LOOKING FOR QUANTUM SUSY SIGNATURES IN TOP QUARK DECAYS AT HADRON COLLIDERS

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ABSTRACT

We discuss the supersymmetric quantum effects on top quark decays within the MSSM. It turns out that $t \to H^+ b$ is the most promising candidate for carrying large quantum SUSY signatures. As a result, the recent $(\tan \beta, M_{H^\pm})$ exclusion plots presented by the CDF Collaboration should be thoroughly revised in the light of the MSSM.

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In this talk I propose to dwell on the supersymmetric phenomenology of top quark decays with an eye on future machine developments such as the upgrade of the Tevatron and the advent of the LHC. In the absence of direct sparticle production, one naturally looks for “quantum signatures” of the new physics by means of the indirect method of high precision measurements. 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 \cite{1}. This fact alone, if we bare in mind the vast amount of high precision data available both from low-energy and high-energy physics, should justify all efforts to search for SUSY in present day particle accelerators. In this respect we wish to stress here the possibility of seeing large virtual effects of SUSY through the interplay between top quark and Higgs boson dynamics at hadron colliders. The typical size of the effects that we are referring to is in general much larger than the tiny few per cent level corrections predicted in all canonical gauge boson observables at LEP.

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 \cite{2}, but unfortunately they are not too large – as typically expected of gauge boson interactions. In the on-shell \( G_F \)-scheme, which is characterized by the set of inputs \((G_F, M_W, M_Z, m_f, M_{SUSY}, \ldots)\), they are negative and of the order of a few per cent (except in some unlikely cases \cite{3}). Therefore, they approximately cancel out against the positive SM contributions of the same order of magnitude and leaving the ordinary QCD effects \( \simeq -10\% \) as the net MSSM corrections. Hence no significant imprint of underlying SUSY dynamics is left on \( \Gamma(t \to W^+ b) \), and we are led to examine other top quark decays.

Among the relevant MSSM top quark decays carrying an interesting SUSY signature, the following two-body modes stand out:

\begin{enumerate}
  \item \( t \to \tilde{t}_i \chi_0^\alpha \),
  \item \( t \to \tilde{b}_i \chi^+_\alpha \),
  \item \( t \to \tilde{t}_i \tilde{g} \),
  \item \( t \to H^+ b \).
\end{enumerate}

\( (1) \)

Therein, \( \tilde{t}_i, \tilde{b}_i, \chi^+_i, \chi_0^\alpha, \tilde{g} (i = 1, 2; \alpha = 1, 2, \ldots, 4) \) denote stop, sbottom, chargino, neutralino and gluino sparticles, respectively. (Also quite a few three-body decays are possible and have been studied \cite{3}.)

While the first three decays in (1) already carry a direct SUSY signature, the third one is meant to involve the charged Higgs boson of the MSSM and it could bring along a significant quantum SUSY signature. In general the direct SUSY decays i)-iii) may also 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 $\chi_0^0$), they may well mimic the standard top quark decay. For example, decay i) 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 $\chi_1^0 b$, and subsequently yield the chain $\chi_1^+ \rightarrow \chi_0^0 W^* \rightarrow \chi_1^0 l^+ \nu$ or $\chi_1^0 + 2$ 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 iii) and most likely also decay ii). Moreover, the typical size of the corrections to the process i) 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, however for measurements to be performed in a hadron environment it is probably not enough to be detected. In contrast, decay iv) 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 are going to focus on that decay. To be sure, $t \rightarrow H^+ b$ has been object of many studies in the past $[2]$, mainly within the context of general two-Higgs-doublet models ($2HDM$), and it is being thoroughly scrutinized in recent analyses at the Tevatron $[3]$. Notwithstanding, no systematic treatment of the MSSM quantum effects existed in the literature until very recently $[5, 7]$.

The basic free parameters of our analysis concerning the electroweak sector are contained in the stop and sbottom mass matrices ($q = t, b$):

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

with

$$\mathcal{M}_{11}^2 = M_{q_L}^2 + m_q^2 + \cos 2\beta(T_q^3 - Q_q \sin^2 \theta_W) M_Z^2,$$

$$\mathcal{M}_{22}^2 = M_{q_R}^2 + m_q^2 + Q_q \cos 2\theta \sin^2 \theta_W M_Z^2,$$

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

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

We denote by $m_{\tilde{t}_i}$ and $m_{\tilde{b}_i}$ the lightest stop and sbottom masses.

Crucial in the treatment of the electroweak SUSY effects is the definition of $\tan \beta$ beyond the tree-level. Following Ref.$[5]$ we define it by means of the $\tau$-lepton decay of

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$2$See $[3]$ and references therein.
$H^\pm$: 
\[ \Gamma(H^+ \rightarrow \tau^+ \nu_\tau) = \frac{\alpha m_t^2 M_H}{8 M_W^2 s_W^2} \tan^2 \beta. \quad (4) \]

This definition generates a counterterm 
\[ \frac{\delta \tan \beta}{\tan \beta} = \frac{1}{2} \left( \frac{\delta M_W^2}{M_W^2} - \frac{\delta g^2}{g^2} \right) - \frac{1}{2} \delta Z_H \]
\[ + \cot \beta \delta Z_{HW} + \Delta. \quad (5) \]

Notice that $\Delta_r$ above stands for 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 $H - W$ wave-function renormalization factors; and the remaining counterterms $\delta g^2$ and $\delta M_W$ are the standard ones in the on-shell scheme \[8\].

The results are conveniently cast in terms of the relative correction with respect to the corresponding tree-level width, $\Gamma_0$:
\[ \delta_{\text{MSSM}} = \frac{\Gamma_{\text{MSSM}}(t \rightarrow H^+ b) - \Gamma_0(t \rightarrow H^+ b)}{\Gamma_0(t \rightarrow H^+ b)}. \quad (6) \]

We will present the numerical results for this quantity in the on-shell $\alpha$-scheme:

\[ (\alpha, M_W, M_Z, m_f, M_{\text{SUSY}}, \ldots). \quad (7) \]

The corresponding results in the $G_F$-scheme are just $\delta_{\text{MSSM}} - (\Delta r)_{\text{MSSM}}$ \[9\]. As it turns out that $\delta_{\text{MSSM}} \gg (\Delta r)_{\text{MSSM}}$ \[9\], the difference between the two schemes is not material in this case, i.e. the bulk of the effect is already contained in the $\alpha$-parametrization.

A fundamental parameter to be numerically tested is $\tan \beta$. It is involved in the basic interaction Lagrangian for our decay:
\[ L_{Htb} = \frac{g}{\sqrt{2} M_W} H^+ \bar{t} [m_t \cot \beta P_L + m_b \tan \beta P_R] b + \text{h.c.}, \quad (8) \]

where $P_{L,R} = 1/2(1 \mp \gamma_5)$ are the chiral projector operators. Furthermore, $\tan \beta$ in supersymmetric theories, like the MSSM, enters the top and bottom quark Yukawa couplings of the superpotential through $1/ \sin \beta$ and $1/ \cos \beta$, respectively:
\[ h_t = \frac{g m_t}{\sqrt{2} M_W \sin \beta}, \quad h_b = \frac{g m_b}{\sqrt{2} M_W \cos \beta}. \quad (9) \]

Notice that the bottom-quark Yukawa coupling may counterbalance the smallness of the bottom mass at the expense of a large value of $\tan \beta$. For a typical choice of parameters, in Fig.1a we plot the various contributions to (\[\]) from SUSY-QCD, SUSY-EW and the MSSM Higgs sector. We also show the standard QCD correction. The full MSSM correction is defined to be the sum of all these individual contributions. In Fig.1b we
display the evolution of the different corrections with $m_{b_1}$; this is a critical parameter governing the size of the leading (SUSY-QCD) corrections. Indeed, the decoupling with the gluino mass is much slower\cite{5}. Still, even when $m_{b_1}$ is very large, there remains an undamped SUSY-EW component (essentially controlled by $m_{t_1}$) which can be sizeable enough for stop masses in the few hundred GeV. The corrections also increase with $A_t$ and $|\mu|$, and change sign with $\mu$. Of course, $\delta_{MSSM} \to 0$ when all sparticle masses increase simultaneously, for the MSSM naturally decouples in the limit $M_{SUSY} \to \infty$.

The definition (4) of $\tan \beta$ allows to renormalize the $H^\pm t b$-vertex in perhaps the most convenient way to deal with our main decay iv). 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\cite{4}. However, we wish to show that this analysis may undergo dramatic changes when we incorporate the MSSM quantum effects\cite{7}. 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\cite{10}.

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. From the non-observation of these events, in Figs.1c and 1d we derive the (95\% C.L.) excluded regions for $\mu < 0$ and $\mu > 0$, respectively. (In the latter case we choose a heavier SUSY spectrum in order that (6) remains perturbative.) Shown are the tree-level, standard QCD-corrected and fully MSSM-corrected results. 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 ($\tan \beta \lesssim 60$).

The lesson to be learnt should be highly instructive: 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 ignored in future searches at the Tevatron and at the LHC.

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**Figure Captions**

- **Fig.1 (a)** The various individual and total MSSM correction eq. (3) as a function of \( \tan \beta \) and given values of the other parameters: \( M_{H^\pm} = 120 \text{ GeV} \), \( \mu = -150 \text{ GeV} \), \( m_{\tilde{g}} = 300 \text{ GeV} \), \( m_{\tilde{b}_1} = 150 \text{ GeV} \), \( m_{\tilde{t}_1} = 100 \text{ GeV} \), \( A_t = A_b = 300 \text{ GeV} \); (b) As in (a), but as a function of \( m_{\tilde{b}_1} \) and two fixed \( \tan \beta \) values; (c) The 95% C.L. exclusion plot in the \((\tan \beta, M_{H^\pm})\)-plane for \( \mu = -90 \text{ GeV} \) and remaining parameters similar to (a). Shown are the tree-level (dashed), QCD-corrected (dotted) and fully MSSM-corrected (continuous) contour lines. The excluded region in each case is the one lying below the curve; (d) As in (c), but for a \( \mu > 0 \) scenario characterized by a heavier SUSY spectrum.
Figure 1