Flavor-changing radiative decay of the t quark

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We study the loop-induced flavor-changing decay $t \to c + \gamma$ in a nonlinear R_ξ gauge for the standard model with three and four generations and in models with two Higgs doublets. We find that this decay could be relevant only in a scenario with two Higgs doublets and the mass of the fourth-generation b' quark larger than m_t. Under these conditions, the flavor-changing electromagnetic decays $T_q^* \to D_q \gamma$ ($T_q \to D_q^* \gamma$) may surpass the standard electromagnetic transition $T_q^* \to T_q \gamma$.

I. INTRODUCTION

Flavor-changing neutral transitions for heavy quarks have received much interest recently. The importance of using rare b decays in probing the standard model (SM) and possible new physics has been pointed out.[1] In particular, it was found that QCD corrections[2] and fourth-generation[3] and supersymmetric[4] effects can be quite significant for the radiative decay $b \to s + \gamma$. It has been suggested also[5] that the b'-quark decay, where b' is the charge $-\frac{1}{3}$ member of the fourth family, could be dominated by loop-induced flavor-changing neutral currents if $m_{b'} < m_t$. In the present paper we report a calculation of the radiative decay $t \to c + \gamma$ in a nonlinear R_ξ gauge for the SM with three and four generations and in models with two Higgs doublets. Our general aim is to find out if this mode could be large enough to be observable and if it is sensitive to new physics beyond SM.

We shall follow the general calculation of the process $q_i \to q_j + \gamma$ for photons of arbitrary momentum and quarks of arbitrary mass made by Deshpande and collaborators[6] in a nonlinear R_ξ gauge. We present in Fig. 1 the Feynman diagrams which contribute to this decay in this gauge. When the masses of the quarks in the external lines are small as compared to M_W, the Glashow-Iliopoulos-Maiani (GIM) mechanism[7] leads to the cancellation of terms of order $G_F e/\pi^2$, resulting[6] in an extra power suppression of the form m_l^2/M_W^2 if the internal quark mass m_l is small, or in a mild logarithmic dependence on m_l for large m_l. Since in the $t \to c + \gamma$ decay one of the masses of the external quarks may be even larger than M_W, it is interesting to know how the GIM mechanism is implemented under this situation.

It is necessary to consider also a possible extra suppression factor arising from the Kobayashi-Maskawa (KM) matrix elements involved in the one-loop contributions. According to the known expectations[8] for three- and four-generation KM matrices, the corresponding suppression factor might be quite important for the $t \to c + \gamma$ decay. Nonetheless, a more accurate estimate of $\Gamma(t \to c\gamma)$ is needed in order to define the chances to observe this decay in collider experiments. In particular, we are interested in finding out if the flavor-changing electromagnetic decays $T_q^* \to D_q \gamma$ ($T_q \to D_q^* \gamma$) where the vector T_q^* and pseudoscalar T_q mesons are composed of a t quark and a light quark q, can compete with the standard electromagnetic transition $T_q^* \to T_q \gamma$. As is well known,[9] weak interactions are expected to dominate the decays of these mesons and the latter electromagnetic transition is strongly suppressed, yielding widths of just a few eV. From this point of view, it is relevant to know if there is a combination of masses, KM matrix elements, or parameters in the two-Higgs-doublet models which could overcome the expected suppression factors in $B(t \to c\gamma)$.

II. THE THIRD- AND FOURTH-GENERATION CONTRIBUTIONS

In the nonlinear R_ξ gauge the vertex $A_\gamma W^\pm \phi^\mp$ vanishes,[6] where $\phi^\pm$ is the unphysical Higgs boson. The

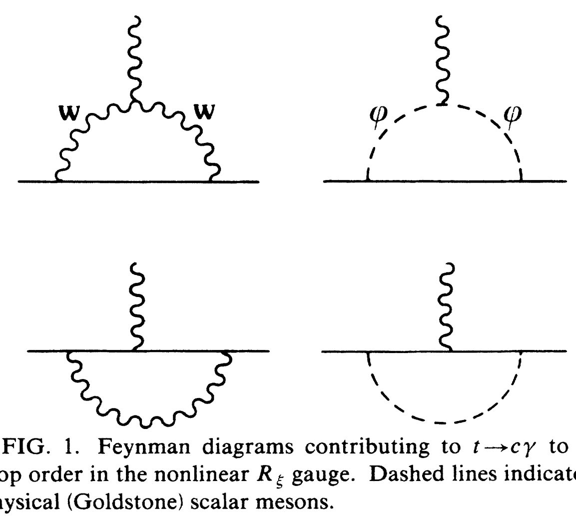

FIG. 1. Feynman diagrams contributing to $t \to c\gamma$ to one-loop order in the nonlinear R_ξ gauge. Dashed lines indicate unphysical (Goldstone) scalar mesons.

number of Feynman diagrams contributing to $\Gamma(t\to c\gamma)$ then reduces to those shown in Fig. 1. Furthermore, since the Ward identities of the electromagnetic vertex in this gauge remain the standard identities in QED, the renormalization of this vertex is rather straightforward.[6] The transition matrix element for the on-shell $t\to c\gamma$ decay is given by

$$M_\mu = i\sigma_{\mu\nu}k^\nu(F_2^L m_c L + F_2^R m_t R)\ , \tag{2.1}$$

where k_ν is the photon momentum, $R,L=(1\pm\gamma_5)/2$, and the magnetic transition form factors $F_2^{L,R}$ are determined by

$$2F_2^L m_c = (B_1+B_3+2B_5)m_t - A_3 m_c + A_{12} - 2A_{11}\ , \tag{2.2}$$

$$2F_2^R m_t = (A_1+A_2+2A_5)m_t - B_3 m_c + B_{12} - 2B_{11}\ . \tag{2.3}$$

The complete expressions for the coefficients A_i and B_i in (2.2) and (2.3) are given in Ref. 6. In particular, if we neglect corrections of order $(m_c/m_W)^2$, the right-hand side of Eq. (2.2) vanishes and the dipole form factor (2.3) reduces to

$$F_2^R = -\frac{ieg^2}{32\pi^2 M_W^2}\int_0^1 dx\int_0^{1-x}dy\left[z_l^2\left[-\frac{x}{X}+\frac{1-x}{Y}\right]+(2+z_l^2)\left[-\frac{x^2+xy-x}{X}+\frac{x^2+xy-x}{3Y}\right]+2\left[-\frac{1-x}{X}+\frac{x}{3Y}\right]\right]\ , \tag{2.4}$$

with

$$X=(1-x)+z_l^2 x+z_t^2 x(x-1)+z_t^2 xy\ , \tag{2.5}$$

$$Y=x+z_l^2(1-x)+z_t^2 x(x-1)+z_t^2 xy\ , \tag{2.6}$$

$z_i=m_i/M_W$, $i=l,t$, and m_l is the mass of the quark in the internal lines of the loops shown in Fig. 1. The integrals in (2.4) can be expressed in terms of Spence functions. However, we have found it much easier to integrate analytically only the variable y and then perform numerically the remaining one-dimensional integrations. In this way the decay rate reduces to

$$\Gamma(t\to c\gamma)=\frac{\alpha G_F^2 m_t^5}{32\pi^4}\sum_l V_l I(z_t,z_l)\ , \tag{2.7}$$

where the function $I(z_t,z_l)$ is given in the Appendix, $V_l=V_{tl}V_{lj}^*$ and V_{ij} are the KM matrix elements.

To calculate the branching ratio we use the relation[2]

$$B(t\to c\gamma)=\frac{\Gamma(t\to c\gamma)}{\Gamma(t\to be\bar{\nu}_e)}B(t\to be\bar{\nu}_e) \tag{2.8}$$

with

$$\Gamma(t\to be\bar{\nu}_e)=\frac{G_F^2 m_t^5}{192\pi^3}\ , \tag{2.9}$$

and an estimated[10] $B(t\to be\bar{\nu}_l)\approx 10\%$.

In Figs. 2 and 3 we present our results for $B(t\to c\gamma)$, with the KM coefficients factored out, for the third and fourth generations as a function of m_t and $m_{b'}$. Using the known constraints of the KM matrix elements for three generations,[11] we can see that $B(t\to c\gamma)$ is of the order 10^{-8}–10^{-6} for $50\le m_t\le 100$ GeV. This result is somewhat low as compared to $B(b\to s\gamma)\sim 10^{-5}$ obtained[2] without QCD corrections. As was the case in the latter decay, we also expect that QCD corrections will enhance the third-generation results shown in Fig. 2 by one order of magnitude.[2] On the other hand, since the b' quark in the fourth-generation contribution to the loop is heavy, QCD corrections are expected to be small in this case.[12] Under these conditions, we can see from Fig. 3 that the fourth-generation contribution may become significant only if the KM matrix elements $V_{tb'}V_{b'c}^*$ are not highly suppressed. According to the conjecture[8] $V_{tb'}\sim\theta^3$, $V_{b'c}\sim\theta^4$, with $\theta\approx 0.2$ the Cabibbo angle, we obtain $10^{-9}\le B(t\to c\gamma)\le 10^{-5}$ for $M_W/2\le m_t, m_b'\le 2M_W$. This result does not represent a clear improvement over the third-generation result and we conclude therefore that the prospects of measuring $B(t\to c\gamma)$ in the SM with three and four generations are not very bright.

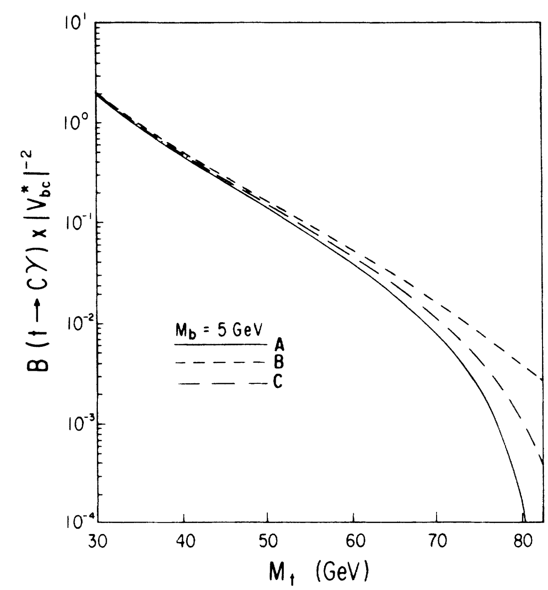

FIG. 2. Results for $B(t\to c\gamma)\times|V_{bc}^*|^{-2}$ as a function of m_t. Curve A: standard model with three quark generations; curves B,C include the contribution of charged Higgs bosons $H^\pm$ according to model II taking $m_H=m_W$ and $m_H=1.5m_W$, respectively. The contribution for model I cannot be distinguished from the SM result, curve A.

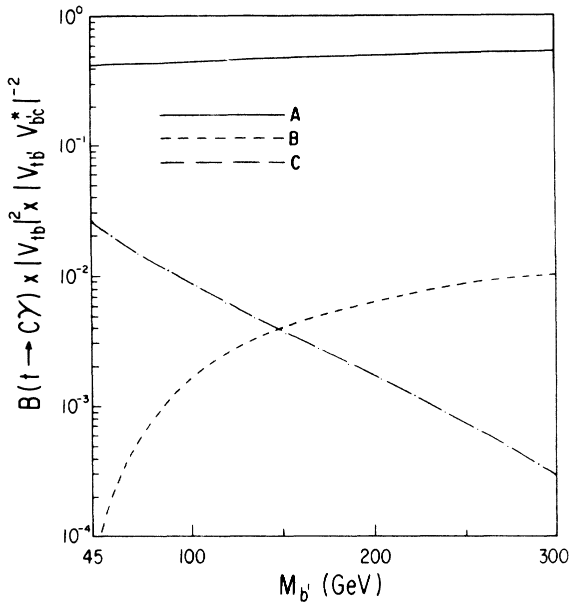

FIG. 3. Results for $B(t \to c\gamma) \times |V_{tb}|^2 \times |V_{tb'}V^*_{b'c}|^{-2}$ as a function of $m_{b'}$ in the frame of the standard model with a fourth generation, taking $m_t = 0.5 m_W$ (curve A), $m_t = m_W$ (curve B), $m_t = 1.5 m_W$ (curve C).

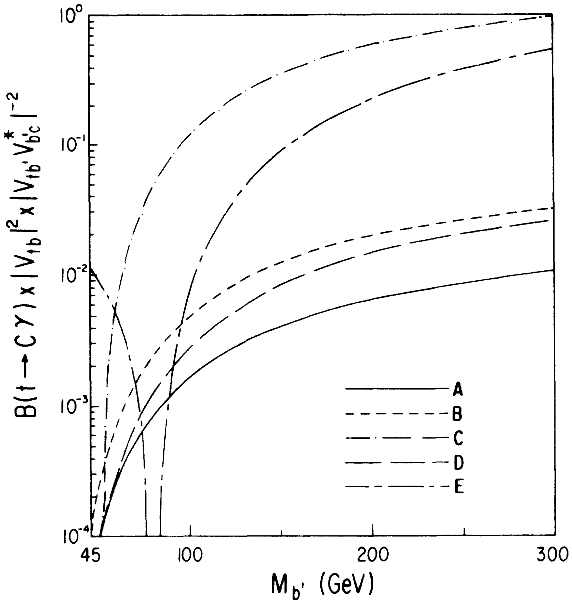

FIG. 4. Results for $B(t \to c\gamma) \times |V_{tb}|^2 \times |V_{tb'}V^*_{b'c}|^{-2}$ as a function of $m_{b'}$, with $m_t = m_W$. Standard model with a fourth generation (curve A). Curves B–E include the contribution of charged Higgs bosons in the frame of (i) model I with $m_H = m_W$ (curve B), $m_H = 1.5 m_W$ (curve D) and (ii) model II with $m_H = m_W$ (curve C), $m_H = 1.5 m_W$ (curve E).

III. CHARGED-HIGGS-BOSON EFFECTS

There are two basic types of two-Higgs-doublet models which avoid tree-level flavor-changing neutral currents.[13] The charged-Higgs-boson Yukawa couplings in these models are given by[13]

$$\frac{g}{\sqrt{2}M_W} H^+ \bar{u}_i V_{ij} (\xi m_i^u L - \xi' m_j^d R) d_j + \text{H.c.} , \quad (3.1)$$

where u_i and d_j represent $+\frac{2}{3}$ and $-\frac{1}{3}$ charge members of each quark generation, respectively, and $\xi = v_2/v_1$, with v_i the vacuum expectation values of the two Higgs doublets. The parameter ξ' in Eq. (3.1) takes the values $\xi' = -\xi^{-1}, \xi$ for models I and II, respectively. The $H^\pm$ contribution to the $t \to c\gamma$ decay induces only one extra diagram to those shown in Fig. 1, with the unphysical scalar $\phi^\pm$ replaced by $H^\pm$ and the corresponding changes for the Yukawa couplings. This is due to the fact that the vertex $H^\pm W^\mp \gamma$ is zero at the tree level in all two-Higgs-doublet models.[14] The $H^\pm$ contributions to $t \to c\gamma$ can be thus expressed in the form (2.7), with just a new $I_H(z_t', z_l')$ function, which is also given in the Appendix, and $z_i' = m_i/m_H$. We shall take for definiteness $\xi = 5$, which is consistent with the restrictions imposed[15] by the B_d^0-$\bar{B}_d^0$ mixing.

In Figs. 1 and 4 we plot $B(t \to c\gamma)$ for the third and fourth generations, respectively, including the $H^\pm$ contribution for two values of m_H and for models I and II. As expected, there is an enhancement of about 3 orders of magnitude of $B(t \to c\gamma)$ in model II. This enhancement is similar to that observed in other processes for this model.[4,5,16]

IV. DECAYS OF t-FLAVORED MESONS

The above scenario could have important consequences for the decays of the t-flavored mesons T_q and T_q^*. In particular, since the weak interactions are expected to dominate their decays and the standard electromagnetic transition $T_q^* \to T_q \gamma$ is expected to be strongly suppressed,[9] the flavor-changing electromagnetic decays $T_q^* \to D_q \gamma$ and $T_q \to D_q^* \gamma$ might play a mayor role. Using the estimate for the hadronic matrix element introduced by Deshpande *et al.*[2] in the context of the constituent quark model, we obtain the scaling factor $\Gamma(T_q^* \to D_q \gamma)$ $[\Gamma(T_q \to D_q^* \gamma)] \sim 10^{-1} \Gamma(t \to c\gamma)$. According to our results shown in Fig. 4, this means that $\Gamma(T_q^* \to D_q \gamma)$ $[\Gamma(T_q \to D_q^* \gamma)]$ could be as large as 10^3 eV in model II with two Higgs doublets and $m_{b'} > m_t$. We have therefore that these flavor-changing electromagnetic decays may compete favorably with the standard electromagnetic decay $T_q^* \to T_q \gamma$, which is expected[9] to have a decay width of the order of few eV.

In conclusion, we have found that the most favorable scenario to observe the flavor-changing transition $t \to c\gamma$ is given in a fourth-generation scheme with two Higgs doublets and $m_{b'} > m_t$. Under these conditions we expect a large enhancement of the decay width $\Gamma(t \to c\gamma)$ and the flavor-changing decays $T_q^* \to D_q \gamma$ $(T_q \to D_q^* \gamma)$ thus surpass the standard transition $T_q^* \to T_q \gamma$. Even more, we expect the former decays to be more easily identifiable than the latter because of the simpler (bottomless) jet structure associated with the hadronic decays of the charmed mesons D_q and D_q^*. Since the photon is monochromatic, it is also possible to distinguish between the flavor-changing and flavor-conserving radiative decays by looking for photons with energies above the appropriate thresholds.[17] This in turn could be also used to discriminate between T_q and T_q^* production[18] in e^+e^- annihilation.

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APPENDIX

Using the integrals

$$I_Z^0=\int_0^1 dx\int_0^{1-x}\frac{dy}{Z}$$
$$=\int_0^1 dx\frac{1}{ax}\ln\left[\frac{c+(a+b)x}{c+bx+ax^2}\right],$$
$$I_Z^1=\int_0^1 dx\int_0^{1-x}\frac{x\,dy}{Z}$$
$$=\frac{1}{a}\int_0^1 dx\,\ln\left[\frac{c+(a+b)x}{c+bx+ax^2}\right],$$
$$I_Z^2=\int_0^1 dx\int_0^{1-x}\frac{x^2+xy}{Z}$$
$$=\frac{1}{a}(1-cI_Z^0-bI_Z^1),$$

with $Z=c+bx+ax^2+axy$ and the GIM restriction $\sum_l V_{il}V_{lj}^*=\delta_{ij}$, the $t\to c\gamma$ decay width takes the form (2.7) with the function

$$I(z_t,z_l)=\frac{1}{z_t^2}\left[-\frac{z_l^2}{2}+(2+z_l^2-2z_t^2)I_X^0\right.$$
$$+[z_l^4+(z_t^2-1)(2-z_l^2)]I_X^1$$
$$+\frac{1}{3}\left[\frac{z_l^2}{2}+z_l^2(z_t^2-2-z_l^2)\right]I_Y^0$$
$$\left.+\tfrac{1}{3}[z_l^4+(z_t^2-1)(2-z_l^2)]I_Y^1\right].$$

Similarly, the $H^\pm$ contribution can be expressed as in Eq. (2.7) with a function

$$I_H(z_t',z_l')=\frac{z_l'^2}{z_t'^2}\left[(\xi'^2(z_l'^2-1)+\xi\xi' z_t'^2)I_{X'}^1+\xi'^2 I_{X'}^0\right.$$
$$-\frac{\xi'^2}{2}+\frac{1}{3}\left[[\xi'^2(z_l'^2-1)+\xi\xi' z_t'^2]I_{Y'}^1\right.$$
$$-(\xi'^2 z_l'^2+\xi\xi' z_t'^2)I_{Y'}^0$$
$$\left.\left.+\frac{\xi'^2}{2}\right]\right],$$

where $\xi'=-\xi^{-1},\xi$ for models I and II, respectively, and X' and Y' are obtained from Eqs. (2.5) and (2.6) substituting z_i by z_i'.

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