Singlet charge $\frac{2}{3}$ quark hiding the top quark: Fermilab Tevatron and CERN LEP implications

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If c and t quarks are strongly mixed with a weak singlet charge $\frac{2}{3}$ quark, $B(t \to \ell\nu + X)$ could be suppressed via the $t \to cH^0$ mode; thereby, the top quark could still hide below M_W, whereas the heavy quark signal observed at the Fermilab Tevatron is due to the dominantly singlet quark Q. This may occur without affecting the small m_c value. Demanding $m_Q \simeq 175$ GeV and $m_t \lesssim M_W$, we find that $B(t \to \ell\nu + X)$ cannot be too suppressed. The heavy quark Q decays via W, H, and Z bosons. The latter can lead to b-tagged $Z + 4$ jet events, while the strong c-Q mixing is reflected in a sizable $Q \to sW$ fraction. $Z \to t\bar{c}$ decay occurs at the tree level and may be at the 10^{-3} order, leading to the signature of $Z \to \ell\nu b\bar{c}$, all isolated and with large p_T, at 10^{-5} order.

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I. INTRODUCTION

Recently, the Collider Detector at Fermilab (CDF) Collaboration has reported [1] some evidence for the production of heavy quarks with mass of order 174 GeV at the Tevatron. The most likely explanation is, of course, the standard model (SM) top quark. However, at present, in principle it is still possible [2] that the signal is due to some other heavy quark, whereas the actual top quark is hiding below M_W. This is because the top quark semileptonic branching ratio (BR) has not yet been measured. If, for some reason, $B_{\rm sl} \equiv B(t \to b\ell\nu) \ll \frac{1}{9}$, the SM expected value, the top quark may have evaded detection. This can arise basically only through scalar-induced interactions [2].

One such scenario was proposed [3] earlier by Mukhopadhyaya and Nandi (MN). Following a suggestion [4] by Barbieri and Hall (BH), MN considered the existence of an SU(2) singlet charge $\frac{2}{3}$ quark Q along side SM fermions. Since the Glashow-Iliopoulos-Maiani (GIM) mechanism is broken, the mixing of Q with up-type quarks induces tree-level flavor-changing neutral couplings of the SM Higgs boson. If $m_{H^0} < m_t$, $t \to cH^0$, transitions may trigger the aforementioned mechanism of suppressing $B_{\rm sl}$. In a subsequent paper, facing criticisms of "naturalness" [5, 6] ($1/m_Q$ suppressions of heavy Q effects), MN retracted, and considered $t \to cH^0$ dominance to be not very likely [7]. In this paper we study the precise conditions in which $t \to cH^0$ dominance can be realized. We find that this requires Q to be strongly mixed with *both* charm and top quarks, which can occur even with a small m_c eigenvalue. However, we find that although $t \to cH^0$ can be dominant, it is unlikely to be overwhelmingly dominant. Thus, $t \to bW^*$ should occur at a reduced but still substantial fraction. This offers hope that, even if $m_t < M_W$, the top quark can be uncovered at the Tevatron by a renewed study with existing data.

The heavy quark Q can decay both via W and Z bosons [8]; hence it could be the heavy quark observed by CDF. Although one could not explain the larger than expected cross section for 174 GeV quarks, one could plausibly account for the b-tagged Z+ 4 jets events [1]. Eventually, events with ZZ+ 2 jets should start to emerge with increased luminosity [9]. Another point of great phenomenological interest is $Z \to t\bar{c}$ decays, first stressed in this context by BH [4]. These occur at the tree level again because the GIM mechanism is violated. Although widely known, the possibility apparently has not been studied with actual data from the CERN e^+e^- collider LEP because of the standard expectation of a very heavy top quark. We estimate that $Z \to t\bar{c}$ could occur at the 10^{-3} level [4], but *with $B(t \to b\ell\nu)$ of order a few percent.* This results in a signal branching ratio of $Z \to \ell^+\nu b\bar{c}$ at the few $\times 10^{-5}$ level, and each LEP experiment could have a few tens of events at present. The background level can probably be managed, and LEP experiments are strongly urged to conduct such a search.

II. SINGLET-QUARK-INDUCED COUPLINGS

In addition to the standard u-type quarks u^0_{iL}, u^0_{iR} (i = 1–3), we add a left-right singlet charge 2/3 quark u^0_{4L}, u^0_{4R}. The left-handed singlet field u^0_{4L} can pair up with the four right-handed fields to form gauge-invariant singlet masses, which we denote as M'_i and M, respec-

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tively. The right-handed singlet field u_{4R}^0 introduces three extra Yukawa couplings, resulting in off-diagonal masses which we denote as m'_i. Thus, the u-type quark mass matrix is

$$\boldsymbol{M} = \boldsymbol{Y} + \boldsymbol{S} = \begin{bmatrix} m & m' \\ M' & M \end{bmatrix}, \tag{1}$$

where

$$\boldsymbol{Y} = \begin{bmatrix} m & m' \\ 0 & 0 \end{bmatrix}, \quad \boldsymbol{S} = \begin{bmatrix} 0 & 0 \\ M' & M \end{bmatrix} \tag{2}$$

are Yukawa and singlet masses. $\boldsymbol{M}$ is diagonalized by a biunitary transform,

$$\overline{\boldsymbol{M}} = U^\dagger \boldsymbol{M} U' = \mathrm{diag}\,(\bar{m}_u,\ \bar{m}_c,\ \bar{m}_t,\ \bar{M}_Q), \tag{3}$$

where, departing from the notation of MN [3],

$$U^\dagger = \begin{bmatrix} K & x^\dagger \\ y^\dagger & z^* \end{bmatrix} \tag{4}$$

and K is a 3×3 matrix. The Yukawa matrix $\boldsymbol{Y}$ is not simultaneously diagonalized,

$$\overline{\boldsymbol{Y}} = U^\dagger \boldsymbol{Y} U' = \overline{\boldsymbol{M}} - U^\dagger \boldsymbol{S} U', \tag{5}$$

and the off-diagonal term controls flavor-changing neutral current (FCNC) H^0 and Z^0 couplings. The apparent freedom due to the presence of a U' rotation matrix on right-handed fields led MN originally to conclude that tcH^0 coupling could easily be rather large. However, from Eq. (3), simple algebra gives

$$-U^\dagger \boldsymbol{S} U' = -\begin{bmatrix} x^\dagger x \overline{\boldsymbol{m}} & x^\dagger z M_Q \\ z^* x \overline{\boldsymbol{m}} & z^* z M_Q \end{bmatrix}, \tag{6}$$

where $\overline{\boldsymbol{m}}$ is the diagonal 3×3 mass matrix [see Eq. (3)]. We see that no reference to U' is left, and the off-diagonal couplings depend only on mass eigenvalues and Q-related mixing elements of U [6,7]. The relevant flavor-changing Higgs couplings are [6] ($i \neq i'$)

$$\begin{aligned} &-\left(\bar{m}_i x_{i'}^* x_i\ \bar{u}_{i'L} u_{iR} + \bar{m}_{i'} x_i^* x_{i'}\ \bar{u}_{iL} u_{i'R}\right) \frac{H}{v}, \\ &-\left(m_i z^* x_i\ \bar{Q}_L u_{iR} + m_Q x_i^* z\ \bar{u}_{iL} Q_R\right) \frac{H}{v}. \end{aligned} \tag{7}$$

The FCNC Z couplings are [3]

$$\begin{aligned} &\frac{g}{2\cos\theta_W}\, x_{i'}^* x_i\ \bar{u}_{i'L} \gamma_\mu u_{iL}\, Z^\mu + \text{H.c.}, \\ &\frac{g}{2\cos\theta_W}\, x_i^* z\ \bar{u}_{iL} \gamma_\mu Q_L\, Z^\mu + \text{H.c.}, \end{aligned} \tag{8}$$

which is simply related to the Higgs couplings. The charged current becomes

$$\begin{aligned} &\frac{g}{\sqrt{2}} V_{ij}\ \bar{u}_{iL} \gamma_\mu d_{jL}\, W^\mu + \text{H.c.}, \\ &\frac{g}{\sqrt{2}} y'_j\ \bar{Q}_L \gamma_\mu d_{jL}\, W^\mu + \text{H.c.}, \end{aligned} \tag{9}$$

where

$$V \equiv K\, U^{(d)}, \qquad y'_j \equiv y_i^*\, U_{ij}^{(d)}. \tag{10}$$

The 3×3 Kobayashi-Maskawa (KM) matrix V is no longer unitary. Both V and y' depend on the 3×3 left-handed down quark rotation matrix $U^{(d)}$.

III. DETAILS

We wish to explore the range of parameter space where tcH coupling could be sizable. To this end we make a special choice of basis to focus on the problem. First, we choose u_R fields such that $M' = 0$ in $\boldsymbol{S}$. Second, we choose u_{iL}^0, $i = 1$–3, such that the matrix m is diagonal; hence the KM matrix largely comes from the down-type quark sector (we have checked that it is not possible to generate the observed KM matrix structure just by introducing u-type singlet quarks). Only the charged current is affected by the d-type quark sector, while the FCNC Higgs and Z couplings depend only on x_i and z.

The u-type quark mass matrix is now in the form

$$\boldsymbol{M} = \begin{bmatrix} m_1 & 0 & 0 & \Delta_1 \\ 0 & m_2 & 0 & \Delta_2 \\ 0 & 0 & m_3 & \Delta_3 \\ 0 & 0 & 0 & M \end{bmatrix}. \tag{11}$$

The relevant freedom introduced by the singlet quark Q is parametrized as three new off-diagonal Yukawa terms, plus the diagonal, gauge-invariant Dirac mass M. The parameters x_i, y_i, and z can be found by diagonalizing $\boldsymbol{M}\boldsymbol{M}^\dagger$. Without loss of generality, we set $\Delta_1 = 0$, and so the u quark decouples from our discussion.

To illustrate the correlation between $\hat{m}_i \equiv m_i/M$ and $\hat{\Delta}_i \equiv \Delta_i/M$, we set $\Delta_2 = 0$ and plot, in Fig. 1, the mass eigenvalues m_t/M, M_Q/M vs $\hat{m}_3$ for different $\hat{\Delta}_3$ values. Level repulsion is evident: $m_t < m_3$ and $M_Q > M$ for m_3, $\Delta_3 < M$. For larger m_3, Δ_3 values, we adopt the convention that, if $x_t > 0.5$, the heavier state is defined as the top quark. Thus, Fig. 1 depicts both the mass eigenvalues and the label for t and Q.

We are more interested in the effect of Δ_2. With finite Δ_2, but negligible m_1, Δ_1, m_2, the heavy mass eigenvalues are

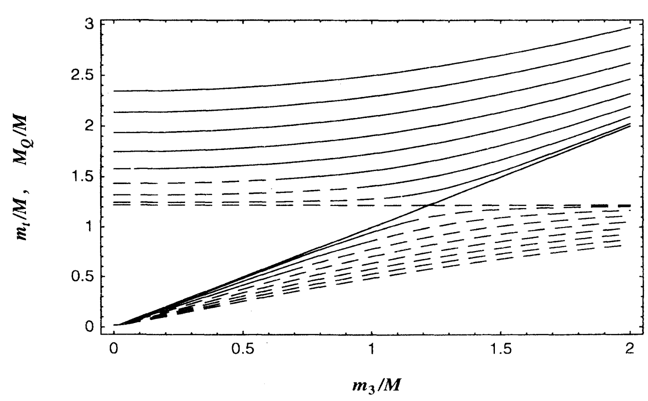

FIG. 1. m_t/M, M_Q/M vs $\hat{m}_3$ for $\hat{\Delta}_2 = 0$ and $\hat{\Delta}_3 = 0$–2 in equal intervals. The solid (dashed) line stands for the physical t (Q) quark. The intersecting straight lines are for $\hat{\Delta}_3 = 0$ while larger $\hat{\Delta}_3$ pushes the eigenvalues apart.

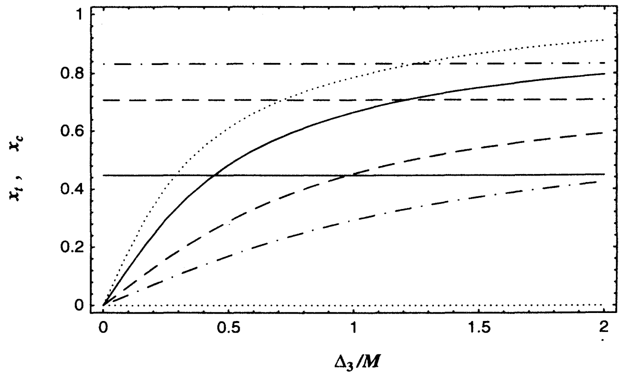

FIG. 2. Mixing parameters x_t, x_c vs $\hat{\Delta}_3$ for $\hat{m}_3 = 0.7$ and $\hat{\Delta}_2 = 0$ (dotted line), 0.5 (solid line), 1 (dashed line), 1.5 (dot-dashed line). Curves for x_c are close to straight lines.

$$m_t^2,\ M_Q^2 = \frac{\Sigma^2 \mp \sqrt{\Sigma^4 - 4m_3^2\,(M^2 + \Delta_2^2)}}{2}, \qquad (12)$$

where $\Sigma^2 = M^2 + m_3^2 + \Delta_3^2 + \Delta_2^2$. For the sake of discussion, we consider the case where m_i, $\Delta_i < M$ (top lighter). Note that in Fig. 1 when $\hat{\Delta}_3$ is not too large, the top mass eigenvalue is close to the diagonal term m_3. This is a generic feature. When other Δ's can be ignored and $\hat{\Delta}_i$ is not too big, the mass eigenvalue and mixing are roughly

$$\bar{m}_i^2 \sim \frac{m_i^2}{1 + \hat{\Delta}_i^2}, \qquad x_i \sim \hat{\Delta}_i. \qquad (13)$$

These relations become affected only when there are *two* $\hat{\Delta}_i$ values that are sizable, which follows largely as a consequence of unitarity of the 4×4 matrix U. In Fig. 2 we plot x_t, x_c as a function of $\hat{\Delta}_3$ for $\hat{m}_3 = 0.7$ and $\hat{\Delta}_2 = 0$, 0.5, 1, 1.5. Note the remarkable feature that x_c is almost independent of $\hat{\Delta}_3$, but when x_c is large, x_t is suppressed for low $\hat{\Delta}_3$ values because of unitarity. The physical reason for this can be traced back to the fact that $V_{cb} \sim 0.04$ is very small compared to 1, and that m_c is small.

Thus, the eigenvalue $\bar{m}_c$ could be made small by choosing a small value for m_2, but this does not forbid Δ_i from being sizable. This is precisely *counter* the "hierarchy principle" [10] advocated in Ref. [6]. However, other than being a prejudice, there is really no reason why Δ_2 cannot be large, since it is an independent parameter. Of course, if $\Delta_2 \sim m_2$, then the conclusions of Ref. [6] would hold.

IV. PHENOMENOLOGY

Inspection of Eq. (7) suggests that $t_L \to c_R$ transitions are suppressed by m_c/v [6], but $t_R \to c_L$ transitions have the effective coupling $m_t\, x_c^* x_t/v$. Since m_t/v is not small, as long as $|x_c x_t|$ is not too suppressed, the $t \to cH^0$ mode has good probability to be dominant over $t \to bW^*$ [3]. The necessary condition is therefore that both x_c and x_t are sizable and neither are suppressed. Hence, *Q, t, and c all become rather arbitrarily mixed* although the charm mass is fixed by m_2. Such an unusual situation is bound to have unusual consequences beyond $t \to cH^0$ being sizable.

It is possible to have $t \to cH^0$ dominance, *and* at same time account for CDF's apparent observation of a heavy quark (whether dominantly doublet or singlet) of mass 174 GeV. The parameters could vary over wide range, subject to the condition that $\hat{\Delta}_2$, $\hat{\Delta}_3$ are sizable, and m_3, M should be chosen such that $m_t \lesssim M_W$ and $m_Q \sim$ 174 GeV. With these conditions, the mechanism largely depends on phase space for $t \to cH^0$. For illustration, we choose $\hat{\Delta}_2$, $\hat{\Delta}_3 = 0.7$, 0.75, $M = 110$ GeV, vary m_3 (to get m_t, M_Q, etc.), and plot, in Fig. 3, $B(t \to cH^0)$ vs m_t (the physical mass) up to 90 GeV, for $m_H = 50$–75 GeV. We allow for m_H below the present LEP bound in case there is more than one Higgs doublet [12]. In producing Fig. 3, we compute the $t \to cH^0$ and bW^* decay width using the couplings of Eqs. (7) and (9). We assume that $U^{(d)}$ amounts to a "small" rotation close to the "standard" KM matrix, ignoring all phases. We have also ignored $t \to cZ^*$ decay as this is a three-body process subdominant compared to $t \to bW^*$. It is clear that, if the Higgs boson mass is sufficiently light, $t \to cH^0$ can be dominant. However, the combined demand of $m_Q \simeq 174$ GeV and $m_t \lesssim M_W$ dictates that the $t \to cH^0$ mode cannot be overwhelmingly dominant. Thus, although suppressed, $B(t \to bW^*)$ should not be vanishingly small. For larger m_H, $t \to cH^0$ dominance quickly fades, and the possibility is ruled out by CDF [13] and D0 [14] Collaborations, since $t \to bW^*$ is not drastically suppressed. In the following, we shall assume that one works in the domain where $B_{\rm sl}$ for the "top quark" (it could be the dominantly singlet quark, since we do not know the scale for M) is suppressed by 1/3 or more. That is, $B_{\rm sl} < 1/27$.

For m_t, $M_Q = 75$, 174 GeV, the parameters of Fig. 3 corresponds to x_c, $x_t \simeq 0.59$, 0.53. Taking $m_H = 60$ GeV, the corresponding $Q \to bW$, sW, tH, cH, tZ, cZ branching ratios are 0.51, 0.24, 0.04, 0.1, 0.02, 0.09,

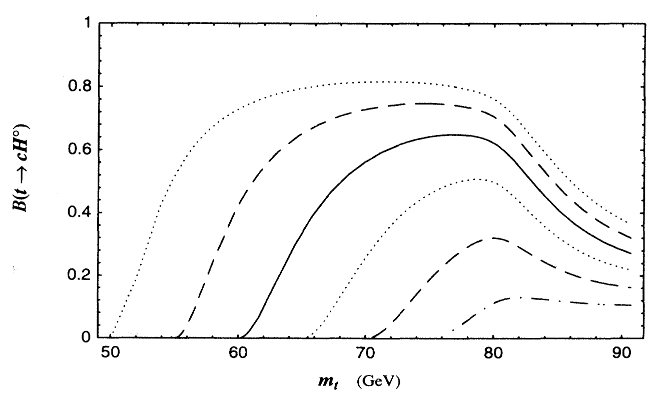

FIG. 3. $B(t \to cH^0)$ vs m_t for $\hat{\Delta}_2$, $\hat{\Delta}_3 = 0.7$, 0.75, $M =$ 110 GeV, and m_H =50–75 GeV in 5 GeV intervals (from left to right).

respectively. Note that, as a consequence of large x_c, $Q \to sW$ decay has a sizable rate. The modes $Q \to tH^0$ and tZ are suppressed by phase space, while $Q \to cH^0$ and cZ are suppressed by an extra power of $|x_c|^2$. Thus, W-induced decays are still dominant, but the b content in the final state is diluted slightly by the $Q \to sW$ mode. Although one cannot account for the large production cross section for the heavy quark (*one could always add another singlet u-type quark for this purpose*), other features reported by CDF can be accounted for [8, 11], in particular, the appearance of b-tagged Z+ 4 jet events. The Z boson comes from $Q \to cZ$, tZ, while a b tag could come from $Q \to bW$ or from $t \to bW^*$ or $H \to b\bar{b}$, etc., in subsequent decays. This could be of order 20% of the $Q\bar{Q}$ cross section, hence consistent with what is observed. On the other hand, single lepton events with a b tag are slightly depressed ($\sim$ 70%) compared to the standard heavy top quark. We therefore conclude that the heavy quark observed by CDF may well be a doublet-singlet mixed state Q. The scenario offers many signatures that can be checked experimentally. The so-far "hidden" light "top quark" with $B_{\rm sl}$ not too suppressed compared to standard model perhaps can be probed with existing Tevatron data [2].

The scenario has a consequence that may be studied at LEP. As first pointed out by Barbieri and Hall [4], $Z \to t\bar{c}$ can be quite sizable with the existence of charge 2/3 singlet quarks. Using Eq. (8) and x_c, x_t values for the example above, we estimate that $B(Z \to t\bar{c} + \bar{t}c)$ is of order a few $\times 10^{-3}$, which is consistent with Ref. [4]. Other phenomenological constraints are not particularly stringent, and can be found in Ref. [4]. For example, the D^0-$\bar{D}^0$ mixing constraint can be satisfied with small Δ_1. Since $B(t \to \ell\nu X)$ does not vanish, we estimate that the potentially observable signal of $Z \to t\bar{c} \to \ell\nu+$ 2 jets (where the jets contain b and c quarks, respectively) could have a branching ratio of order a few $\times 10^{-5}$. Since the lepton and neutrino should be well isolated with sizable (15–20 GeV) p_T or missing energy (they are *bona fide* virtual W decay events), and that ℓ, ν, and one of the jets should pair up to be the top mass, there should be sufficient handles to suppress the background. The latter presumably comes from events with $Z \to b\bar{b}$ plus gluon bremsstrahlung.

In summary, with the addition of a weak singlet charge 2/3 quark, it is possible to have c, t, and Q all strongly mixed with each other while $m_c \simeq 1.5$ GeV is maintained. We arbitrarily label the quarks such that $m_t < m_Q$. If the Higgs boson is lighter than the top quark, the singlet content of t and Q could allow the top quark to be lighter than the W boson but decaying dominantly via $t \to cH^0$, thereby remaining hidden at the Tevatron. The heavy quark Q would decay via W, H, and Z bosons, with charged current still the dominant agent. Thus, the CDF observation of a heavy quark with mass of order 174 GeV can be accounted for, while the H- and Z-induced decays could explain some of the apparent oddities. An important consequence of strong (GIM rule breaking) c-t-Q mixing is the presence of a sizable $Q \to sW$ branching fraction. Another consequence is *tree-level Z-t-c* couplings. Since $m_t < M_W$ remains possible, besides studies at the Tevatron, one could also search for $Z \to t\bar{c}$ decays at LEP. The signal mode of $Z \to t\bar{c} \to \ell\nu b\bar{c}$ is expected at 10^{-5} order.

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