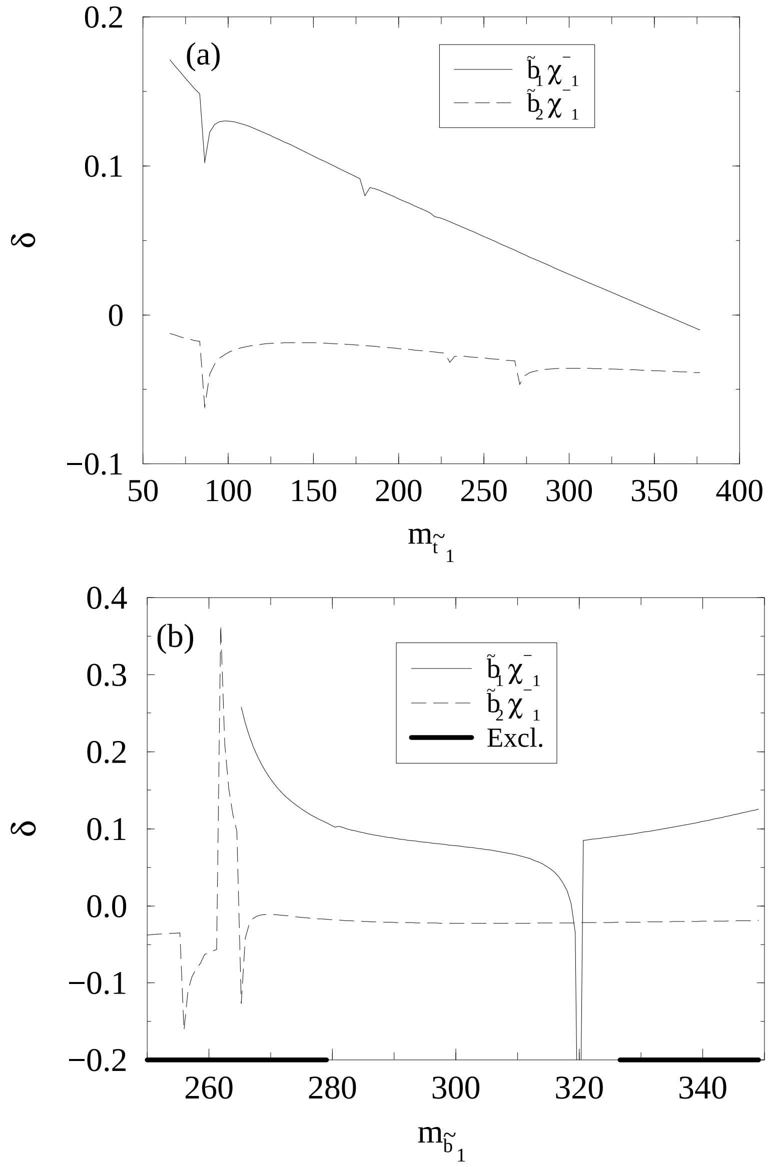
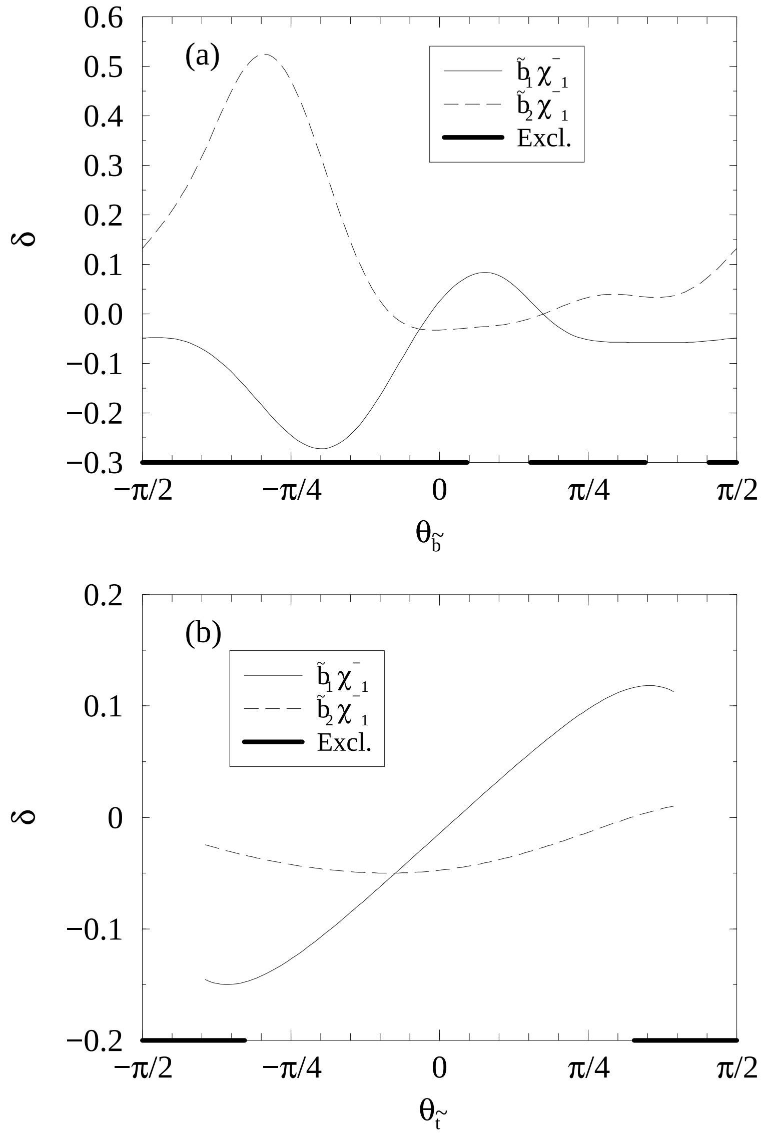

(a)
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
δ
0.2
0.1
0
−0.1
50
100
150
200
250
300
350
400
$m_{\tilde{t}_1}$
(b)
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
Excl.
δ
0.4
0.3
0.2
0.1
0.0
−0.1
−0.2
260
280
300
320
340
$m_{\tilde{b}_1}$

Fig. 3



(a)
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
Excl.
δ
0.6
0.5
0.4
0.3
0.2
0.1
0.0
−0.1
−0.2
−0.3
−π/2
−π/4
0
π/4
π/2
$\theta_{\tilde{b}}$
(b)
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
Excl.
δ
0.2
0.1
0
−0.1
−0.2
−π/2
−π/4
0
π/4
π/2
$\theta_{\tilde{t}}$

Fig. 4

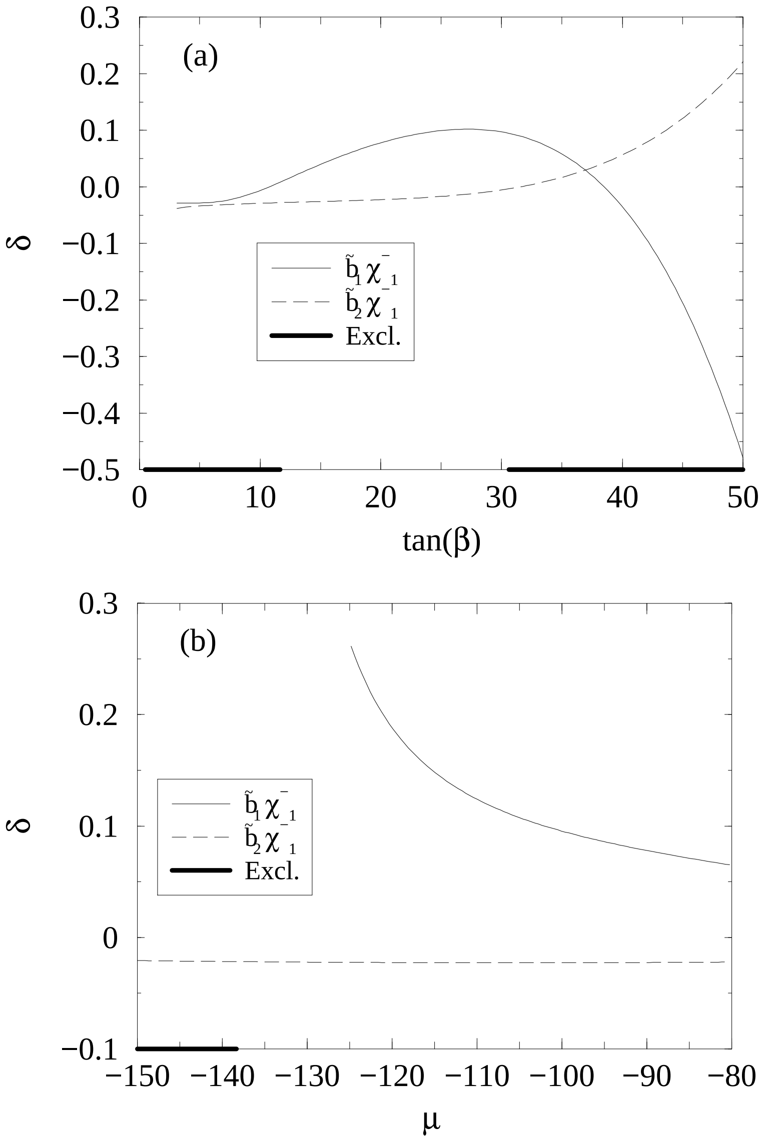

(a)
0.3
0.2
0.1
0.0
−0.1
−0.2
−0.3
−0.4
−0.5
0
10
20
30
40
50
tan(β)
δ
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
Excl.
(b)
0.3
0.2
0.1
0
−0.1
−150
−140
−130
−120
−110
−100
−90
−80
μ
δ
$\tilde{b}_1 \chi^-_1$
$\tilde{b}_2 \chi^-_1$
Excl.

Fig. 5