$t \to W^+ b$ AND $t \to H^+ b$ AT THE QUANTUM LEVEL IN THE MSSM

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ABSTRACT

We compare the standard top quark decay and the charged Higgs decay of the top quark, $t \to W^+ b$ and $t \to H^+ b$, at the quantum level in the MSSM. While the SUSY loop corrections to the standard top quark decay are only of a few percent, it turns out that $t \to H^+ b$ is a most promising candidate for carrying large quantum SUSY signatures. As a result, the $(\tan \beta, M_{H^\pm})$ exclusion plots presented by the CDF Collaboration should be thoroughly revised in the light of the MSSM.
Theoretically, Supersymmetry (SUSY) is perhaps the only known framework beyond the Standard Model (SM) which is capable of extending non-trivially the quantum field theoretical structure of the conventional strong and electroweak interactions while keeping all the necessary ingredients insuring internal consistency, such as gauge invariance and renormalizability. In particular, the Minimal Supersymmetric Standard Model (MSSM) \[1, 2\] has been able to accommodate all known high precision measurements to a similar degree of significance as the Standard Model \[2\]. Remarkably, SUSY has been able to survive over the years and it has become a “fact” of live for many physicists. Most likely this situation will remain invariable at least until the Tevatron II and LHC eras have explored in full the experimental feasibility of the MSSM.

In this talk I propose to dwell on the supersymmetric phenomenology of top quark decays with an eye on the Tevatron and LHC phenomenological capabilities. Among the relevant MSSM top quark decays potentially carrying a direct or indirect SUSY signature, the following two-body modes stand out:

\begin{align*}
&i) \quad t \to W^+ b \\
&ii) \quad t \to H^+ b \\
&iii) \quad t \to \tilde{t}_a \chi_0^a, \\
&iv) \quad t \to \tilde{b}_a \chi_i^+, \\
v) \quad t \to \tilde{t}_a \tilde{g}.
\end{align*}

Therein, \(\tilde{t}_a, \tilde{b}_a, \chi_i^+, \chi_0^a, \tilde{g}_r\) \((a, i = 1, 2; \alpha = 1, 2, ..., 4; r = 1, 2, ..., 8)\) denote stop, sbottom, chargino, neutralino and gluino “sparticles”, respectively. (Also quite a few three-body decays are possible and have been studied \[3\].) Of course, decay i) is the SM top quark decay, and decay ii) into a charged Higgs need not to be a SUSY decay. However, by studying the possible MSSM quantum effects on these decay modes one may hope to unveil indirect traces of the underlying SUSY dynamics. On the other hand the last three decays in (1) do carry an explicit SUSY signature. In general also these decays may require a higher order treatment, the reason being that some of the final state signatures, after the sparticles have decayed into conventional particles and the LSP (typically the lightest neutralino \(\chi_0^1\)), they may well mimic the standard top quark decay. For example, decay iii) may lead to a signature similar to the standard top quark decay into the final states \(b l^+ \nu\) or \(b + 2\) jets; for the stop could decay into \(\chi_i^+ b\), and subsequently yield the chain \(\chi_i^+ \to \chi_0^1 W^* \to \chi_0^1 l^+ \nu\) or \(\chi_0^1 + 2\) jets. Therefore, a detailed treatment of these direct SUSY modes is in principle desirable to help disentangling the nature of the complicated final configurations and to enable a reliable determination of the top quark cross-section within the MSSM. Barring a light gluino window, which is nowadays harder and harder to maintain, current limits on squark and gluino masses already rule out decay v) and
most likely also decay iv). However, even keeping alive this last decay, unfortunately the
typical size of the corrections to the two processes iii) and iv) is not too significant (at the
ten per cent level at most [4]). While this would amply suffice in a high precision machine
such as LEP, nevertheless for measurements to be performed in a hadron environment it
is probably not enough to be detected.

The fact that sparticles seem to be rather heavy makes direct SUSY searches more
and more difficult. For this reason it may be advisable to hunt for "quantum signatures"
by means of the indirect method of high precision measurements of less exotic, and so
more manageable, processes. Thus we shall report here on the behaviour of the more
conventional decays i) and ii) at the quantum level. To start with, we recall that the
supersymmetric strong (SUSY-QCD) and the supersymmetric electroweak (SUSY-EW)
corrections to the standard top quark decay $t \to W^+ b$ are well understood [5]. The
leading one-loop vertex functions are shown in Fig. 1 (taking the external dashed line as
the $W^+$). The possibility of large non-standard quantum effects lies to a great extent on
the influence of the parameter $\tan \beta$ [6]. In supersymmetric theories like the MSSM, $\tan \beta$
enters the top and bottom quark Yukawa couplings of the superpotential through $1/\sin \beta$
and $1/\cos \beta$, respectively:

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

(2)

and therefore one may expect an enhancement of the Yukawa couplings as compared to
the gauge couplings both at low and high $\tan \beta$. Notice that the bottom-quark Yukawa
coupling may counterbalance the smallness of the bottom mass at the expense of a large
value of $\tan \beta$.

Apart from $\tan \beta$, the basic free parameters of our analysis concerning the electroweak
sector are contained in the stop and sbottom mass matrices ($\tilde{q} = \tilde{t}, \tilde{b}$):

$$\mathcal{M}_{\tilde{q}}^2 = \begin{pmatrix} \mathcal{M}_{11}^2 & \mathcal{M}_{12}^2 \\ \mathcal{M}_{12}^2 & \mathcal{M}_{22}^2 \end{pmatrix},$$

(3)

with

$$\mathcal{M}_{11}^2 = M_{\tilde{q}_L}^2 + m_q^2 + \cos 2\beta (T_3 - Q_q \sin^2 \theta_W) M_Z^2,$$

$$\mathcal{M}_{22}^2 = M_{\tilde{q}_R}^2 + m_q^2 + Q_q \cos 2\beta \sin^2 \theta_W M_Z^2,$$

$$\mathcal{M}_{12}^2 = m_q M_q^{LR},$$

$$M_{LR}^{(t, b)} = A_{(t, b)} - \mu \{ \cot \beta, \tan \beta \}.$$  

(4)

We denote by $m_{\tilde{t}_1}$ and $m_{\tilde{b}_1}$ the lightest stop and sbottom masses.
Figure 1: SUSY-QCD and SUSY-EW one-loop vertices for $t \rightarrow W^+ b$ and $t \rightarrow H^+ b$. The EW “inos” $\Psi_{i,\alpha}$ are unphysical mass-eigenstates related to the physical states $\chi_{i,\alpha}$ (Cf. Ref.[3]).

The numerical results are conveniently cast in terms of the relative correction with respect to the corresponding tree-level width $\Gamma_0$:

$$\delta = \frac{\Gamma(t \rightarrow (W^+, H^+) b) - \Gamma_0(t \rightarrow (W^+, H^+) b)}{\Gamma_0(t \rightarrow (W^+, H^+) b)}.$$ 

Let us project out the total SUSY correction (5) for decay i), i.e. the total MSSM correction after subtracting the SM part. The latter is defined (in the MSSM context) as the one obtained by decoupling the sparticle effects and leaving only the lightest CP-even Higgs contribution ($h^0$) in the limit of infinite CP-odd Higgs mass ($M_A \rightarrow \infty$) [6]. In the on-shell $G_F$-scheme, which is characterized by the set of inputs ($G_F, M_W, M_Z, m_f, M_{SUSY}, ...$), the total SUSY correction $\delta_{SUSY}$ is negative and of the order of a few per cent (except in some unlikely cases [5]). In Fig. 2a it is shown the dependence of the corrections on the crucial parameter $\tan \beta$, for a typical set of parameters. Here $\mu, M$ are the higgsino and $SU(2)$-gaugino mass parameters, respectively, and $A$ is the value of a universal trilinear coupling. In spite of the enhancement at high $\tan \beta$ the negative shift of the decay amplitude is below 4% even for $\tan \beta = 50$. Therefore, for $1 \lesssim \tan \beta \lesssim 30$ the negative SUSY effects approximately cancel out against the (positive) electroweak SM contributions, which are of the same order of magnitude ($\lesssim +2\%$), leaving the ordinary QCD effects [7] ($\simeq -10\%$) as the net MSSM corrections. As a result no significant imprint of the underlying SUSY dynamics is left behind the standard top quark decay $t \rightarrow W^+ b$ and we are thus led to examine other top decays beyond the SM.

In contrast to decay i), decay ii) may receive spectacularly large SUSY quantum corrections, namely of the order of 50%, which certainly could not be missed – if SUSY is there at all. For this reason, we concentrate on that decay. To be sure, $t \rightarrow H^+ b$ has been object of many studies in the past, mainly within the context of general two-Higgs-
Figure 2: (a) The total SUSY correction to $t \to W^+ b$ for given set of parameters, as a function of $\tan \beta$; (b) The SUSY-EW, standard EW, SUSY-QCD, standard QCD and full MSSM corrections as a function of $\tan \beta$. Remaining inputs as in (a); (c) As in (b), but as a function of $m_{b_1}$; (d) As a function of $m_{\tilde{g}}$. 
doublet models (2HDM). From the experimental point of view that decay has been thoroughly scrutinized at the Tevatron [8]. Recently a systematic study has been made on multilepton and multijet signatures for the charged Higgs decay of the top quark at the Tevatron that could be useful to constraint the 2HDM parameter space [4]. However, it is shown that the current CDF data [10] on 2 b-jets and 1 lepton channel do not pose any real restriction on the charged Higgs decay of the top quark. On the other hand no systematic treatment of the MSSM quantum effects of the decay \( t \to H^+ b \) existed in the literature until the works of Refs. [11] and [12]. Moreover, remember that in the MSSM, in contrast to the general 2HDM, the charged Higgs can elude the stringent lower mass bounds following from radiative \( B \)-decays (\( b \to s \gamma \)) [13] that would preclude the decay under consideration. Admittedly, the situation with radiative \( B \)-decays is not completely clear since there are many sources of error that deserve further experimental consideration. Still this information can be used to single out the SUSY nature of the Higgs sector. Thus in the MSSM the existence of decay ii) is more tenable than in the framework of an unconstrained 2HDM. Next we briefly review the results concerning the important MSSM quantum corrections potentially affecting its decay width.

We present our results for the decay ii) also in the on-shell scheme. In considering the various parameter dependences, again a fundamental parameter to be tested is \( \tan \beta \). This parameter is involved explicitly in the relevant interaction Lagrangian for the decay ii), namely

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

where \( P_{L,R} = 1/2(1 \mp \gamma_5) \) are the chiral projector operators. Therefore, crucial for the treatment of the SUSY-EW effects on the decay ii) is the definition of \( \tan \beta \) beyond the tree-level. Following Ref.[11] we define it by means of the \( \tau \)-lepton decay of \( H^\pm \):

\[
\Gamma(H^+ \to \tau^+ \nu_\tau) = \frac{\alpha m_\tau^2 M_H}{8 M_W^2 s_W^2} \tan^2 \beta .
\]

This definition generates a counterterm

\[
\frac{\delta \tan \beta}{\tan \beta} = \frac{1}{2} \left( \frac{\delta M_W^2}{M_W^2} - \frac{\delta g^2}{g^2} \right) - \frac{1}{2} \delta Z_H \cot \beta \delta Z_{HW} + \Delta_\tau .
\]

Here \( \Delta_\tau \) comprises the complete set of MSSM one-loop effects on the \( \tau \)-lepton decay of \( H^\pm \); \( \delta Z_H \) and \( \delta Z_{HW} \) stand respectively for the charged Higgs and mixed \( H - W \) wavefunction renormalization factors; and the remaining counterterms \( \delta g^2 \) and \( \delta M_W \) are the standard ones in the on-shell scheme [2].

For a typical choice of parameters, in Fig.2b we plot the various contributions to (5) from SUSY-QCD, SUSY-EW and the MSSM Higgs sector, as a function of \( \tan \beta \).
The standard QCD correction is also shown [14]. The full MSSM correction $\delta_{\text{MSSM}}$ is defined to be the sum of all these individual contributions. As it turns out that for this decay $\delta_{\text{MSSM}} >> \langle \Delta r \rangle_{\text{MSSM}}$ it follows that the difference between the results in the $G_F$-scheme and the $\alpha$-scheme [11] is not material in this case; hence the bulk of the effect is already contained in the $\alpha$-parametrization. This was certainly not the case with decay i). In Fig.2c we display the evolution of the different corrections with $m_{\tilde{b}_1}$; this is a critical parameter governing the size of the leading (SUSY-QCD) corrections. Although $\delta_{\text{SUSY-QCD}}$ dies away relatively fast with increasing $m_{\tilde{b}_1}$, for large sbottom masses there remains an undampened SUSY-EW component (essentially controled by $m_{\tilde{g}}$) which can be sizeable enough for stop masses in the few hundred GeV. The decoupling with the gluino mass is much slower, and with the remarkable property that before entering the decoupling regime it has a long sustained local maximum around $m_{\tilde{g}} = 500$ GeV (Cf. Fig. 2d). Finally we mention that the corrections also increase with $A_t$ and $|\mu|$, and change sign with $\mu$. Of course, $\delta_{\text{MSSM}} \to 0$ when all sparticle masses increase simultaneously.

The definition (7) of $\tan \beta$ allows to renormalize the $H^\pm t \bar{b}$-vertex in perhaps the most convenient way to deal with our main decay ii). Indeed, from the practical point of view, we should recall the excellent methods for $\tau$-identification developed by the Tevatron collaborations and recently used by CDF to study the existence region of the decay iv) in the $(\tan \beta, M_H)$-plane [8]. However, we wish to show that this analysis may undergo dramatic changes when we incorporate the MSSM quantum effects [12]. Although CDF utilizes inclusive $\tau$-lepton tagging, for our purposes it will suffice to focus on the exclusive final state $(l, \tau)$, with $l$ a light lepton, as a means for detecting an excess of $\tau$-events [15]. To be precise, we are interested in the $t \bar{t}$ cross-section leading to the decay sequences $t \bar{t} \to H^+ b, W^- \bar{b}$ and $H^+ \to \tau^+ \nu_\tau, W^- \to l \bar{\nu}_l$, and vice versa. The relevant quantity
can be easily derived from the measured value of the canonical cross-section \( \sigma_{t\bar{t}} \) for the standard channel \( t \to b l \nu_l, \bar{t} \to b q q' \), after inserting appropriate branching fractions, namely [12]

\[
\sigma_{\tau\tau} = \left[ \frac{4}{81} \epsilon_1 + \frac{4}{9} \frac{\Gamma(t \to H b)}{\Gamma(t \to W b)} \epsilon_2 \right] \sigma_{t\bar{t}}. \tag{9}
\]

The first term in the bracket comes from decay i), and for the second term we assume (at high \( \tan \beta \)) 100\% branching fraction of \( H^+ \) into \( \tau \)-lepton, as explained before. Finally, \( \epsilon_i \) are detector efficiency factors. Thus, in most of the phase space available for top decay the bulk of the cross-section (9) is provided by the contribution of decay ii). Consequently, the observable (9) should be highly sensitive to MSSM quantum effects. In fact, from the non-observation of any excess of \( \tau \)-events at the Tevatron, in Figs. 3a and 3b we derive the (95\% C.L.) excluded regions for \( \mu < 0 \) and \( \mu > 0 \), respectively. We point out that the region of MSSM parameter space considered in this analysis can be shown to be compatible with the \( b \to s \gamma \) constraints mentioned above [12]. From inspection of these figures it can hardly be overemphasized that the MSSM quantum effects can be dramatic. In particular, while for \( \mu < 0 \) the MSSM-corrected curve is significantly more restrictive than the QCD-corrected one, for \( \mu > 0 \) the bound essentially disappears from the perturbative region (\( \tan \beta \lesssim 60 \)). Notice that in the latter case the SUSY correction is negative and so it adds up to the ordinary QCD effects. That is why in this case we have used a heavier SUSY spectrum than in Fig. 3a in order that the results remain perturbative.

We conclude by pointing out the recent work of Ref.[16]. Using Tevatron data in the \( b\bar{b}\tau^+\tau^- \) channel these authors improve the bound in the \( (\tan \beta, M_H) \)-space. Nonetheless this analysis was performed only at the tree-level and hence it could undergo significant MSSM radiative corrections. The potentially large effects not included in that paper stem from the production mechanism of the CP-odd Higgs boson \( A^0 \) (through \( b\bar{b} \)-fusion) before it decays into \( \tau^+\tau^- \) pairs. Indeed, the \( b\bar{b}A^0 \) vertex is known [17] to develop important MSSM corrections in the relevant regions of the \( (\tan \beta, M_H) \)-plane purportedly “excluded” by the tree-level analysis of Ref.[16]. Clearly, a detailed re-examination of the excluded region at the quantum level is in order within the context of the MSSM [18]. The lesson to be learnt should be highly instructive: namely, in contrast to the tiny corrections to gauge boson observables, the MSSM quantum effects on top-Higgs boson physics can be rather large and should not be neglected in future searches at the Tevatron and at the LHC.

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