

SUPERSYMMETRIC QCD CORRECTIONS TO THE CHARGED HIGGS BOSON DECAY OF THE TOP QUARK

Jaume GUASCH, Ricardo A. JIMÉNEZ, Joan SOLÀ

Grup de Física Teòrica

and

Institut de Física d'Altes Energies

Universitat Autònoma de Barcelona

08193 Bellaterra (Barcelona), Catalonia, Spain

ABSTRACT

The one-loop supersymmetric QCD quantum effects on the width of the unconventional top quark decay mode $t \rightarrow H^+ b$ are evaluated within the MSSM. The study of this process is useful to hint at the supersymmetric nature of the charged Higgs emerging from that decay. Remarkably enough, recent calculations of supersymmetric corrections to Z -boson observables have shown that the particular conditions by which the decay $t \rightarrow H^+ b$ becomes competitive with the standard decay $t \rightarrow W^+ b$ have a chance to be realized in nature. This further motivates us to focus our attention on the dynamics of $t \rightarrow H^+ b$ as an excellent laboratory to unravel Supersymmetry at the quantum level in future experiments at Tevatron and at LHC.

The recent discovery of the top quark at Tevatron [1, 2] is reckoned to be the latest major event in the world of elementary particle physics. The weighted average of the *CDF* and *D0* measurements yields a mass, $m_t = 180 \pm 12 \text{ GeV}$, which is in good agreement with the Standard Model (SM) global analyses of the electroweak precision data. It can also be used to sharpen the prediction of the SM Higgs mass, M_H , in a range centered at $\sim 100 \text{ GeV}$ or below [3]. Far from being a final confirmation of the SM, the finding of the top quark raises many questions on the nature of the spontaneous symmetry-breaking mechanism (SSB) that go beyond the SM, in particular whether the SSB is caused by fundamental scalars or by some dynamical mechanism postulating a new species of strongly-interacting fermions (such as in technicolour like models [4]), or perhaps involving the condensation of the top quark itself (such as in topcolour models [5]). In this paper we shall adhere to the extensions of the SM associated with elementary scalars, specifically to the Minimal Supersymmetric Standard Model MSSM [6]. The global fit analyses of precision data within the MSSM chart an increasing trend of compatibility with the MSSM [7]. They lead to a top quark mass, $m_t = 165 \pm 10 \text{ GeV}$, consistent with the experimental measurements. Moreover, the richer structure of the top and Higgs sector of the MSSM gives rise to a very stimulating top/stop-Higgs/higgsino dynamics. The archetype example of it could be the non-standard decay of the top quark into a charged Higgs: $t \rightarrow H^+ b$. Although there are in principle many new exotic decays of the top quark in the MSSM, perhaps the latter mode is (if open) the closest one to the canonical mode $t \rightarrow W^+ b$ in the SM and therefore the less difficult to handle from the experimental point of view.

Apart from its obvious interest on its own, a further motivation [8] to consider the analysis of the charged Higgs decay of the top quark stems from the recent results of *Z* physics, particularly of the observed anomaly in the ratios R_b , R_c as well as in the value of $\alpha_s(M_Z)$ [9]. These anomalies are discrepancies (at the $2 - 3\sigma$ level) between the measured values of these observables and the corresponding predictions of the SM. As shown by early calculations of supersymmetric (SUSY) radiative corrections to R_b [10], the theoretical prediction of this observable can be made in better agreement with experiment. More recently, there has been a flood of renewed interest in this subject (see e.g. [11]-[15]) and it has been possible to establish in a more precise way the particular conditions by which the MSSM is able to cure, or at least to alleviate, the “ R_b crisis”. The corresponding MSSM analysis of the ratio R_c (in correlation with the ratio R_b) was first presented in Ref.[14]. The upshot is that both the R_b and R_c “crises” can be solved within the MSSM provided that $\tan\beta$ is large enough ¹ and there exists a supersymmetric pseudoscalar

¹A possible solution also exists for $\tan\beta < 1$, but it is not favoured by model building [16] and it is not relevant to the present analysis.

Higgs as well as some superpartner all of them in the 50 GeV range –in agreement with the most recent global fit analyses [7]. In these conditions it turns out that one can simultaneously provide a supersymmetric explanation [17] of the longstanding mismatch between the low-energy and high-energy determinations of $\alpha_s(M_Z)$ [18]. All in all, these results bring in an independent incentive from the high precision world of Z physics to choose our SUSY parameters in the region of large $\tan\beta$ and moderate charged Higgs mass. Indeed, from the well-known Higgs mass relations in the MSSM [19] and assuming that a light CP-odd (“pseudoscalar”) Higgs mass m_{A^0} exists in the 50 GeV ballpark², it follows that there must be a charged Higgs companion of $M_{H^\pm} \simeq 100\text{ GeV}$. In these circumstances the decay $t \rightarrow H^+ b$ becomes competitive with the SM decay $t \rightarrow W^+ b$ and it should be possible to identify it by tagging violations of lepton universality caused by the presence of an excess of final state τ -leptons associated to the subsequent Higgs decay –as commented at the end of Ref.[14].

In view of the potential interest of the decay mode $t \rightarrow H^+ b$, one would naturally like to address the computation of the strong virtual corrections to its partial width. Of these, the conventional QCD corrections have already been considered in detail in Ref.[20] and they turn out to be sizeable and negative (of order -10%). Although they are blind to the nature of the underlying Higgs model, they need to be subtracted from the experimentally measured number in order to be able to probe the existence of new sources of quantum effects beyond the SM. These effects may ultimately reveal whether the charged Higgs emerging from that decay is supersymmetric or not.

In this paper we compute the strong SUSY radiative corrections to the partial width by paying special attention to the aforementioned privileged region of the MSSM parameter space. The analysis of the larger and more complex body of SUSY electroweak corrections, namely the corrections mediated by squarks, sleptons, chargino-neutralinos and the Higgs bosons themselves, will be presented elsewhere [21]³. To compute the one-loop QCD corrections to $\Gamma_t \equiv \Gamma(t \rightarrow H^+ b)$ in the MSSM, we shall adopt the on-shell renormalization scheme where the fine structure constant, α , and the masses of the gauge bosons, fermions and scalars are the renormalized parameters: $(\alpha, M_W, M_Z, M_H, m_f, M_{SUSY}, \dots)$ [25]. The interaction Lagrangian describing the $t b H^\pm$ -vertex in the MSSM reads as follows:

$$\mathcal{L}_{Hbt} = \frac{g V_{tb}}{\sqrt{2}M_W} H^- \bar{b} [m_t \cot\beta P_R + m_b \tan\beta P_L] t + \text{h.c.}, \quad (1)$$

where $P_{L,R} = 1/2(1 \mp \gamma_5)$ are the chiral projector operators, $\tan\beta$ is the ratio between

²It is noteworthy that at high $\tan\beta$ the approximate phenomenological lower limit $m_{h^0} \gtrsim 50\text{ GeV}$ on the light CP-even Higgs mass of the MSSM translates into $m_{A^0} \gtrsim 50\text{ GeV}$ for the CP-odd mass. Recent global fit analyses also favour a light Higgs mass in the SM and a light CP-even Higgs mass in the MSSM [3].

³The study of the corresponding supersymmetric quantum corrections to the canonical decay $t \rightarrow W^+ b$ has been presented by some of the authors in Refs.[22] and [23] (See also Ref.[24]).

the vacuum expectation values of the two Higgs doublets of the MSSM [6] and V_{tb} is the corresponding Kobayashi-Maskawa matrix element—hereafter we set $V_{tb} = 1$ ($V_{tb} = 0.999$ within $\pm 0.1\%$, from unitarity of the KM-matrix *and* assuming 3 quark families).

There are no oblique strong supersymmetric corrections at 1-loop order. The non-oblique vertex corrections originating from gluinos and squarks (stop and sbottom species) are depicted in Fig.1. The SUSY-QCD interaction Lagrangian relevant to our calculation is given, in four-component notation, by

$$\mathcal{L} = -\frac{g_s}{\sqrt{2}} \left[\tilde{q}_L^{i*} (\lambda_r)_{ij} \tilde{g}^r P_L q^j - \tilde{q}^i (\lambda_r)_{ij} P_L \tilde{g}^r \tilde{q}_R^j \right] + \text{h.c.}, \quad (2)$$

where \tilde{g}^r ($r = 1, 2, \dots, 8$) are the Majorana gluino fields, $(\lambda_r)_{ij}$ ($i, j = 1, 2, 3$) are the Gell-Mann matrices, and $\tilde{q}'_a = \{\tilde{q}_L, \tilde{q}_R\}$ are the weak-eigenstate squarks associated to the two chiral components $P_{L,R} q$; they are related to the corresponding mass-eigenstates $\tilde{q}_a = \{\tilde{q}_1, \tilde{q}_2\}$ by a rotation 2×2 matrix (we neglect intergenerational mixing):

$$\begin{aligned} \tilde{q}'_a &= \sum_b R_{ab}^{(q)} \tilde{q}_b, \\ R^{(q)} &= \begin{pmatrix} \cos \theta_q & \sin \theta_q \\ -\sin \theta_q & \cos \theta_q \end{pmatrix} \quad (q = t, b). \end{aligned} \quad (3)$$

These rotation matrices diagonalize the corresponding stop and sbottom mass matrices:

$$\mathcal{M}_t^2 = \begin{pmatrix} M_{t_L}^2 + m_t^2 + \cos 2\beta \left(\frac{1}{2} - \frac{2}{3} s_W^2 \right) M_Z^2 & m_t M_{LR}^t \\ m_t M_{LR}^t & M_{t_R}^2 + m_t^2 + \frac{2}{3} \cos 2\beta s_W^2 M_Z^2 \end{pmatrix}, \quad (4)$$

$$\mathcal{M}_b^2 = \begin{pmatrix} M_{b_L}^2 + m_b^2 + \cos 2\beta \left(-\frac{1}{2} + \frac{1}{3} s_W^2 \right) M_Z^2 & m_b M_{LR}^b \\ m_b M_{LR}^b & M_{b_R}^2 + m_b^2 - \frac{1}{3} \cos 2\beta s_W^2 M_Z^2 \end{pmatrix}, \quad (5)$$

with

$$M_{LR}^t = A_t - \mu \cot \beta, \quad M_{LR}^b = A_b - \mu \tan \beta, \quad (6)$$

μ being the SUSY Higgs mass parameter in the superpotential⁴. The $A_{t,b}$ are the trilinear soft SUSY-breaking parameters and the $M_{\tilde{q}_{L,R}}$ are soft SUSY-breaking masses [6]. By $SU(2)_L$ -gauge invariance we must have $M_{\tilde{t}_L} = M_{\tilde{b}_L}$, whereas $M_{\tilde{t}_R}$, $M_{\tilde{b}_R}$ are in general independent parameters. Finally, we also need the interaction Lagrangian involving the charged Higgs and the stop and sbottom squarks

$$\mathcal{L}_{H\tilde{t}\tilde{t}} = -\frac{g}{\sqrt{2} M_W} H^- \left(g_{LL} \tilde{b}_L^* \tilde{t}_L + g_{RR} \tilde{b}_R^* \tilde{t}_R + g_{LR} \tilde{b}_R^* \tilde{t}_L + g_{RL} \tilde{b}_L^* \tilde{t}_R \right) + \text{h.c.}, \quad (7)$$

where

$$\begin{aligned} g_{LL} &= M_W^2 \sin 2\beta - (m_t^2 \cot \beta + m_b^2 \tan \beta), \\ g_{RR} &= -m_t m_b (\tan \beta + \cot \beta), \\ g_{LR} &= -m_b (\mu + A_b \tan \beta), \\ g_{RL} &= -m_t (\mu + A_t \cot \beta). \end{aligned} \quad (8)$$

⁴Its sign is relevant in the numerical analysis. We fix it as in eq.(3) of Ref.[22].

The one-loop renormalized vertex, Λ , is derived from the renormalized Lagrangian plus counterterms, $\mathcal{L} \rightarrow \mathcal{L} + \delta\mathcal{L}$, following the standard procedure [25]. It can be parametrized in terms of two form factors F_L, F_R and the corresponding mass and wave-function renormalization constants $\delta m_f, \delta Z_{L,R}^f$ associated to the external quarks, viz.

$$\Lambda = \frac{ig}{\sqrt{2}M_W} [m_t \cot\beta (1 + \Lambda_R) P_R + m_b \tan\beta (1 + \Lambda_L) P_L], \quad (9)$$

with

$$\begin{aligned} \Lambda_R &= F_R + \frac{\delta m_t}{m_t} + \frac{1}{2} \delta Z_L^b + \frac{1}{2} \delta Z_R^t, \\ \Lambda_L &= F_L + \frac{\delta m_b}{m_b} + \frac{1}{2} \delta Z_L^t + \frac{1}{2} \delta Z_R^b. \end{aligned} \quad (10)$$

In the on-shell scheme we have⁵

$$\frac{\delta m_q}{m_q} = - \left[\frac{\Sigma_L^q(m_q^2) + \Sigma_R^q(m_q^2)}{2} + \Sigma_S^q(m_q^2) \right] \quad (11)$$

and

$$\delta Z_{L,R}^q = \Sigma_L^q(m_f^2) + m_q^2 [\Sigma_L^{q'}(m_q^2) + \Sigma_R^{q'}(m_q^2) + 2\Sigma_S^{q'}(m_q^2)]. \quad (12)$$

In these equations we have decomposed the (real part of the) quark self-energy according to

$$\Sigma^f(p) = \Sigma_L^f(p^2) \not{p} P_L + \Sigma_R^f(p^2) \not{p} P_R + m_f \Sigma_S^f(p^2), \quad (13)$$

and used the notation $\Sigma'(p) \equiv \partial\Sigma(p)/\partial p^2$.

From the renormalized amplitude (9), the width $\Gamma = \Gamma(t \rightarrow H^+ b)$ including the one-loop SUSY-QCD corrections is the following:

$$\Gamma = \Gamma_0 \left\{ 1 + \frac{N_L}{N} [2 \operatorname{Re}(\Lambda_L)] + \frac{N_R}{N} [2 \operatorname{Re}(\Lambda_R)] + \frac{N_{LR}}{N} [2 \operatorname{Re}(\Lambda_L + \Lambda_R)] \right\}, \quad (14)$$

where the corresponding lowest-order result is

$$\Gamma_0 = \left(\frac{G_F}{8\pi\sqrt{2}} \right) \frac{N}{m_t} \lambda^{1/2} \left(1, \frac{m_b^2}{m_t^2}, \frac{M_{H^+}^2}{m_t^2} \right). \quad (15)$$

We have defined

$$\lambda^{1/2}(1, x^2, y^2) \equiv \sqrt{[1 - (x+y)^2][1 - (x-y)^2]} \quad (16)$$

and

$$\begin{aligned} N &= (m_t^2 + m_b^2 - M_{H^+}^2) (m_t^2 \cot^2\beta + m_b^2 \tan^2\beta) + 4m_t^2 m_b^2, \\ N_L &= (m_t^2 + m_b^2 - M_{H^+}^2) m_b^2 \tan^2\beta, \\ N_R &= (m_t^2 + m_b^2 - M_{H^+}^2) m_t^2 \cot^2\beta, \\ N_{LR} &= 2m_t^2 m_b^2. \end{aligned} \quad (17)$$

⁵Our sign conventions for the self-energy functions are those of Ref.[22].

Notice that, in contradistinction to electroweak one-loop calculations [22], an additional correction term Δr does not appear in (14) due to the absence of one-loop SUSY-QCD corrections in μ -decay.

The explicit contribution to the form factors F_L, F_R from the vertex diagram of Fig.1 is given by

$$\begin{aligned} F_L &= 8\pi\alpha_s iC_F \frac{G_{ab}}{m_b \tan\beta} [R_{2b}^{(t)*} R_{2a}^{(b)} (C_{11} - C_{12}) m_t + R_{1b}^{(t)*} R_{1a}^{(b)} C_{12} m_b + R_{1b}^{(t)*} R_{2a}^{(b)} C_0 m_{\tilde{g}}], \\ F_R &= 8\pi\alpha_s iC_F \frac{G_{ab}}{m_t \cot\beta} [R_{1b}^{(t)*} R_{1a}^{(b)} (C_{11} - C_{12}) m_t + R_{2b}^{(t)*} R_{2a}^{(b)} C_{12} m_b + R_{2b}^{(t)*} R_{1a}^{(b)} C_0 m_{\tilde{g}}]. \end{aligned} \quad (18)$$

Here $C_F = (N_C^2 - 1)/2N_C = 4/3$ is a colour factor. We have furthermore defined:

$$G_{ab} = R_{1b}^{(t)} R_{1a}^{(b)*} g_{LL} + R_{2b}^{(t)} R_{2a}^{(b)*} g_{RR} + R_{1b}^{(t)} R_{2a}^{(b)*} g_{LR} + R_{2b}^{(t)} R_{1a}^{(b)*} g_{RL}. \quad (19)$$

The three-point function notation is as in Ref.[22, 27], with the following arguments:

$$C = C(p, p', m_{\tilde{g}}, m_{\tilde{t}_b}, m_{\tilde{t}_a}). \quad (20)$$

As for the self-energies,

$$\begin{aligned} \Sigma_L^q(p^2) &= 8\pi\alpha_s iC_F |R_{1a}^{(q)}|^2 (B_0 - B_1), \\ \Sigma_R^q(p^2) &= 8\pi\alpha_s iC_F |R_{2a}^{(q)}|^2 (B_0 - B_1), \\ \Sigma_S^q(p^2) &= -8\pi\alpha_s iC_F \frac{m_{\tilde{g}}}{m_q} \text{Re}(R_{1a}^{(q)} R_{2a}^{(q)*}) B_0, \end{aligned} \quad (21)$$

where the two-point functions –defined also as in Ref.[22]– have the following arguments:

$$B = B(p^2, m_{\tilde{q}_a}^2, m_{\tilde{g}}^2). \quad (22)$$

(A summation over squark indices is understood in eqs.(18) and (21).) It is easy to convince oneself that the form factors (18) are to be UV -finite in SUSY-QCD, as indeed they are. The remaining contributions to the renormalized amplitude (9) are immediately seen to cancel UV -divergences each other out.

The numerical analysis of the strong supersymmetric corrections to $\Gamma(t \rightarrow H^+ b)$ is exhibited in Figs.2-5. We present the results both in terms of the corrected Γ and in terms of the relative correction with respect to the tree-level width, i.e.

$$\delta_{\tilde{g}} = \frac{\Gamma - \Gamma_0}{\Gamma_0}. \quad (23)$$

The free parameters at our disposal lie in the mass matrices (4)-(5). Among the sfermion mass matrices, the stop mass matrix is the only one where a non-diagonal structure caused

by a sizeable mixing term is most likely to arise. For $M_{LR}^t = 0$ this matrix is trivial, but since m_t is large a nonvanishing M_{LR}^t naturally leads to a light mass eigenvalue, denoted by $m_{\tilde{t}_1}$, whereas the other eigenvalue, $m_{\tilde{t}_2}$, can be much heavier. As a matter of fact a light stop with a mass $m_{\tilde{t}_1} = \mathcal{O}(M_Z/2)$ is still phenomenologically allowed [26]. In contrast, the off-diagonal element of the sbottom mass matrix (5), being proportional to $m_b \simeq 4.5 \text{ GeV}$, is expected to be small (unless M_{LR}^b is very large). We shall treat the sbottom mass matrix in the simplest possible way compatible with the phenomenological bounds on squark masses [26]—only scaped perhaps by the lightest stop. For definitiveness, unless stated otherwise, we shall assume that $M_{LR}^b = 0$ (equivalently $A_b = \mu \tan \beta$) and that the two mass eigenvalues are equal ($m_{\tilde{b}_1} = m_{\tilde{b}_2} \equiv m_{\tilde{b}}$) and constrained to satisfy $m_{\tilde{b}} \geq 150 \text{ GeV}$. As for A_t , it will be treated either as an input parameter or it will be fixed once we are given M_{LR}^t , μ and $\tan \beta$. We remind that M_{LR}^t is expected to preserve the inequality

$$M_{LR}^t \leq 3 m_{\tilde{b}_L}, \quad (24)$$

which roughly corresponds to a necessary, though not sufficient, condition to avoid colour-breaking vacua [28]. In the conditions described above, once $m_{\tilde{b}}$, μ and $\tan \beta$ are fixed, the stop mass matrix depends on only 2 parameters, e.g. $(A_t, M_{\tilde{t}_R}^t)$, $(M_{LR}^t, m_{\tilde{t}_1})$, etc. For the strong coupling constant we used the value

$$\alpha_s = \frac{g_s^2}{4\pi} = 0.11 \quad (25)$$

which remains essentially constant within the CDF-D0 ranges mentioned above. Whenever m_t needs to be fixed, we take the central value $m_t = 180 \text{ GeV}$.

A crucial parameter to be explored in our analysis is $\tan \beta$. In Fig.2a we plot the SUSY-QCD corrected $\Gamma = \Gamma(t \rightarrow H^+ b)$, eq.(14), versus $\tan \beta$ for $\mu = +100 \text{ GeV}$ and $\mu = -100 \text{ GeV}$, and for given values of the other parameters. We see that Γ is very sensitive to $\tan \beta$ and that, barring the narrow interval $\tan \beta \leq 1$, the process $t \rightarrow H^+ b$ becomes steadily competitive with the standard process, $t \rightarrow W^+ b$, in the large $\tan \beta$ region, i.e. when $\tan \beta$ is of the order of $30 - 40 \simeq m_t/m_b$. As mentioned in the beginning, this is precisely the range singled out by the Z boson observables [14]. We also see from Fig.2a that, for $\tan \beta \gtrsim 30$, Γ starts to deviate from Γ_0 quite manifestly. Therefore we shall choose $\tan \beta = 30$ as a representative value in the other plots.

Highly remarkable is also the incidence of the parameter μ both of its value and of its sign. In fact, the sign of $\delta_{\tilde{g}}$ happens to be opposite to the sign of μ and the respective corrections for μ and for $-\mu$ take on approximately the same absolute value. In Fig.2b we deliver the correction itself, $\delta_{\tilde{g}}$, for $\mu = -100 \text{ GeV}$ and for different values of the squark and gluino masses. The sign dependence of $\delta_{\tilde{g}}$ suggests that two extreme scenarios could take place with the SUSY-QCD corrections to $\Gamma(t \rightarrow H^+ b)$: namely, they could either

significantly enhance the, negative, conventional QCD corrections [20], or on the contrary they could counterbalance them and even result in opposite sign.

In Figs.3a-3b we display $\delta_{\tilde{g}}$ as a function of $m_{\tilde{g}}$ and m_t , respectively, for three values of the sbottom masses. We learn from Fig.3a that light gluinos of $\mathcal{O}(1) GeV$ [29] yield, contrary to naive expectations, a rather small correction as compared to heavy gluinos of $\mathcal{O}(100) GeV$. It should be clear that these corrections do eventually decouple –as we have checked– for larger and larger gluino masses. Notwithstanding, the decoupling rate of the gluinos is particularly noticeable, for it happens to be so slow (Fig.3a) that it fakes for a while a non-decoupling behaviour. This trait is caused by the presence of a long sustained local maximum (or minimum, depending on the sign of μ) spreading over a wide range of heavy gluino masses centered at $\sim 300 GeV$. For this reason, heavy gluinos are in the present instance preferred to light gluinos. In Fig.4, we test the sensitivity of $\delta_{\tilde{g}}$ to M_{H^\pm} ; it turns out to be very small, except near the uninteresting vicinity of the phase space border where Γ_0 is about to vanish.

In all previous figures we have fixed $A_t = 0$, $\theta_t = \pi/4$ and $m_{\tilde{b}} \leq 150 GeV$, so that the lightest stop mass was always $m_{\tilde{t}_1} > 190 GeV$. In Fig.5 we relax these conditions and plot contour isolines $\delta_{\tilde{g}} = \text{const.}$ in the $(M_{LR}^t, m_{\tilde{t}_1})$ -plane, where A_t and θ_t are variable. In particular, we approach the region of the lightest possible stop masses compatible with the strict LEP phenomenological bound. As expected, the corrections are larger the smaller is $m_{\tilde{t}_1}$. Finally, let us mention that a non-vanishing mixing in the sbottom mass matrix does not alter at all the typical size of the corrections obtained here. It was only to avoid much cluttering of free parameters that we have treated so far that matrix in the most simple-minded form ($M_{LR}^b = 0$ and $m_{\tilde{b}_1} = m_{\tilde{b}_2}$). For example, if the mass-eigenvalues are $m_{\tilde{b}_{1,2}} = 200, 250 GeV$ and $\mu = -100 GeV$, $\tan \beta = 30$, then one has $A_b = -500 GeV$ and the corresponding correction reads $\delta_{\tilde{g}} = +30\%$. In the same conditions, but choosing $A_b = 0$, μ is determined to be $-83 GeV$ and $\delta_{\tilde{g}} = +25\%$, etc.

Some words on previous work are in order. We chart significant differences in our complete analysis as compared to preliminary calculations in the literature. In Ref.[30] a first study of the SUSY-QCD corrections to $t \rightarrow H^+ b$ was presented, but they neglect the bottom quark Yukawa coupling and as a consequence they are incorrectly sensitive to the high $\tan \beta$ effects. Indeed, the form factor F_L which is associated to that coupling is by far the dominant piece of our numerical analysis in the large $\tan \beta$ region. Furthermore, the impact from mixing effects and the incidence of the various parameter dependences were completely missed and only the simplest situation, characterized by degenerate masses, was considered⁶. The study of Ref.[31] also neglects the bottom quark Yukawa coupling. Thus the purported large effects claimed in the large $\tan \beta$ region are not correctly justi-

⁶Notice that the assumption of stop masses equal to sbottom masses is incompatible with eqs.(4)-(5).

fied. Moreover, in the framework of these two references, the lowest-order width is fully proportional to $\cot\beta$; thus, in such a context, finding quantum effects increasing with $\tan\beta$ is rather useless since they result in corrections to an uninteresting, vanishingly small, tree-level width. Finally, the latter reference also neglects the wave-function and mass renormalization contributions which, numerically, are of the same order of magnitude as the vertex contributions in the intermediate and small $\tan\beta$ region. However, where it has been possible to force an overlapping, we have found numerical coincidence.

In summary, the SUSY-QCD contributions to the partial width of $t \rightarrow H^+ b$ could be quite large (several 10%); and what is more, these corrections apply to a decay mode which, in the region of the MSSM parameter space prompted by the high precision Z boson observables [14, 17], has an appreciable branching ratio as compared to the standard decay $t \rightarrow W^+ b$. Furthermore, we have found that the impact of the gluino corrections on $t \rightarrow H^+ b$ could occur in two opposite ways, either by reinforcing the conventional QCD corrections or on the contrary by severely cancelling them out—perhaps even to the extent of reversing their sign!. Most remarkable, the potential size of these effects stems not only from the strong interaction character of the SUSY-QCD corrections, but also from the high sensitivity of $t \rightarrow H^+ b$ to the (weak-interaction) SSB parameter $\tan\beta$. At the end of the day we must conclude that $t \rightarrow H^+ b$ could reveal itself as the ideal environment where to study the nature of the SSB mechanism. It could even be the right place where to target our long and unsuccessful search for large, *and* slowly decoupling, quantum supersymmetric effects. In this respect it should not be understated the fact that the typical size of our corrections is maintained even for sparticle masses well above the LEP 200 discovery range. These features are in stark contrast to the standard decay of the top quark, $t \rightarrow W^+ b$, whose SUSY-QCD corrections are largely insensitive to $\tan\beta$ [23]. Fortunately, the next generation of experiments at Tevatron and the future high precision experiments at LHC may well acquire the ability to test the kind of effects considered here (Cf. Refs. [32, 33, 34]). Thus, in favorable circumstances, we should be able to disentangle the potential supersymmetric nature of the charged Higgs decay of the top quark out of a measurement of the top quark width⁷ at a modest precision of $\sim 5 - 10\%$.

⁷Or related observables, such as e.g. the –double and single– top quark production cross-sections, or the differential distributions of the lepton final states in given exclusive decay channels, etc.

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Figure Captions

- **Fig.1** SUSY-QCD Feynman diagrams, up to one-loop order, correcting the partial width $\Gamma(t \rightarrow H^+ b)$. Each one-loop diagram is summed over the mass-eigenstates of the stop, sbottom squarks $(\tilde{b}_a, \tilde{t}_b; a, b = 1, 2)$ and gluinos $\tilde{g}_r; r = 1, 2, \dots, 8$.
- **Fig.2** (a) SUSY-QCD corrected $\Gamma(t \rightarrow H^+ b)$ as a function of $\tan \beta$, for two opposite values of μ , compared to the corresponding tree-level width, Γ_0 . The framed set of inputs is common to Figs. 2,3. The horizontal (dashed) line marks $\Gamma_0(t \rightarrow W^+ b) = 1.71 \text{ GeV}$ –the tree-level width of the standard process ($m_t = 180 \text{ GeV}$); its SUSY-QCD corrections are generally small and essentially independent of $\tan \beta$ [23]; (b) The relative correction $\delta_{\tilde{g}}$ as a function of $\tan \beta$ for $\mu = -100 \text{ GeV}$ and three values of $m_{\tilde{b}} = m_{\tilde{g}}$.
- **Fig.3** Dependence of $\delta_{\tilde{g}}$ upon (a) $m_{\tilde{g}}$, including the light gluino region, for the same squark masses and μ as in Fig.2b; (b) $\delta_{\tilde{g}}$ as a function of m_t (within the *CDF-D0* limits) and remaining parameters as in (a).
- **Fig.4** $\delta_{\tilde{g}}$ as a function of M_{H^\pm} . Rest of inputs as in Fig.3.
- **Fig.5** Contour plots of $\delta_{\tilde{g}}$ in the $(M_{LR}^t, m_{\tilde{t}_1})$ -plane for $\mu = -100$ and $\tan \beta = 30$. The mixing angle θ_t and A_t are variable and the other fixed parameters are as in Fig.2a. The shaded area is excluded by the triple condition $M_{\tilde{t}_R}^2 > 0$, $m_{\tilde{t}_1} \geq M_Z/2$ and eq.(24).