THE TOP-QUARK WIDTH IN THE LIGHT OF 
Z-BOSON PHYSICS

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ABSTRACT

We discuss possible non-standard contributions to the top-quark width, particularly the virtual effects on the standard decay $t \rightarrow W^+ b$ within the context of the MSSM. We also place a renewed emphasis on the unconventional mode $t \rightarrow H^+ b$ in the light of recent analyses of $Z$-boson observables. It turns out that in the region of parameter space highlighted by $Z$-boson physics, the charged Higgs mode should exhibit an appreciable branching fraction as compared to the standard decay of the top quark. Remarkably enough, the corresponding quantum effects in this region are also rather large, slowly decoupling, and most likely resolvable in the next generation of experiments at Tevatron and at LHC.

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1Based on the talk presented at the Workshop on Physics of the Top Quark, Iowa State University, May 1995.
The recent discovery of a heavy top quark \((m_t \sim 180\, \text{GeV})\) at Tevatron \([1]\) constitutes, paradoxically as it may sound, both a reassuring confirmation of a long-standing prediction of the Standard Model (SM) of the electroweak interactions and, with no less emphasis, the consolidation of an old and intriguing suspicion, namely, that the SM cannot be the last word in elementary particle physics. Needless to say, the ultimate proof of this conjecture can only be substantiated in the experimental ring.

In the meanwhile, the projected ten-fold increase of the Tevatron luminosity via the Main Injector and Recycler facilities in combination with a \(\sim 40\%\) increase of the top-quark production cross-section at a \(2\, \text{TeV}\) running energy, as compared to the \(1.8\, \text{TeV}\) run (Run I), augur a brilliant Run II performance of this machine; perhaps also an exciting new era of top-quark physics that may explicitly reveal the long sought-after signs of newness and non-standardness. In fact, the finding of such a heavy quark poses some questions of fundamental nature that go beyond the SM, may be the most obvious one being the following: is the top quark an “abnormal” fermion or on the contrary it is the only “normal” fermion in the SM?. Supporters of the first contention may argue that the top quark is the only superheavy quark in the SM, whereas opponents may adduce that the top quark is the only quark whose Yukawa coupling with the SM Higgs boson is of the same order as the electroweak gauge coupling. Be as it may, this same question formulated in a non-SM context may result in a richer panoply of answers. For instance, in topcolour models \([2]\) the huge mass of the top quark is linked to the postulation of a new strong force responsible for the formation of a dynamical quark mass; thereupon it is not inconceivable to think of the spontaneous symmetry breaking mechanism (SSB) in the SM as a phenomenon of top-quark condensation. In more sophisticated versions, like in the so-called topcolour-assisted models \([3]\), whether technicolour \([4]\)-like or Higgs-like, one relaxes the requirement that a top condensate accounts for the full SSB and makes allowance for a hierarchy of electroweak mass scales. Thus, on the one hand, the large mass \(m_t\) is almost entirely driven by a new (non-confining) strong interaction, preferentially coupled to the third quark generation, which gives rise to a dynamical condensate naturally tilted in the top-quark direction; light fermion masses, on the other hand, are generated by the underlying (extended \([5]\), perhaps even walking \([6]\)) technicolour or by fundamental Higgs interactions.

Quite in contrast, in supersymmetric theories, like the Minimal Supersymmetric
Standard Model (MSSM)\[7\], one sticks altogether to fundamental scalars. However, the corresponding spectrum of higgses and of Yukawa couplings is far and away richer than in the SM; and, in such a framework, the bottom-quark Yukawa coupling may counterbalance the smallness of the bottom mass at the expense of a large value of \(\tan \beta\) – the ratio of the vacuum expectation values (VEV’s) of the two Higgs doublets – the upshot being that the top-quark and bottom-quark Yukawa couplings standing in the superpotential \[7\],

\[
\begin{align*}
    h_t &= \frac{g m_t}{\sqrt{2} M_W \sin \beta}, \\
    h_b &= \frac{g m_b}{\sqrt{2} M_W \cos \beta},
\end{align*}
\]

(1)
can be of the same order of magnitude, perhaps even showing up in “inverse hierarchy”: \(h_t < h_b\) for \(\tan \beta > m_t/m_b\). Notice that due to the perturbative bound \(\tan \beta \lesssim 60−70\) one never reaches a situation where \(h_t << h_b\). In a sense, \(h_t \sim h_b\) could be judged as a natural relation in the MSSM, and thus some of the criticism raised above loses its meaning since \(g \sim h_t \sim h_b\) can be made to coexist with \(m_t >> m_b\) in the MSSM even at the electroweak scale. As a matter of fact one ends up by rephrasing the same problem by posing a different question; one no longer asks about why \(m_t\) is much larger than \(m_b\), though one has to cope with a rather large value of \(\tan \beta\) that must be explained – may be by invoking some unification model that subsumes the general structure of the MSSM.

On the phenomenological side, one should not dismiss the possibility that the bottom-quark Yukawa coupling could play a momentous role in the physics of the top quark, to the extent of drastically changing standard expectations on top-quark observables, particularly on the top-quark width. In this respect it should be mentioned that the present measurements of the branching fraction of the standard decay of the top quark \[8\],

\[
\frac{\Gamma(t \rightarrow W^+ b)}{\Gamma(t \rightarrow \text{all})} = 0.87^{+0.13+0.13}_{-0.30-0.11},
\]

(2)
do still leave room enough to accommodate non-standard decays; hence we may expect additional decay modes of the top quark into bottom jets plus a new charged pseudoscalar particle subsequently disintegrating into fermion pairs. A non-standard mode that has been suggested along this line in the framework of Refs.\[2, 3\] is \(t \rightarrow \tilde{\pi}^+ b\), where \(\tilde{\pi}^+\) is a charged member of the “top-pion” triplet, one of the firmest predictions of topcolour models\[3\]. However, in this case the coupling strength is governed by a

\[\text{Top-pions are predicted to be in the top-quark mass range, so it is not clear which one of the}\]
Goldberger-Trieman type relation: \( g_{tb} \sim m_t/\sqrt{2}f_\pi \sim 2.5 \) (\( f_\pi \) being the top-pion decay constant), and as a consequence top-pion physics is basically sensitive to the top-quark mass, not to the bottom-quark mass.

In this talk we wish to emphasize the possibility that the charged pseudoscalar involved in a potential unconventional top-quark decay be the charged Higgs of the MSSM: \( t \to H^+ b \)\(^3\). In contrast to \( \tilde{\pi}^+ \), the charged Higgs can be, as noted above, very sensitive to bottom-quark interactions. Specifically, after expressing the two-doublet Higgs fields of the MSSM in terms of the corresponding mass-eigenstates, the interaction Lagrangian describing the \( t b H^\pm \)-vertex reads as follows [13]:

\[
L_{Htb} = \frac{g V_{tb}}{\sqrt{2} M_W} H^- \bar{b} [m_t \cot \beta P_R + m_b \tan \beta P_L] t + \text{h.c.} .
\] (3)

Similarly, from the \( D \)-type terms of the MSSM Lagrangian the relevant interaction vertices involving the charged Higgs and the stop and sbottom squarks take on the form

\[
L_{H\tilde{b}\tilde{b}} = -\frac{g}{\sqrt{2} M_W} H^- \left( g_{LL} \tilde{b}^*_L \tilde{t}_L + g_{RR} \tilde{b}^*_R \tilde{t}_R + g_{LR} \tilde{b}^*_R \tilde{t}_L + g_{RL} \tilde{b}^*_L \tilde{t}_R \right) + \text{h.c.} ,
\] (4)

with

\[
\begin{align*}
g_{LL} &= M_W^2 \sin 2\beta - (m_t^2 \cot \beta + m_b^2 \tan \beta) , \\
g_{RR} &= -m_t m_b (\tan \beta + \cot \beta) , \\
g_{LR} &= -m_b (\mu + A_b \tan \beta) , \\
g_{RL} &= -m_t (\mu + A_t \cot \beta) ,
\end{align*}
\] (5)

\( A_{t,b} \) being the trilinear soft SUSY-breaking parameters [4]. Notice that \( \tilde{q}'_a = \{\tilde{q}_L, \tilde{q}_R\} \) are the weak-eigenstate squarks associated to the two chiral fermion components \( P_{L,R} q = \frac{1}{2} (1 \mp \gamma_5) q \); they are related to the corresponding mass-eigenstates \( \tilde{q}_a = \{\tilde{q}_1, \tilde{q}_2\} \) by a rotation \( 2 \times 2 \) matrix (we neglect intergenerational mixing):

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

\(^3\)In the MSSM there are several additional 2-body decays [10, 11] of the top quark and also a host of exotic 3-body final states worth studying [12].
From these Lagrangians it is clear that the parameter $\tan \beta$ is called to play a fundamental role in Higgs-top-bottom interactions and, of course, in any SUSY-like version of them. In the MSSM we have a full plethora of additional Higgs-like interactions originally involving the same Yukawa couplings; namely, the interactions of matter fermions and sfermions with higgsinos, the spin 1/2 (super-) companions of higgses. Because of SSB, the higgsinos mix with the gauginos (the fermion partners of the gauge bosons) and in the mass-eigenstate basis form the so-called charginos and neutralinos. In the end one obtains a set of supersymmetric vertices of the type fermion-sfermion-chargino/neutralino involving a fairly complicated admixture of top and bottom Yukawa couplings affecting the loop structure of both the conventional decay $t \rightarrow W^+ b$ and the unconventional mode $t \rightarrow H^+ b$. On the face of it, it is patent that a $\tan \beta$-enhanced bottom-quark Yukawa coupling, i.e. $\tan \beta \sim m_t/m_b$, may have a drastic impact on the MSSM phenomenology of the top-quark.

In spite of the already abundant literature on the various aspects, whether theoretical or experimental, of the charged Higgs decay of the top quark, we believe that nowadays it definitely deserves a renewed interest. There are in part new theoretical reasons, but above all there are intriguing phenomenological motivations grounded on the most recent results from the high precision world of Z-boson physics. Theoretically, the large $\tan \beta$ regime is naturally suggested in top-bottom-tau Yukawa coupling unification models as well as in many string-like unification schemes where, in order to get the radiative electroweak symmetry-breaking pattern, one is forced to depart from the canonical universal boundary conditions on the scalar masses at the GUT scale. Most remarkable, the large $\tan \beta$ regime in conjunction with a moderate value of the CP-odd Higgs mass around $m_{A^0} \approx 50$ GeV has been insistently projected by phenomenological analyses of Z-boson observables within the context of the MSSM, such as in comprehensive global fits of electroweak precision data and in thorough scrutinies of the MSSM parameter space. These studies were aimed at solving, or at least alleviating, the discrepancies (at the $2 - 3 \sigma$ level) between the strict SM theoretical prediction and the corresponding experimental measurements of several Z-boson observables, most conspicuously the $R_b$, $R_c$ branching ratios and the lineshape value of the strong coupling constant at the scale of the $Z$-boson.

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4 Notice that the SUSY interactions specified in the Lagrangian enter the dynamics of $t \rightarrow H^+ b$ through virtual loop corrections (Cf. Fig.1 below).

5 A detailed interaction Lagrangian is given e.g. in eqs.(18)-(19) of Ref. [14].
boson mass, $\alpha_s(M_Z)$; and they have prompted some speculations on new physics \[24\], e.g. claiming the existence of relatively light sparticles \[26, 27\].

In the aforementioned region of parameter space one finds, at the tree-level, the ratio

$$\frac{\Gamma_0(t \to H^+ b)}{\Gamma_0(t \to W^+ b)} = \left(1 - \frac{M_{H^\pm}^2}{m_t^2}\right)^2 \frac{m_t^2 \tan^2 \beta + \cot^2 \beta}{m_t^2 \left(1 - \frac{M_W^2}{m_t^2}\right)^2 \left(1 + 2 \frac{M_W^2}{m_t^2}\right)}.$$  \hspace{1cm} (7)

We see from it that if $M_{H^\pm} \simeq M_W$ (by the way, a situation perfectly compatible with $m_{A^0} \simeq 50$ GeV \[13\]), there are two regimes of $\tan \beta$ where the width of the charged Higgs decay becomes of the same order as (or larger than) the conventional decay width: namely, i) for $\tan \beta \leq 1$, and ii) for $\tan \beta \geq m_t/m_b \sim 36$. The critical status of the decay $t \to H^+ b$ occurs at the intermediate value $\tan \beta = \sqrt{m_t/m_b} \sim 6$, where its partial width has a pronounced dip. Around this value, the charged Higgs mode is overwhelmed by the canonical mode $t \to W^+ b$. Sufficiently away from the dip, however, $t \to H^+ b$ becomes competitive with $t \to W^+ b$. As mentioned above, we do have some theoretical and experimental motivations \[22\],\[27\] to contend both that $M_{H^\pm} = \mathcal{O}(M_W)$ and that at least one of the two $\tan \beta$ regimes i) or ii) applies, most probably the latter.

In view of the potential interest of the decay mode $t \to H^+ b$, one would naturally like to address the computation of the virtual loop corrections to its partial width. Of these, the conventional QCD corrections have already been considered in detail in Ref.\[28\] 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. Similarly, the SM one-loop corrections to $t \to W^+ b$ are known; they are basically dominated by the QCD gluonic contributions ($\simeq -8\%$) \[29\] plus small ($\simeq +1\%$) electroweak corrections \[30\]. Thus we may concentrate on just the MSSM additional loop diagrams. Here we shall focus on the strong and electroweak SUSY corrections

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\[6\] For alternative –topcolour and related– approaches to the $R_b$ anomaly, see e.g. Ref.\[25\].

\[7\] Although the analysis of Ref.\[19\] also emphasizes the role played by light sparticles, it misses –in contrast to the systematic approach of Refs.\[20\]-\[22\],\[27\]– crucial quantum effects, originated in the large $\tan \beta$ region from the neutral components of the MSSM Higgs sector, which should by no means be understated.
to the standard top-quark decay \( t \to W^+ b \) \[14\], \[31\], and on the strong SUSY corrections to the unconventional decay \( t \to H^+ b \) \[32\]. The direct, process-dependent, SUSY Feynman diagrams contributing to these decays are sketched in Fig.1. The analysis of the larger and far more complex body of SUSY-electroweak Feynman diagrams contributing to \( t \to H^+ b \), namely the corrections mediated by squarks, sleptons, chargino-neutralinos and the Higgs bosons themselves, is currently under study and will soon be available \[33\].

The basic free parameters of our analysis, in the electroweak sector, are contained in the stop and sbottom mass matrices:

\[
\mathcal{M}_t^2 = \begin{pmatrix}
M_{tL}^2 + m_t^2 + \cos 2\beta \left( \frac{1}{2} - \frac{2}{3} s_W^2 \right) M_Z^2 & m_t M_{tLR} \\
\frac{1}{2} - \frac{2}{3} s_W^2 \end{pmatrix},
\]

\[
\mathcal{M}_b^2 = \begin{pmatrix}
M_{bL}^2 + m_b^2 + \cos 2\beta \left( -\frac{1}{2} + \frac{1}{3} s_W^2 \right) M_Z^2 & m_b M_{bLR} \\
\frac{1}{2} - \frac{2}{3} s_W^2 \end{pmatrix},
\]

These mass matrices are diagonalized by means of the rotation matrices \[3\]. We have defined

\[
M_{tLR} = A_t - \mu \cot \beta, \quad M_{bLR} = A_b - \mu \tan \beta,
\]

\( \mu \) being the SUSY Higgs mass parameter in the superpotential\[8\]. The \( \tilde{q}_L,R \) are soft SUSY-breaking masses\[7\]; by \( SU(2)_L \)-gauge invariance we must have \( M_{tL} = M_{bL} \), whereas \( M_{tR}, M_{bR} \) are in general independent parameters. In the strong supersymmetric sector, the basic parameter is the gluino mass, \( m_{\tilde{g}} \), and the interaction Lagrangian defining the SUSY-QCD gluino interactions with squarks is the following:

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

where \( \tilde{g}^r (r = 1, 2, ..., 8) \) are the Majorana gluino fields, and \( \left( \lambda_r \right)_{ij} (i,j = 1,2,3) \) are the Gell-Mann matrices.

Next we shortly review\[9\] the actual results of our comparative analysis of the SUSY quantum effects on the partial widths \( \Gamma(t \to W^+ b) \) and \( \Gamma(t \to H^+ b) \) in the on-shell renormalization scheme \[34\]. For convenience we define the relative correction,

\[
\delta = \frac{\Gamma - \Gamma_0}{\Gamma_0},
\]

\[8\]Its sign is relevant in the numerical analysis. We have corrected a known inconsistency in this sign as it appears in Ref.\[13\], and have fixed it as in eq.(3) of Ref.\[14\].

\[9\]For renormalization niceties, detailed one-loop formulae and exhaustive numerical analyses, see Refs.\[14\],\[31\]-\[33\].
Figure 1: Feynman diagrams for the one-loop electroweak and strong SUSY corrections to $t \to W^+ b$ and $t \to H^+ b$. Loop diagrams are summed over all possible values of the mass-eigenstate charginos ($\Psi^{\pm}_i; i = 1, 2$), neutralinos ($\Psi^{0}_\alpha; \alpha = 1, 2, \ldots, 4$), stop and sbottom squarks ($\tilde{b}_a, \tilde{t}_b; a, b = 1, 2$) and gluinos ($\tilde{g}_r; r = 1, 2, \ldots, 8$). Virtual Higgs contributions are not depicted (Cf. Ref.[33]).
with respect to the corresponding tree-level width, $\Gamma_0$. In Fig.2 $\delta^{SUSY}$ and $\delta_\tilde{g}$ refer to the SUSY-electroweak and SUSY-gluino (i.e. SUSY-QCD) corrections, respectively, both given in the $G_F$-scheme [14], viz. parametrized in terms of $G_F$ (Fermi’s constant in $\mu$-decay) by using

$$G_F = \frac{\sqrt{2}}{2M_W s_W^2} (1 + \Delta r^{MSSM}), \tag{13}$$

where $\Delta r^{MSSM}$ is the prediction of the $\mu$-decay parameter $\Delta r$ in the MSSM. Clearly, the term $\Delta r^{MSSM}$ is only relevant for the electroweak corrections (Fig.2a) and is does not affect the strong contributions (Figs.2b-2c) due to the absence of one-loop QCD effects in $\mu$-decay.

A crucial parameter to be tested is $\tan \beta$. In Fig.2a we plot the SUSY-electroweak corrections to the standard decay $\Gamma(t \to W^+ b)$ as a function of $\tan \beta$ for given choices of the other parameters [14]. We see that they can be of order $-10\%$ for very large $\tan \beta$. This is in contrast to the corrections from the two-doublet Higgs sector of the MSSM where in comparable conditions they are one order of magnitude smaller [36], as it is also the case with those from the one-doublet Higgs sector of the SM [30]. Noteworthy is also the fact that the supersymmetric electroweak corrections are of the same (negative) sign and could be of the same order of magnitude as the conventional QCD corrections [29]. On the whole the standard QCD plus SUSY-electroweak corrections to $t \to W^+ b$ could reduce this partial width by about $10-15\%$. Consequently, a measurable reduction beyond $\sim 8\%$ (QCD) could be attributed to a “genuine” SUSY effect. The fact that the chargino-neutralino sector of the MSSM could afford a non-negligible quantum correction to the top quark decay width, in contradistinction to the inappreciable yield from the scalar Higgs sector of the MSSM, can be traced to the highly constrained structure of the Higgs potential as dictated by SUSY [13].

In Fig.2b we study the SUSY-QCD corrections to the standard decay $t \to W^+ b$. Here we have fixed $m_{\tilde{b}} = m_{\tilde{g}} = 120 \, GeV$ for the sbottom and gluino masses and plot contour lines of constant $\delta_\tilde{g}$ in the $(M_{LR}^t, m_{\tilde{t}_1})$-plane [31]. The excluded zone in Fig.2b violates the condition $M_{LR}^2 > 0$ in the stop mass matrix [8]. Notice that there is a threshold (pseudo) singularity (dashed line) associated to the wave-function renormalization of the top-quark field at $m_{\tilde{t}_1} = 54 \, GeV$ for $m_t = 174 \, GeV$, where $m_{\tilde{t}_1}$ is the lightest stop mass. We cannot arbitrarily approach this line from above with-

\[A\] A dedicated study of $\Delta r^{MSSM}$ has been presented in Ref [35].
Figure 2: (a) Electroweak SUSY corrections to $\Gamma(t \to W^+ b)$ as a function of $\tan \beta$ for three values of $(\mu, M)$ and fixed sfermion masses; (b) Isolines in the $(m_{\tilde{t}_1}, M_{t_{LR}})$-plane for the SUSY-QCD corrections $\delta_g$ (in %) to the decay $t \to W^+ b$. (c) SUSY-QCD corrected width of $t \to H^+ b$ versus $\tan \beta$ for $\mu = \pm 100 GeV$. Also indicated are the tree-level widths for the standard and charged Higgs decay modes.
out violating perturbation theory, but we see that even staying prudentially away from it the SUSY-QCD corrections are appreciably high (\(\lesssim 8\%\)) and negative. On the contrary, if we approach the threshold line from below (a non-singular limit), the correction is positive and of order 5%. Nevertheless, it should be clear that the most interesting scenario for our decay corresponds to \(\delta_{\tilde{g}} < 0\), since the alternative two-body supersymmetric decay into stop and gluino, \(t \to \tilde{t}_1, \tilde{g}\), is then phase-space blocked up. This situation is further preferred by the fact that the strong supersymmetric corrections could be reinforced by the additional negative contributions from the electroweak supersymmetric sector of the MSSM (Fig.2a). Last but not least, the case \(\delta_{\tilde{g}} < 0\) is especial in that the SUSY-QCD loops would add up to the conventional QCD corrections (\(\delta_{QCD} \simeq -8\%)\) [24], so that in favourable circumstances the total strong correction could reach \(-(15-18)\%\). Therefore, as the strong SUSY corrections to \(t \to W^+ b\) are insensitive to \(\tan \beta\) [31], we may envision an scenario with large \(\tan \beta\) (\(\gtrsim \frac{m_t}{m_b}\)) in which the electroweak supersymmetric corrections, being also negative, are of the same order of magnitude as the SUSY-QCD contributions studied here; hence the total MSSM pay-off to the top quark width –the Higgs correction being negligible [36]– could result in an spectacular reduction of \(\Gamma(t \to W^+ b)\) by about 25%.

In Fig.2c we turn our attention to the alternative decay \(t \to H^+ b\) and plot the corresponding SUSY-QCD corrected width, \(\Gamma = \Gamma(t \to H^+ b)\), versus \(\tan \beta\) for \(\mu = +100\, GeV\) and \(\mu = -100\, GeV\), and for given values of the other parameters [32]. In particular, the stop and sbottom mixing mass terms are \(M^t_{LR} = -\mu \cot \beta\) and \(M^b_{LR} = 0\), respectively, and we assume that the two diagonal elements in \(M^t_{2}\) (resp. in \(M^b_{2}\)) are equal. We see that \(\Gamma\) rapidly increases with \(\tan \beta\), the preferred range singled out by the \(Z\) lineshape observables [22, 27]. Highly remarkable is also the incidence of the parameter \(\mu\). Indeed, the sign of \(\delta_{\tilde{g}}\) happens to be opposite to the sign of \(\mu\) and the respective corrections for \(\mu\) and for \(-\mu\) take on approximately the same absolute value. The sign dependence of \(\delta_{\tilde{g}}\) suggests that two extreme scenarios could take place with the SUSY-QCD corrections to \(\Gamma(t \to H^+ b)\): namely, they could either significantly enhance the, negative, conventional QCD corrections [28], or on the contrary they could counterbalance them and even result in opposite sign. Worth noticing is also the dependence of these corrections on the gluino mass. In Fig.2c we have fixed \(m_{\tilde{g}} = 200\, GeV\), which is rather heavy. As a matter of fact, we have checked [32] that the decoupling rate of the gluinos is very slow, to the extent that
it fakes for a while, so to speak, a non-decoupling behaviour. This trait is caused by
the presence of a long sustained local maximum (or minimum, depending on the sign
of $\mu$) spreading over a wide range of heavy gluino masses centered at $\sim 300 GeV$ \[32\].
For this reason, heavy gluinos are in the present instance preferred to light gluinos,
contrary to naive expectations\[11\].

Let us mention that we chart significant differences in our analysis as compared to
preliminary calculations in the literature. In Refs.\[38\] a first study of the SUSY-QCD
corrections to $t \rightarrow H^+ b$ was presented, but they neglect a crucial piece of the analysis,
viz. the bottom-quark Yukawa coupling, and as a consequence they are incorrectly
sensitive to the relevant high tan $\beta$ corrections (see Fig.2c). A similar situation occurs
with the incomplete SUSY treatment of $t \rightarrow W^+ b$ in Ref.\[33\]. Furthermore, the
impact from mixing effects and the incidence of the various parameter dependences
were completely missed and only the simplest pattern, characterized by degenerate
masses, was considered\[12\]. Moreover, in the framework of these references, the setting
$m_b = 0$ makes the lowest-order width 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.

In summary, the SUSY contributions to the partial widths of $t \rightarrow W^+ b$ and
$t \rightarrow H^+ b$ could be sizeable, especially in the latter decay where they may comfortably
reach several $\pm 10\%$ even for $\mathcal{O}(100) GeV$ sparticle masses. In the former case an
average $\sim 5\%$ (negative) correction seems more realistic; and although one could also
attain the 10% level (and beyond) one has to wrestle harder with the parameters.
Moreover, it should not be understated the fact that the quantum corrections to
these decays hold in a region of the MSSM parameter space prompted by the high
precision $Z$-boson observables \[22, 27\], and in this region $t \rightarrow H^+ b$ has an appreciable
branching ratio as compared to the standard decay $t \rightarrow W^+ b$. The potential size of
the SUSY effects on $t \rightarrow H^+ b$ stems not only from the strong interaction character
of the SUSY-QCD corrections, but also from the high sensitivity of this decay mode
to the (weak-interaction) SSB parameter tan $\beta$. Barring the tan $\beta << 1$ regime –
considered as very unlikely from the point of view of model building–, we see that
the relevance of the charged Higgs decay mode of the top quark is ultimately linked

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\[11\] The existence of light gluinos of $\mathcal{O}(1) GeV$ is not yet completely excluded. They could decisively
influence the MSSM phenomenology in other instances, as shown in Ref.\[37\].

\[12\] Notice that the assumption of stop masses equal to sbottom masses is incompatible with the
structure of the mass matrices \[8\]-\[10\].
to the dynamics of the bottom-quark Yukawa coupling.

All in all, we believe that the Higgs mode reveals itself as an ideal environment where to study the nature of the SSB mechanism. We might even venture into saying that $t \rightarrow H^+ b$ stands among the best candidate processes where to target our long and (yet) unsuccessful search for “virtual Supersymmetry”; that is, it could be an optimal place where to enquire for huge, and slowly decoupling, quantum supersymmetric effects. In this respect it should be stressed that the typical size of our corrections is maintained even for sparticle masses well above the LEP 200 discovery range. Theses 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 $\tan \beta$ and the corresponding electroweak corrections are only moderately sensitive to this parameter. 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. As a very promising example, we remark the future measurement of the cross-section for single top-quark production, which is directly sensitive to the top-quark width. Thus, in favorable circumstances, we should be able to unravel the existence of new physics out of a precise measurement of the top-quark width, or related observables, at a modest –and attainable– precision of $\sim 5 - 10\%$.

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$^{13}$In the event that the charged Higgs boson is heavier than the top quark, the alternative mode $H^+ \rightarrow t b$ is expected to exhibit similar properties.
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