Hidden top quark with charged-Higgs-boson decay

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(Received 17 August 1989)

In top-quark decays, there may be competition between decays mediated by real or virtual W bosons and decays into a charged Higgs boson. We evaluate the branching fractions to delineate the regions where standard top signals would be suppressed for two-Higgs-doublet models. We also show that both the standard semileptonic-decay signals $t \to bW^+ \to bl\nu$ and unconventional signatures such as $t \to bH^+ \to b\tau\nu$ may be simultaneously suppressed with suitable parameter choices for any top-quark mass; in such circumstances the top quark would be very difficult to detect at hadron colliders.

Extensive experimental searches for the top quark at $p\bar{p}$ colliders[1] have so far not detected its standard-model semileptonic decay signal, indicating that its mass is more than 70 GeV if it has conventional W-mediated decays. However, it may still happen that the t quark is light but its standard semileptonic decay modes are suppressed by competition with additional nonstandard decays, which can occur in two-Higgs-doublet models[2] through the decay[3,4] $t \to bH^+$, where H^+ is a charged Higgs boson. The H^+ would decay dominantly via $H^+ \to \tau^+\nu_\tau$ and $H^+ \to c\bar{s}$ modes, with branching fractions determined by the parameter $\tan\beta = v_2/v_1$, where v_2 and v_1 are the vacuum expectation values of the doublets giving masses to charge $\frac{2}{3}$ and charge $-\frac{1}{3}$ quarks, respectively. For $\tan\beta \gtrsim 1$ the $H^+ \to \tau^+\nu_\tau$ branching fraction is greater than 30% and $t \to bH^+$ decays have new signatures[4] such as missing p_T (denoted $\not{p}_T$) and isolated τ jets that can be identified at $p\bar{p}$ colliders. The competition between $t \to bW^*$ (where W^* denotes real or virtual W) and $t \to bH^+$ depends both on $\tan\beta$ and the masses $m_t, m_{H^\pm}$. It is an important quantitative issue to determine whether the top quark can remain hidden at hadron colliders, with respect to $t \to bW^* \to be\nu$ or $t \to bH^+ \to b\tau\nu$ or both. In the present work we evaluate these branching fractions for a wide range of $\tan\beta$ and $m_t, m_{H^\pm}$ masses.

In the class of two-Higgs-doublet models that we discuss here, flavor-changing neutral currents are avoided[5] by one doublet ϕ_1 giving mass to charge $Q=-\frac{1}{3}$ quarks and the other doublet ϕ_2 giving mass to $Q=\frac{2}{3}$ quarks via vacuum expectation values (VEV's) v_1 and v_2, respectively. The Yukawa couplings of the physical charged Higgs fields $H^\pm = (v_1\phi_2^\pm - v_2\phi_1^\pm)/(v_1^2+v_2^2)^{1/2}$ are given by

$$\mathcal{L} = g(\sqrt{2}M_W)^{-1}H^+(\tan\beta\, \bar{U}_L V M_D D_R + \cot\beta\, \bar{U}_R M_U V D_L + \tan\beta\, \bar{N}_L M_L L_R) + \text{H.c.} , \quad (1)$$

where $U=(u,c,t)$, $D=(d,s,b)$, $L=(e,\mu,\tau)$, $N=(\nu_e,\nu_\mu,\nu_\tau)$, V is the Kobayashi-Maskawa (KM) matrix, M_D and M_U are the $Q=-\frac{1}{3}$ and $Q=\frac{2}{3}$ diagonal quark mass matrices, M_L is the charged-lepton mass matrix, $\tan\beta = v_2/v_1$, and g is the SU(2) electroweak coupling. The mass $m_{H^\pm}$ and the ratio v_2/v_1 are *a priori* unknown in a general two-doublet model. In minimal supersymmetry (SUSY) models,[6] $m_{H^\pm}$ is constrained to be greater than M_W but in E_6 superstring models the bound[7] is $m_{H^\pm} > 53$ GeV. The dominant partial decay widths of $H^\pm$ are

$$\Gamma(H^+ \to \tau^+\nu) = \frac{g^2 m_{H^\pm}}{32\pi M_W^2} m_\tau^2 \tan^2\beta , \quad (2)$$

$$\Gamma(H^+ \to c\bar{s}) = \frac{3g^2 m_{H^\pm}}{64\pi M_W^2}(|A_{cs}|^2 + |B_{cs}|^2) , \quad (3)$$

where

$$A_{cs} = (m_c\cot\beta + m_s\tan\beta)V_{cs} ,$$
$$B_{cs} = (m_c\cot\beta - m_s\tan\beta)V_{cs} , \quad (4)$$

neglecting corrections of order $m_c^2/m_H^2, m_\tau^2/m_H^2, m_s^2/m_H^2$.

The corresponding expressions for $H^+ \to \mu^+\nu_\mu, u\bar{s}, u\bar{b}, c\bar{d}, c\bar{b}$ partial widths are directly analogous to those for $H^+ \to \tau^+\nu, c\bar{s}$. For calculations we take masses (in GeV) $m_u = m_d = m_e = 0$, $m_s = 0.2$, $m_c = 1.5$, $m_b = 5.2$, $M_W = 80.6$, $|V_{cs}| = 0.975$, $|V_{us}| = |V_{cd}| = 0.22$, $|V_{cb}| = 0.043$, and $|V_{ub}| = 0.011$. The branching fractions have a strong β dependence:

$\tan\beta=\frac{1}{2}$	$\tan\beta=1$	$\tan\beta=2$
$B(\tau\nu)=3\%$	31%	82%
$B(c\bar{s})=92\%$	64%	13%
$B(c\bar{d})=5\%$	3.3%	0.5%
$B(c\bar{b})=0.3\%$	1.5%	3.7%

In the approximation of $m_c \simeq m_\tau$, $m_s \simeq 0$, and a diagonal KM matrix, the H^+ branching fractions become

$$B(H^+ \to \tau^+ \nu) \simeq 1/(1+3\cot^4\beta) ,$$
$$B(H^+ \to c\bar{s}) \simeq 1/(1+\tfrac{1}{3}\tan^4\beta) . \quad (5)$$

The partial width for the top-quark decay to a real Higgs boson is

$$\Gamma(t \to bH^+) = \frac{g^2}{128\pi M_W^2 m_t} \{ |A_{tb}|^2[(m_t+m_b)^2 - m_H^2] + |B_{tb}|^2[(m_t-m_b)^2-m_H^2] \} \times \lambda^{1/2}\left[1, \frac{m_b^2}{m_t^2}, \frac{m_H^2}{m_t^2}\right] , \quad (6)$$

where $\lambda(a,b,c)=a^2+b^2+c^2-2ab-2bc-2ca$ and A_{tb}, B_{tb} are defined analogously to Eq. (4). The partial width for the top-quark decay via virtual H^+ to the $\tau^+\nu$ channel is

$$\Gamma(t \to b\tau\nu) = N \int dm_X^2 \frac{\Gamma(t \to bX^+)(m_X^2 - m_\tau^2 - m_{\nu_\tau}^2)}{(m_X^2 - m_H^2)^2 + (\Gamma_H m_H)^2} \times \lambda^{1/2}\left[1, \frac{m_\tau^2}{m_X^2}, \frac{m_{\nu_\tau}^2}{m_X^2}\right] , \quad (7)$$

where X^+ denotes an off-shell Higgs boson with invariant mass squared m_X^2 and $\Gamma(t \to bX^+)$ is defined by Eq. (6) with m_H^2 replaced by m_X^2. The limits of the m_X^2 integration are $(m_\tau + m_{\nu_\tau})^2 \le m_X^2 \le (m_t - m_b)^2$. The total H^+ width Γ_H is found by summing the partial widths Eqs. (2) and (3) and their analogs. The normalization factor N is

$$N = g^2 m_\tau^2 \tan^2\beta / (32\pi^2 M_W^2) . \quad (8)$$

The corresponding partial widths for top-quark decay via real and virtual W bosons are

$$\Gamma(t \to bW^+) = \frac{g^2 m_t^3}{64\pi M_W^2} |V_{tb}|^2 \lambda^{1/2}\left[1, \frac{M_W^2}{m_t^2}, \frac{m_b^2}{m_t^2}\right] \times \left[\left[1+\frac{m_b^2}{m_t^2}\right]^2 + \frac{M_W^2}{m_t^2}\left[1+\frac{m_b^2}{m_t^2}\right] - 2\frac{M_W^4}{m_t^4} \right] , \quad (9)$$

$$\Gamma(t \to b\tau\nu) = \frac{g^2 M_W^2}{48\pi^2} \int \frac{dm_X^2 \Gamma(t \to bX^+)}{(m_X^2 - M_W^2)^2 + (\Gamma_W M_W)^2} , \quad (10)$$

where X^+ now denotes an off-shell W^+ and $\Gamma(t \to bX^+)$ is defined by Eq. (9) with $M_W^2 \to m_X^2$. Analogous formulas with color factors apply to $t \to bc\bar{s}$ decays. With the above formulas we have calculated the branching fractions for top decay including both W^+-mediated and H^+-mediated modes.

The first question of interest is the degree to which competition from the $t \to bH^+$ modes suppresses the standard $t \to bW^*$ signals. Figure 1 shows the $t \to bW^* \to$(all final states) branching fraction versus $m_{H^\pm}$, for various top-quark mass choices with $\tan\beta = \frac{1}{2}, 1, 2, 4$. The conventional $t \to bW^*$ signals are

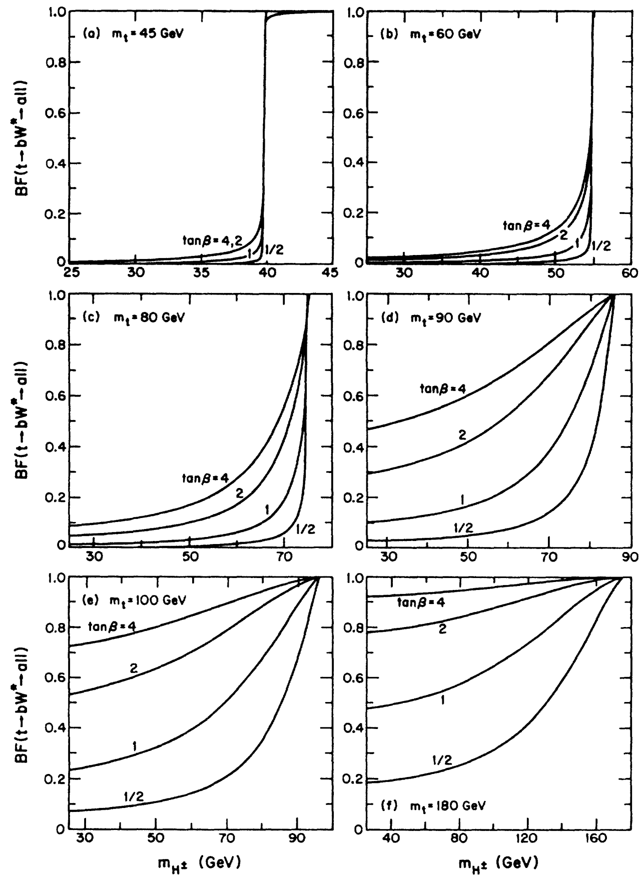

FIG. 1. The branching fraction for standard W-mediated top decays,

$$B(t \to bW^* \to \text{all}) = \frac{\Gamma(t \to bW^* \to \text{all})}{[\Gamma(t \to bW^* \to \text{all}) + \Gamma(t \to bH^+ \to \text{all})]} ,$$

is shown vs the charged-Higgs-boson mass $m_{H^\pm}$ for $\tan\beta = \frac{1}{2}, 1, 2, 4$. The top-quark masses in GeV are (a) 45, (b) 60, (c) 80, (d) 90, (e) 100, and (f) 180.

only suppressed when $m_{H^\pm} < (m_t - m_b)$; decays via a virtual $H^\pm$ are essentially never competitive. Subject to this mass condition, there is a strong suppression of $t \to bW^*$ modes for $m_t \sim 45$–60 GeV essentially regardless of the $\tan\beta$ value, as shown in Figs. 1(a) and 1(b). In the region $m_t \sim M_W$ this suppression is still present but is less severe since $t \to bW^*$ modes have resonant enhancement; see Figs. 1(c) and 1(d) for $m_t = 80$ and 90 GeV. For $m_t > m_W + m_b$, the H modes may still dominate over the W^* modes but only if $m_{H^\pm}$ is well below $m_t - m_b$ and $\cot\beta$ is large [to enhance the tbH coupling, see Eqs. (4) and (6)]; Figs. 1(e) and 1(f) illustrate the cases $m_t = 100$ and $m_t = 180$ (the upper limit allowed by radiative correction analyses[8]).

The next question is whether $t \to bH^+$ modes can be experimentally identified, in cases where $t \to bW^*$ modes are suppressed. At $p\bar{p}$ colliders, only the $H^+ \to \tau^+\nu$ modes have good potential for identification,[4] using high missing-p_T and isolated τ-jet signatures principally; $H^+ \to c\bar{s}$ modes would be very difficult to separate from

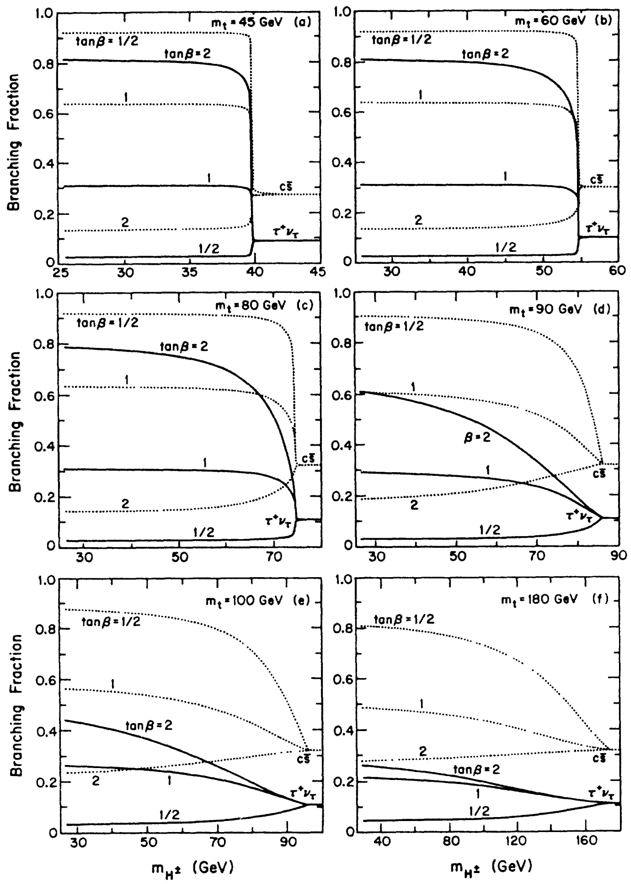

FIG. 2. The net top-quark branching fractions $B(t \to b\tau^+\nu)$ and $B(t \to bc\bar{s})$ of the two modes characteristic of $t \to bH^+$ decays are shown vs $m_{H^\pm}$ for $\tan\beta = \frac{1}{2}1,2$. Both $t \to bW^*$ and $t \to bH^+$ contributions are included. The top-quark masses in GeV are (a) 45, (b) 60, (c) 80, (d) 90, (e) 100, and (f) 180. Solid curves denote $B(t \to b\tau^+\nu)$, dashed curves denote $B(t \to bc\bar{s})$.

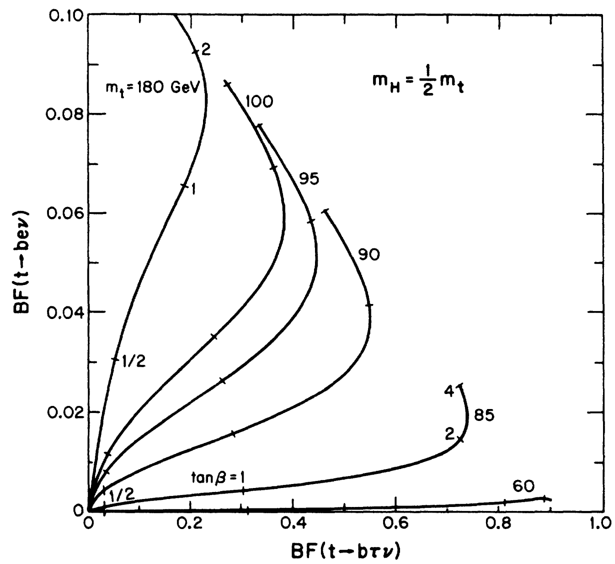

FIG. 3. The correlation of the top-quark branching fractions $B(t \to be\nu)$ and $B(t \to b\tau\nu)$ for the leptonic modes that are identifiable at hadron colliders, for a range of m_t and $\tan\beta$ values with $m_{H^\pm} = \frac{1}{2}m_t$. Both $t \to bW^*$ and $t \to bH^+$ contributions are included. Note that $B(t \to b\mu\nu) \simeq B(t \to be\nu)$ but $B(t \to b\tau\nu)$ is not equal to these branching fractions when $t \to bH^+$ contributes.

charm hadroproduction backgrounds. At e^+e^- colliders the production of a heavy quark can be detected via an excess of high-sphericity events; identification as $t \to bH$ decay can be made for both $H^+ \to \tau^+\nu$ and $H^+ \to c\bar{s}$ modes, using microvertex detectors and invariant-mass peaks. (Of course $t \to W^* \to \tau\nu$ and $t \to W^* \to c\bar{s}$ modes also exist, but the special feature of $t \to bH^+$ decays is that $\tau\nu$ and $c\bar{s}$ dominate and all other channels are suppressed.)

Figure 2 presents the net $t \to b\tau^+\nu$ and $t \to bc\bar{s}$ branching fractions (including both W^* and H contributions) versus $m_{H^\pm}$. In all $m_t, m_{H^\pm}$ mass scenarios, a substantial $B(t \to bH^+ \to \tau^+\nu)$ can be obtained only if $\tan\beta \gtrsim 1$, as we expect from Eq. (5). For $m_t \lesssim M_W$ and $\tan\beta \gtrsim 2$, very large branching fractions can result. For $m_t > M_W$ the $t \to bH^+$ channel is competitive only if $\tan\beta \lesssim \frac{1}{2}$ [because $\Gamma(t \to bW)$ is large for real W and $\Gamma(t \to bH^+) \sim \cot^2\beta$], but then $B(t \to b\tau\nu)$ cannot be large. This latter scenario shows the possibility that all the identifiable decay modes (both $t \to bW^* \to be\nu$ or $b\mu\nu$ and $t \to bH^+ \to b\tau\nu$) may be simultaneously suppressed, leaving top quarks hidden at hadron colliders.

To summarize the possible undetectability of top via $t \to be\nu$ $(b\mu\nu)$ or $t \to b\tau\nu$ decays, we show in Fig. 3 a plot of the correlation of the two branching fractions for various m_t and $\tan\beta$ values (taking $m_{H^\pm} = \frac{1}{2}m_t$ as a representative case). The lower left corner of this plot is the region of maximum undetectability. For example, top-quark detection would be extremely difficult if $B(t \to be\nu) \lesssim 2\%$ (compared to 10% for the standard model) and simultaneously $B(t \to b\tau\nu) \lesssim 18\%$ (judging from the UA1 experiment[9] where the efficiency for detecting $W \to \tau\nu$ was only one-ninth of that for detecting $W \to e\nu$ events). For any m_t, there is a range of $m_{H^\pm}$ and $\tan\beta$ for which the top signals would be in this hidden region; e.g., $m_{H^\pm} \lesssim 0.5m_t$ with $\tan\beta < 0.8$ for $m_t < 95$ GeV, or with $\tan\beta < 0.4$ for $m_t < 180$ GeV. Given such a choice of $\tan\beta$, the signal suppressions are actually much stronger at lower values of m_t, partially compensating for the larger top-quark production rates here. We note incidentally that $\tan\beta < 1$ may be favored for light top quarks, since charged-Higgs-boson box diagrams can then explain the observed B_d^0-$\bar{B}_d^0$ oscillations[10] (see Glashow and Jenkins[3]).

We therefore conclude the following.

(i) The competition of H-mediated and W-mediated decays is a complicated question, depending on several parameters $m_{H^\pm}$, m_t, and $\tan\beta$. However, Figs. 1–3 give a clear overall view of the possible physical outcomes and their relation to the parameter values.

(ii) So long as the charged-Higgs-boson mass $m_{H^\pm} < m_t - m_b$, it is always possible for the standard

$t \to be\nu$ (or $b\mu\nu$) semileptonic top-quark signals to be suppressed. For $m_t \sim 45-60$ GeV strong suppression occurs for a wide range of $\tan\beta$. For larger values of m_t, suppression may require $\tan\beta$ to be small when $m_t > M_W$.

(iii) The observable leptonic signature $t \to bH^+ \to b\tau^+\nu$ of the charged-Higgs-boson-mediated decays may be simultaneously suppressed, leaving the top quark essentially undetectable at $p\bar{p}$ colliders. This occurs, for example, if $\tan\beta \lesssim 0.4$ and $m_{H^\pm} \lesssim \frac{1}{2}m_t$ for any $m_t < 180$ GeV.

(iv) An inequality of the branching fractions $B(t \to be\nu)$ and $B(t \to b\tau\nu)$ is a characteristic consequence of $t \to bH^+$ contributions.

This research was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, and in part by the U.S. Department of Energy under Contract No. DE-AC02-76ER00881.

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