Technicolor contribution to the rare decays of heavy quarks

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In this paper we calculate some rare decays of heavy quarks (b and t) using the ideas of technivector meson dominance and TC spectral functions. The low energy ETC dynamics that gives rise to the large mass of the top quark inevitably induces such FCNC decays of b and t quarks. We find that in an $SU(4)_{TC}$ model with 1 and 10 doublets of technifermions the $b \to s\gamma$ decay rate, for $m_t \approx 150$ GeV, will be modified by 11% and 26%, respectively, from the SM value, whereas for the same top mass the branching ratios for the rare decays $t \to c\gamma$ and $t \to cZ$ get contributions of $\sim 10^{-10} - 10^{-9}$ and $\sim 10^{-6}$ from the ordinary TC scenario. These values are greater than the corresponding SM predictions by 2–3 and 7 orders of magnitude, respectively. In particular, for TC models the rare decay $t \to cZ$ has a good prospect of being observed at the CERN LHC. The SSC, which has recently been abandoned, would also have produced an observable number of such decays.

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I. INTRODUCTION

Rare decays of heavy quarks have been under intense theoretical and experimental study for quite some time. The recent observation [1] of the electromagnetic penguin diagram decay $B \to K^*\gamma$ has stimulated more interest in this field. There are many reasons behind this. First, many rare decays are sensitive to various parameters of the standard model (SM). For example, the mass of the charm quark was estimated from the K_L-K_S mass difference by using the Glashow-Iliopoulos-Maiani (GIM) mechanism [2] and it proved to be approximately correct when the charm quark was later discovered. Second, the smallness of rare decay rates within the context of the SM makes them excellent probes for the effects of new physics beyond the SM, such as technicolor (TC), supersymmetry, multi-Higgs-boson models, or extra fermions. In fact, some rare decays such as $K \to \pi\nu\bar{\nu}$ and $B \to K^*\gamma$ can receive large enhancements from physics beyond the SM. Finally, heavy quark decays are expected to be short distance interaction dominated and in particular the decays of the top quark should be relatively free from hadronic uncertainties and unambiguously calculable. All these reasons justify the use of rare decays of heavy quarks as a unique window in the search for effects of new physics.

TC is a dynamical mechanism [3] for breaking the electroweak (EW) gauge symmetry and giving mass to the weak gauge bosons. In TC models the elementary Higgs sector of the SM is replaced by composite Goldstone bosons made out of technifermions and antitechnifermions. However, to give mass to ordinary fermions one has to introduce extended technicolor [4] (ETC) which gives rise to transitions between ordinary fermions and technifermions. Light fermions (e.g., u, d, s) get their masses from ETC interactions at higher energies and heavy fermions (e.g., b, t) presumably get their masses from ETC dynamics at lower energies [5]. In particular, a top mass of 150 GeV can arise from ETC dynamics at a few TeV. In this paper we shall show that the same low energy ETC dynamics can give rise to rare decays of b and t quarks at sizable rates.

Recently CLEO [1] has reported the observation of the exclusive EM penguin diagram decay $B \to K^*\gamma$. The observed value of $(4.5 \pm 1.5 \pm 0.9) \times 10^{-5}$ is consistent with the theoretical prediction of $(1-15) \times 10^{-5}$ within the experimental and theoretical uncertainties involved in measuring and estimating exclusive branching ratios. The uncertainties involved, however, leave ample room for contributions from beyond SM physics. CLEO has also set upper and lower bounds on the total inclusive rate of $0.4 \times 10^{-4} < B(b \to s\gamma) < 19 \times 10^{-4}$. These bounds are expected to improve in the near future and can be used to severely constrain the effects of new physics such as technicolor, supersymmetry, or extended Higgs sectors [6]. Although CLEO has not reported the total inclusive decay rate for $b \to s\gamma$, the prospects for its measurement are quite promising. The Collider Detector at Fermilab (CDF) Collaboration has recently placed a lower bound of 113 GeV on m_t by using the dilepton channel as predicted by the SM [7]. Further, an indirect value of $m_t = 164^{+16+18}_{-17-21}$ GeV has been estimated by comparing precision measurements of the Z boson properties at the CERN e^+e^- collider LEP with calculations of EW radiative corrections within the context of the SM [8]. In this range of values of m_t the branching fractions of the rare decays $t \to c\gamma$ and $t \to cZ$, within the SM, are extremely small and are of the order of 10^{-12} and $10^{-12} - 10^{-13}$, respectively [9]. Although the CERN Large Hadron Collider (LHC) is expected to produce 10^8 $t\bar{t}$ pairs per year, it is clear that the SM-induced rare top decays cannot be observed at the same collider. Similar conclusions also hold for the recently abandoned Superconducting Super Collider (SSC). However, contributions from new physics could make some of them observable. Further, since the top mass is the largest fermion mass scale in the three generation SM, its flavor-changing neutral current (FCNC) decays might hold the key to many puzzling aspects of flavor physics. In this paper we have calculated the TC contribution to the branching ratios of the rare decays $b \to s\gamma$ [10], $t \to c\gamma$, and $t \to cZ$ by using the ideas of technivector meson dominance [11] and TC spectral functions [12]. We find that in an $SU(4)_{TC}$ [13] model with 1 and 10 doublets the $b \to s\gamma$ decay rate, for $m_t \approx 150$ GeV, will be modified by 11% and 26%, respectively, from the SM value, whereas for the same top mass, the branching ratios for the rare decays $t \to c\gamma$ and $t \to cZ$ get contributions of $\sim 10^{-10} - 10^{-9}$, and $\sim 10^{-6}$, respectively. Since our computation does not take into account the full effects of walking dynamics on the technifermion loop, we shall not make any conclusive statement about decay rates for walking technicolor (WTC) relative to ordinary TC. Also we shall not calculate the TC contribution to the decay rate for $b \to sZ$ since the dominant chiral contribution to it has already been correctly estimated in Ref. [10] and we have nothing new to report on it.

The contents of this paper are as follows. In Sec. II we present the ETC-induced fundamental Lagrangian and use the ideas of technivector meson dominance to find the resulting couplings of EW gauge bosons with heavy quarks. The transition $t \to bW^+$ (although it is not a rare decay mode of the top quark) in TC models is then examined in Sec. III. Ordinary TC contributions to the recently observed $b \to s\gamma$ decay rate are presented in Sec. IV. The decay rates for $t \to c\gamma$ and $t \to cZ$ in TC models are considered in Sec. V. Finally, Sec. VI contains a brief discussion of our results and the conclusions of our study.

II. THE ETC-INDUCED LAGRANGIAN AND TECHNIVECTOR MESON DOMINANCE

Consider a technifermion doublet $T_L = (P_L, N_L)$, P_R, N_R which transforms under $SU(2)_L \otimes U(1)_Y$ as $(P_L, N_L) = (2,0), P_R = (1,1)$, and $N_R = (1,-1)$. We shall calculate heavy quark decay amplitudes in minimal TC as well as in TC models with a family of technifermions. Unfortunately, our understanding of strong interactions is based on our experience with QCD, which is an SU(3) gauge theory. Vector meson dominance (VMD), in particular, is also based on QCD. The best we can hope to do is to consider a general $SU(N)_{TC}$ gauge theory, assume QCD-inspired VMD to be valid for it, and determine its parameters by large N_{TC} scaling. For gauge groups other than $SU(N)_{TC}$ it is neither reasonable to assume the validity of QCD-inspired VMD, nor can we determine its parameters. We shall therefore consider only $SU(N)_{TC}$ models in this article. For arbitrary $SU(N)_{TC}$ gauge groups a good walking condition can be achieved by adjusting the number of technifermion doublets N_D so that the TC β function is very small. In this article we shall evaluate all rare decay amplitudes for arbitrary N_{TC} and some properly chosen N_D to achieve walking. However, to present some numerical estimates, in the end we shall specialize to $SU(4)_{TC}$ with N_D=1 and 10 technifermion doublets in the fundamental representation.

Since there are no realistic or phenomenologically viable ETC models, we shall consider ETC-induced transitions between quarks and technifermions only in the four-fermion approximation and avoid reference to any specific ETC model. In a wide variety of ETC models, in the absence of fine-tuning, the top quark mass arises from a light ETC gauge boson. We shall assume that this ETC gauge boson is an $SU(2)_L$ singlet, but carries TC and color and couples to left-handed (LH) and RH currents in the following manner [14]:

$$g_{\rm ETC}\sum_{1}^{N_D}[\xi\bar{\psi}'^{\omega}_{Li}\gamma^{\mu}T^{\alpha x}_{Li}+\xi'\bar{t}'^{\omega}_{R}\gamma^{\mu}U^{\alpha x}_{R}]. \tag{1}$$

Here $g_{\rm ETC}$ is the ETC coupling constant, ξ and ξ' are ETC-induced group theory factors of order 1, w and α are color and TC indices, i and j are $SU(2)_L$ indices, and x is an index for labeling the technifermion doublets. In the above we have assumed that the relevant ETC interaction is $SU(N_D)$ invariant. The primes on ordinary fermion fields indicate that they are gauge eigenstates. The top mass arises from a four-fermion operator that couples the LH and RH currents of Eq. (1). Information about the top mass obtained from high energy colliders can be used to estimate the value of $g^2_{\rm ETC}/M^2_{\rm ETC}$ for the ETC gauge boson under consideration. In contrast, the FCNC amplitudes of heavy quarks b and t will arise from an $SU(N_D)$-symmetric four-fermion operator that couples two LH currents of Eq. (1) in the following manner:

$$L_{\rm eff}=-\frac{g^2_{\rm ETC}}{M^2_{\rm ETC}}\sum_{1}^{N_D}[\hat{\bar{\psi}}'^{\omega}_{Li}\gamma^{\mu}T^{\alpha x}_{Li}\bar{T}^{\alpha x}_{Lj}\gamma_{\mu}\psi'^{\omega}_{Lj}]. \tag{2}$$

[FCNC amplitudes involving the top quark can also arise from the product of two RH currents. However, they will be of the same order of magnitude as those arising from Eq. (2).]

By Fierz transformation the above can be put in the form of a sum of products of TC singlet currents. One of the terms contains a product of $SU(2)_L$ singlet currents and the other a product of $SU(2)_L$ triplet currents. The second term can be further decomposed into a sum of products of charged currents and neutral currents. The final form is [J_a are the generators of SU(2)]

$$\begin{aligned}L_{\rm eff}=-\frac{g^2_{\rm ETC}}{2M^2_{\rm ETC}}\sum_{1}^{N_D}\big[&\bar{\psi}'_L\gamma^{\mu}\psi'_L\bar{T}^x_L\gamma_{\mu}T^x_L\\&+4\bar{\psi}'_L\gamma^{\mu}J_3\psi'_L\bar{T}^x_L\gamma_{\mu}J_3T^x_L\\&+(4\psi'_L\gamma^{\mu}J_+\psi'_LT^x_L\gamma_{\mu}J_-T^x_L+\text{H.c.})\big].\end{aligned} \tag{3}$$

The technifermions associated with the above four-fermion vertex and the vertex of neutral EW gauge bosons will form a closed loop. By integrating out the strongly interacting technifermions in some approximation such as vector meson dominance or naive dimensional analysis (NDA) we can obtain a rough estimate of the TC contribution to the FCNC amplitudes of heavy quarks. In a full ETC model there could possibly be other four-fermion operators that also contribute to the same FCNC amplitudes. But the ETC gauge bosons that lead to these additional operators are expected to be heavier than those associated with the top mass generation. In the absence of any physically compelling reason that could keep these ETC gauge bosons light (1–10 TeV), their contributions to the FCNC amplitudes are expected to be smaller.

Below the confinement scale of TC, the currents carried by the technifermions must be replaced by appropriate SU(2) covariant currents carried by low energy technihadrons. In the technivector meson dominance [11] approach the relevant technihadrons are the technivector (ρ^a) and axial-vector (β^a) mesons. This approach will project out the needed spin-1 channel of low energy technihadrons and simplify the calculations. In addition, the heavy quark decays considered in this article are associated with the production of an EW gauge boson. The mixing of the technivector mesons with the EW gauge bosons as described in the framework of VMD therefore simplifies the calculation. Nevertheless, the use of VMD to integrate out the technifermion loop is only an approximation, the reason being that, at the energy scale where the contribution of vector mesons is important, other parts of the TC spectrum could also contribute significantly but are ignored in VMD. An exact calculation of the technifermion loop would involve integration of the complete TC spectrum, which is beyond the scope of this work. It must be noted that the results derived here could, alternatively, be derived by the rules of NDA which have been generalized for arbitrary SU($N_{\rm TC}$) gauge theories. However, the results derived on the basis of NDA are usually uncertain to $O(1)$ because of strong TC corrections. In addition, NDA, like VMD, has similar limitations towards WTC. For WTC the use of QCD-motivated VMD or NDA to integrate out the technihadrons is not justified, since their properties in a walking scenario might deviate considerably from those of ordinary hadrons. In the absence of any theoretical or phenomenological input we shall blindly apply QCD-inspired VMD also to TC models which contain enough technifermions to achieve walking.

The phenomenological SU(2) currents carried by the vector and axial-vector particles are given by [15]

$$\begin{aligned} gJ^V_\mu &= \partial^\nu F'^V_{\nu\mu} + ig[V'^\nu, F'^V_{\nu\mu}] + ig[A'^\nu, F'^A_{\nu\mu}], \\ gJ^A_\nu &= \partial^\nu F'^A_{\nu\mu} + ig[V'^\nu, F'^A_{\nu\mu}] + ig[A'^\nu, F'^V_{\nu\mu}], \end{aligned} \tag{4}$$

where

$$\begin{aligned} F'^V_{\nu\mu} &= \partial_\nu V'_\mu - \partial_\mu V'_\nu + ig[V'_\nu, V'_\mu] + ig[A'_\nu, A'_\mu], \\ F'^A_{\nu\mu} &= \partial_\nu A'_\mu - \partial_\mu A'_\nu + ig[A'_\mu, V'_\nu] + ig[V'_\nu, A'_\mu]. \end{aligned} \tag{5}$$

In the above equations $V'_\mu = \rho'^a_\mu J^a$ and $A'_\mu = \beta'^a_\mu J^a$. g is the strong interaction strength of flavor SU(2) gauge fields V'_μ and A'_μ. It can be estimated from QCD by large $N_{\rm TC}$ expansion [16], which gives $g^2/4\pi = 3\alpha_{\rho\pi\pi}/N_{\rm TC}$. The primes on the vector and axial-vector meson fields in the above equations indicate that they are eigenstates of isospin SU(2). Note that the above expressions for J^V_μ and J^A_μ are in agreement with the current field identity [11] $\langle 0|J^\nu_\mu|\rho^a\rangle = \varepsilon^a_\mu M_V^2/g$, where M_V is the mass of the vector meson.

Considering only the terms that are linear in V'_μ and A'_μ, the effective low energy Lagrangian describing their interactions with heavy quarks becomes

$$\begin{aligned} L_{\rm eff} = -\frac{g^2_{\rm ETC}\sqrt{N_D}}{gM^2_{\rm ETC}} [\{&\bar\psi'_L\gamma^\mu J_+\psi'_L\partial^\nu(\partial_\nu\rho'^+_\mu - \partial_\mu\rho'^+_\nu) - \bar\psi'_L\gamma^\mu J_+\psi'_L\partial^\nu(\partial_\nu\beta'^+_\mu - \partial_\mu\beta'^+_\nu) + \text{H.c.}\} \\ &+\bar\psi'_L\gamma^\mu J_3\psi'_L\partial^\nu(\partial_\nu\rho'^3_\mu - \partial_\mu\rho'^3_\nu) - \bar\psi'_L\gamma^\mu J_3\psi'_L\partial^\nu(\partial_\nu\beta'^3_\mu - \partial_\mu\beta'^3_\nu) + \tfrac{1}{2}\bar\psi'_L\gamma^\mu\psi'_L\partial^\nu(\partial_\nu\omega'_\mu - \partial_\mu\omega'_\nu)] \,. \end{aligned} \tag{6}$$

When the EW interactions are turned on, the technivector mesons will mix with the EW gauge bosons and this will cause the gauge eigenstates to be orthogonally related to the mass eigenstates. To lowest order in EW gauge couplings we have [11]

$$\begin{aligned} \beta'^\pm &= \beta^\pm - xR^2W^\pm, \\ \rho'^\pm &= \rho^\pm + xW^\pm, \\ \beta'^3 &= \beta^3 - \frac{xR^2}{\cos\theta}Z, \\ \rho'^3 &= \rho^3 + \frac{x\cos 2\theta}{\cos\theta}Z + \frac{e\sqrt{N_D}}{g}A_\gamma, \\ \omega' &= \omega - x\bar Y\sin\theta\tan\theta Z + \frac{e\sqrt{N_D}}{2g}\bar Y A_\gamma, \end{aligned}$$

where R^2 is related to the vector meson mass through a Kawarabayaski-Suzuki-Riazuddin-Fayyazuddin- (KSRF-) [17] like relation $1 - R^2 = g^2F^2_{\rm TC}M^2_V \approx 0.5$, $x = g_L\sqrt{N_D}/2g$, and $\bar Y^2 = (2/N_D){\rm Tr}Y^2_L$. Y_L is the $2N_D \times 2N_D$ hypercharge matrix for the LH technifermions. The unprimed fields on the right-hand side (RHS) stand for mass eigenstates. The ETC-induced effective Lagrangians for the charged and neutral EW gauge bosons become

$$L^{\rm CC}_{\rm eff} = \frac{g^2_{\rm ETC}\sqrt{N_D}}{\sqrt2 gM^2_{\rm ETC}}(1+R^2)x[\partial^\nu(\bar t'_L\gamma^\mu b'_L)W^+_{\nu\mu} + \text{H.c.}] \tag{7}$$

and

$$\begin{aligned} L^{\rm NC}_{\rm eff} = \frac{g^2_{\rm ETC}\sqrt{N_D}}{2gM^2_{\rm ETC}}\Bigg[&\frac{x}{\cos\theta}\{2(R^2+\cos 2\theta)\partial^\nu(\bar\psi'_L\gamma^\mu J_3\psi'_L) \\ &-\bar Y\sin^2\theta\,\partial^\nu(\bar\psi'_L\gamma^\mu\psi'_L)\}Z_{\nu\mu} \\ &+\frac{e\sqrt{N_D}}{g}\Bigg\{\frac{\bar Y}{2}\partial^\nu(\bar\psi'_L\gamma^\mu\psi'_L) \\ &+2\partial^\nu(\bar\psi'_L\gamma^\mu J_3\psi'_L)\Bigg\}F_{\nu\mu}\Bigg], \end{aligned} \tag{8}$$

where $W^+_{\nu\mu} = \partial_\nu W^+_\mu - \partial_\mu W^+_\nu$, $Z_{\nu\mu} = \partial_\nu Z_\mu - \partial_\mu Z_\nu$, and $F_{\nu\mu}$ is the EM field strength tensor. In the above we have done an integration by parts by considering the action functional for the ETC-induced effective Lagrangian. The value of $g^2_{\rm ETC}/M^2_{\rm ETC}$ can be estimated from m_t by using the relation $m_t = -(g^2_{\rm ETC}/M^2_{\rm ETC})\langle\bar{T}T\rangle_\mu \exp[\int_\mu^{M_{\rm ETC}} \gamma_m(\alpha_{\rm rc})\frac{dQ}{Q}]$. Here $\mu \approx 2M_V$ and γ_m is the anomalous dimension of the technifermion condensate. $g^2_{\rm ETC}$ can be estimated by setting it equal to $4\pi\alpha_{\rm TC}(M_{\rm ETC})$. For ordinary TC we shall approximate $\gamma_m(\alpha_{\rm TC})$ by the lowest order term in its perturbative expansion, whereas for WTC $\gamma_m(\alpha_{\rm TC})$ can be approximated by 1 for $\mu \le Q \le M_{\rm ETC}$. For a good walking scenario and low $M_{\rm ETC}$ this is a reasonable approximation. The anomalous scaling factor becomes large for WTC. This effect tends to lower the decay amplitudes for WTC by a small factor of $\frac{1}{3}$ relative to those for ordinary TC. The condensate enhancement effect in this case is rather small, because there is so little room for walking between μ and a low $M_{\rm ETC}$. Note that the estimates of the decay rates given in this article are just order of magnitude ones. This is mainly due to uncertainties connected with ETC-induced group theory factors and approximate evaluation of the technifermion loop. Therefore it does not make sense to attach much importance to small reduction factors of $\frac{1}{3}$ which can easily get buried under the above uncertainties.

III. $t \to bW^+$ IN TC MODELS

In order to get the magnetic moment type operators for heavy quarks from Eqs. (7) and (8) we shall make use of the two identities

$$\bar{\psi}_L\overrightarrow{\partial^\nu}\gamma^\mu\psi_L X_{\nu\mu} = -\tfrac{1}{2}m\bar{\psi}_L\sigma^{\mu\nu}\psi_R X_{\nu\mu} + \cdots \tag{9a}$$

and

$$\bar{\psi}_L\overleftarrow{\partial^\nu}\gamma^\mu\psi_L X_{\nu\mu} = -\tfrac{1}{2}m\bar{\psi}_R\sigma^{\mu\nu}\psi_L X_{\nu\mu} + \cdots\ , \tag{9b}$$

where the ellipsis stands for operators that are nonlinear in the vector field and $X_{\nu\mu}$ is any of the above field strength tensors. Transforming the gauge eigenstates of quark fields into mass eigenstates and using the above identities, the charged current Lagrangian becomes

$$L^{\rm CC}_{\rm eff} = \frac{g^2_{\rm ETC}\sqrt{N_D}}{2\sqrt{2}gM^2_{\rm ETC}}(1+R^2)x[(\bar{U}_{Li}S^{u+}_{i3}m^j_d\sigma^{\mu\nu}S^d_{3j}D_{Rj} + \bar{U}_{Ri}S^{u+}_{i3}m^i_u\sigma^{\mu\nu}S^d_{3j}D_{Lj})W^+_{\mu\nu} + \text{H.c.}]\ . \tag{10}$$

Here $t'_L = S^u_{3i}U_{Li}$, $b'_L = S^d_{3j}D_{Lj}$, $U_L = (u_L, c_L, t_L)$, and $D_L = (d_L, s_L, b_L)$ are mass eigenstates. For the $W\bar{b}t$ vertex we need to put $i = j = 3$. Further, since $m_t \gg m_b$ we can drop terms containing m_b. We then get

$$\delta^2 L_{W\bar{b}_l} \approx \frac{g^2_{\rm ETC}\sqrt{N_D}}{2\sqrt{2}gM^2_{\rm ETC}}m_t(1+R^2) \times x[V_{tb}\bar{t}_R\sigma^{\mu\nu}b_L W^+_{\mu\nu} + \text{H.c.}]\ , \tag{11}$$

where we have set $S^{u+}_{33}S^d_{33} \approx V_{tb}$. The above operator correction is of dimension 5 and it does not give the leading TC contribution to the $W\bar{b}t$ vertex. The leading TC contribution mainly comes from the Goldstone pole in the axial vector channel, the reason being that the diagram to be evaluated consists of a closed technifermion loop with a W line attached to one point and b and t lines attached to another point. The amplitude could therefore be expressed in terms of vector and axial-vector two-point functions for TC which can be modeled following QCD. The contribution of the Goldstone pole in the axial vector channel turns out to be many times greater [12] than the contribution arising from other parts of the TC spectrum in both channels. For ordinary TC and $m_t \approx 150$ GeV it is given by $\delta L_{W\bar{b}t} \approx 0.024 L_{W\bar{b}t}$ [12,14]. Here $L_{W\bar{b}t}$ is the SM contribution to the $W\bar{b}t$ vertex at tree level. Hence to probe the ETC-induced correction to the $W\bar{b}t$ vertex, collider measurements of the same must reach a precision of about 1% for ordinary TC.

The decay amplitudes of WTC can differ from those in ordinary TC for three reasons. First, the values of $g^2_{\rm ETC}/M^2_{\rm ETC}$ estimated from m_t tend to be slightly smaller for WTC because for the condensate enhancement effect. This effect tends to lower the decay amplitudes for WTC. Secondly, WTC models in general contain a large number of technifermion doublets. The decay amplitudes therefore become proportional to N_D and tend to increase for WTC. Finally, the value of the technifermion loop per doublet in WTC may differ from ordinary TC because the TC spectrum may be different in the two cases. The last effect of walking dynamics will not be considered in this article. It should be noted that, although the ETC-induced correction enhances the coupling constant for the $W\bar{b}t$ vertex relative to that of the SM by a few percent, there could be other competing channels for top decay in these models. For example, in ETC models with extra chiral symmetries, a heavy top quark could decay into a charged pseudoscalar P^+ and b. For large enough m_t the coupling constant of this vertex could be large. So although in ETC models the partial width for $t \to bW^+$ is slightly enhanced, its branching ratio could decrease significantly from the SM value.

IV. $b \to s\gamma$ IN TC MODELS

The recent observation of the EM penguin-mediated rare decay $B \to K^*\gamma$ by CLEO has rekindled the interest in rare B decays. It experimentally confirmed the existence of FCNC loop decays in the context of the SM. The rare decay $b \to s\gamma$ is important from several viewpoints. First, within the SM the top quark contribution is the largest one in the loop diagram. Hence the loop decay provides information about m_t and the relevant Cabibbo-Kobayashi-Maskawa (CKM) matrix elements V_{td}, V_{ts}, and V_{tb}. Second, QCD corrections enhance the decay rate by a factor of 3–5 making them experimentally more accessible. In the absence of any kind of new physics, precise determination of the decay rate could therefore provide a good testing ground for QCD corrections. Finally, the rate for the rare decay $b \to s\gamma$ is one of the most sensitive tests of new physics beyond the SM. Within the SM the branching ratio for $b \to s\gamma$ is $(2–4)\times10^{-4}$ for $m_t \approx 150$ GeV [18]. The major uncertainty comes from QCD corrections. The average CLEO value [1] for $B(B \to K^*\gamma)$ is $(4.5 \pm 1.5 \pm 0.9) \times 10^{-5}$. It is consistent with the theoretical predictions for the exclusive branching ratio which are in the range $(1\text{–}15)\times 10^{-5}$, with the largest uncertainty coming from $f_{K^*} = \Gamma(B \to K^*\gamma)/\Gamma(b \to s\gamma) = (4\text{–}40)\%$ [1] CLEO has also looked for $b \to s\gamma$ inclusive decays and has established an upper bound on $B(b \to s\gamma)$ of 5.4×10^{-4} (90% C.L.). This is within a factor of 2 from the SM prediction. Although the theoretical and experimental values are in agreement at present, the uncertainties involved in both leave sufficient room for large contributions from new physics like TC.

The TC contribution to $b \to s\gamma$ can be obtained from the effective neutral current Lagrangian that we have derived. We consider the down quark sector, transform the quark gauge eigenstates to mass eigenstates, and use the two identities given in Eq. (9). Next we choose the initial and final quark operators to correspond to $b \to s\gamma$ and drop the terms containing m_s. The final result is

$$\delta L_{b\bar{s}\gamma} \approx -\frac{g_{\rm ETC}^2 e m_b N_D N_{\rm TC}(2-\bar{Y})}{96\pi^2 M_{\rm ETC}^2} \times [V_{ts}^* \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu} + \text{H.c.}] \; . \tag{12}$$

In the above equation we have set $S_{23}^{d+} S_{33}^{d} \approx V_{ts}$. Our result agrees with that given in Ref. [10] if and only if we use the QCD value for g. However, one should bear in mind that for QCD-like TC models with large $N_{\rm TC}$, the value of g could be significantly different from the QCD value used. In addition, for WTC we do not know how to estimate the value of g. A large $N_{\rm TC}$ expansion is certainly not reliable for WTC. Note that although the above operator is of dimension 5 it gives the leading TC contribution to the $b \to s\gamma$ transition. The unbroken EM gauge invariance forbids any dimension 4 ($d = 4$) operator correction to the $b\bar{s}\gamma$ vertex. The SM induces an effective Lagrangian at the one-loop level (i.e., without QCD corrections) which is given by [19]

$$L_{b\bar{s}\gamma} \approx -A\left(\frac{m_f^2}{m_W^2}\right)\frac{m_b e}{(4\pi v)^2}[V_{ts}^* \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu} + \text{H.c.}] \; . \tag{13}$$

This operator will have to be corrected for short distance QCD effects. In the above equation A ranges from 0.59 for $m_t \approx 120$ GeV to 0.68 for $m_t \approx 210$ GeV. In particular, for $m_t \approx 150$ Gev, $A \approx 0.63$. The TC contribution to $b \to s\gamma$ will therefore be suppressed relative to the SM value by a factor of order $(\mu N_D^{3/2}/M_{\rm ETC})(1-\bar{Y}/2)m_t/4\pi A F_{\rm TC}$ for a TC model containing N_D doublets of technifermions. Here μ is approximately twice the chiral symmetry-breaking scale for TC. For $SU(4)_{\rm TC}$ with $N_D = 1$ and $m_t \approx 150$ GeV this will give rise to an 11% change in the $b \to s\gamma$ decay rate from the SM value, whereas for $N_D = 10$ the decay rate will be modified by 26% relative to the SM value. The present CLEO precision for measuring $b \to s\gamma$ is, however, not sufficient to detect the above deviations. It should not be inferred from the above that the WTC contribution to $b \to s\gamma$ will necessarily be large. Because although the decay amplitude gets multiplied by N_D, the value of the technifermion loop per doublet could decrease for WTC. In particular, WTC models typically contain a large number of light pseudo Goldstone bosons and their effect on the decay amplitude has to be considered. This comment is also valid for the decay amplitude of $t \to c\gamma$ considered in the next section.

V. $t \to c\gamma$ AND $t \to cZ$ IN TC MODELS

The rare top decay $t \to cV$ where $V = \gamma, Z$ have several attractive features. First, like $b \to s\gamma$ it is sensitive to the CKM parameters V_{tb} and V_{cb}. However, unlike $b \to s\gamma$ it does not receive large enhancements from a heavy fermion in the loop diagram or due to short distance QCD corrections. Rare top decays are therefore not much affected by uncertainties due to QCD corrections. Finally, within the context of the SM and for $m_t \geq 110$ GeV, $B(t \to c\gamma) \approx 10^{-12}$ and $B(t \to cZ) \approx 10^{-12}\text{–}10^{-13}$ [9]. This should be compared with the SM value $B(b \to s\gamma) \approx (2\text{–}4) \times 10^{-4}$. The big difference between the SM values for rare b and t decay rates is due to reasons mentioned above and since $|V_{tb}|^2 \approx 1$ whereas $|V_{cb}|^2 \approx 2 \times 10^{-3}$. The last two features make $t \to c\gamma$ and $t \to cZ$ very good probes for possible new physics. However, background from other SM processes can be a serious problem and detailed studies will have to be carried out before any claims can be made. If one of the two processes is found to occur at a rate much larger than the SM value and if that cannot be explained by SM background processes, then that would be a signature of new physics. Proceeding as in Secs. III and IV we get

$$\delta L_{t\bar{c}\gamma} \approx \frac{g_{\rm ETC}^2 e N_D N_{\rm TC}(1+\bar{Y}/2)m_t}{48\pi^2 M_{\rm ETC}^2} \times [V_{ts}^* \bar{c}_L \sigma^{\mu\nu} t_R F_{\mu\nu} + \text{H.c.}] \tag{14}$$

and

$$\delta^2 L_{t\bar{c}Z} \approx \frac{g_{\rm ETC}^2 g_L N_D N_{\rm TC} m_t (R^2 + \cos 2\theta - \bar{Y}\sin^2\theta)}{96\pi^2 \cos\theta M_{\rm ETC}^2} \times [V_{ts}^* \bar{c}_L \sigma^{\mu\nu} t_R Z_{\mu\nu} + \text{H.c.}] \; . \tag{15}$$

Manifest EM gauge invariance forbids any $d = 4$ operator correction to the $t\bar{c}\gamma$ vertex. However, the $t\bar{c}Z$ vertex is not associated with any gauge invariance and hence $d = 4$ operator correction to it is allowed. Indeed, the leading TC contribution to $t\bar{c}Z$ vertex is of dimension 4 and it arises from the Goldstone pole in the axial-vector channel. This contribution can be evaluated by using chiral Lagrangian [14] or TC spectral functions [12]. For the ordinary TC scenario it is given by

$$\delta L_{t\bar{c}Z} \approx -0.014 g_L [V_{ts}^* \bar{c}_L \gamma^\mu t_L Z_\mu + \text{H.c.}] \; . \tag{16}$$

This result can be obtained from Refs. [12,14] by transforming the quark gauge eigenstates into mass eigenstates and supplying the needed CKM element to effect the transformation. The above expressions for $\delta L_{tc\gamma}$ and δL_{tcZ} can be used to find the TC contributions to the branching ratios of $t \to c\gamma$ and $t \to cZ$. In order to get an estimate of the likely number of rare top decays that could be observed at the LHC we shall assume that the SM charged current two-body decay $t \to bW^+$ gives the dominant decay mode of the top quark also in the TC models considered in this article. This assumption may become somewhat susceptible for TC models with extra chiral symmetries. However, it might still be good enough to give us an order of magnitude estimate of the number of rare top decays that could be expected at the LHC. For an SU($N_{\rm TC}$) model with $N_D = 1, m_t \approx 150$ GeV, and $|V_{ts}| \approx 0.054$ we get

$$B(t \to c\gamma) \approx \frac{\Gamma(t \to c\gamma)}{\Gamma(t \to bW^+)} \approx 8.5 \times 10^{-11} \qquad (17a)$$

and

$$B(t \to cZ) \approx \frac{\Gamma(t \to cZ)}{\Gamma(t \to bW^+)} \approx 2.4 \times 10^{-6} , \qquad (17b)$$

whereas for $N_D = 10$ the branching ratio for $t \to c\gamma$ will be given by 2.34×10^{-9}. The TC contribution to the decay rate of $t \to c\gamma$ is therefore greater than the SM value by 2 orders of magnitude for $N_D = 1$ and by nearly 3 orders of magnitude for $N_D = 10$. On the other hand, the TC contribution to the decay rate of $t \to cZ$ is independent of N_D (except for effects of walking dynamics) and it is greater than the SM value by nearly 7 orders of magnitude. This makes rare top decays and, in particular, $t \to cZ$ an excellent probe for low energy effects of TC dynamics. The decay $t \to cZ$, with the Z subsequently decaying into hadrons, has a huge background from the dominant decay mode $t \to bW^+$. This problem can be avoided by looking for $t \to cZ$ with the Z subsequently decaying into $e^+e^- + \mu^+\mu^-$, which has a branching ratio of 7%. The rare decay $t \to cZ \to c + e^+e^-(c + \mu^+\mu^-)$ will therefore have a branching ratio of 1.7×10^{-7} in ordinary TC. It should be noted that the contribution of any kind of new physics to rare t decays is suppressed relative to rare b decays because the CKM matrix elements $|V_{tb}|$ and $|V_{bc}|$ associated with the dominant decay modes $t \to bW^+$ and $b \to W^-c$ are of order 1 and 0.05, respectively. Since the branching ratios of rare t decays are naturally very small the best places to look for them are high energy and high luminosity colliders, which will produce a large number of top quarks per year. The LHC is expected to produce about 10^8 $t\bar{t}$ pairs per year for $m_t \approx 150$ GeV. Hence, even with 100% detection and tagging efficiencies, rare top decays induced by the SM will be unobservable at the LHC. Among the TC-induced rare t decays, $t \to c\gamma$ will be unobservable at the LHC unless the luminosity can be increased by at least one order of magnitude. On the other hand, TC dynamics can in principle produce a few hundred decays of $t \to cZ$ and a few decays of $[t \to c + e^+e^-(c + \mu^+\mu^-)]$ at the same collider with the projected luminosity. The SSC, which was recently abandoned, would have produced the same number of events. However, a detailed background study needs to be done, but it is beyond the scope of this work. It is also necessary to estimate rare top decay rates in all known extensions of the SM, to compare their predictions, and to determine which scenario gives the largest enhancement over the SM value.

VI. CONCLUSIONS

In this article we have calculated the TC contribution to the rare decays of heavy quarks. We have shown that the light ETC dynamics that gives rise to a large top mass also gives rise to sizable FCNC radiative transitions of heavy quarks. TC contributions to FCNC EM transitions have been calculated by using technivector meson dominance. On the other hand, the Z-mediated FCNC transitions in TC models have been calculated by using TC spectral functions. We find that for $m_t \approx 150$ GeV the TC contribution to the decay rate of $b \to s\gamma$ amounts to an 11% (26%) deviation from the SM value in TC models with 1 and 10 technifermion doublets. We also find that minimal TC ($N_D = 1$) contributions to the branching ratios of $t \to c\gamma$ and $t \to cZ$ are of order 10^{-10} and 10^{-6}, respectively. These values are greater than the corresponding SM contributions by 2 and 7 orders of magnitude. As in the case of the SM, the TC-induced rare decay $t \to c\gamma$ will be unobservable at the LHC unless there is an increase in luminosity. However, although the decay $t \to cZ$ in the context of the SM cannot be seen at the LHC, ordinary TC models are expected to produce a few hundred events at the same collider with the projected luminosity. A 500 GeV or a TeV e^+e^- collider would be the ideal place to study rare t decays, because the t quark events can be cleanly identified. But the production cross section for $t\bar{t}$ pairs at e^+e^- collider ranges between (1 and 0.1) pb for $\sqrt{s} \approx 0.5 - 1$ TeV and $m_t \approx 150$ GeV. Since this cross section is lower than the corresponding value at hadronic colliders by a factor of 10^{-3}–10^{-4}, a very high integrated luminosity would have to be achieved before observation of these rare decays could become feasible.

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