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Study on the rare radiative decay $B_c \to D_s^* \gamma$ in the standard model and multiscale walking technicolor model

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Applying the perturbative QCD method, we study the decay $B_c \to D_s^* \gamma$ in the standard model (SM) and multiscale walking technicolor model (MWTCM). In the SM, we find that the contribution of weak annihilation is more important than that of the electromagnetic penguin diagram. The presence of pseudo Goldstone bosons in the MWTCM leads to a large enhancement in the rate of $B_c \to D_s^* \gamma$, but this model is in conflict with the branching ratio of $Z \to b\bar{b}(R_b)$ and the CLEO data on the branching ratio $B(b \to s\gamma)$. If top-color is further introduced, the calculated results in the top-color-assisted MWTCM can be suppressed and be in agreement with the CLEO data for a certain range of parameters. [S0556-2821(96)03219-5]

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I. INTRODUCTION

The inclusive rare decay $B \to X_s \gamma$ was studied several years ago [1]. Recently, the physics of the B_c meson has gained intensive attention [2]. The B_c meson is believed to be the next and the final family of B mesons; it provides a unique opportunity to examine various heavy quark fragmentation models, heavy quark spin-flavor symmetry, different quarkonium bound state model and properties of inclusive decay channels. Furthermore, the radiative weak decays of the B_c meson also offer a rich source to measure Cabibbo-Kobayashi-Maskawa (CKM) matrix elements of the standard model (SM). In this paper, we will address B_c radiative decay $B_c \to D_s^* \gamma$.

Different from the decay $B \to X_s \gamma$, which is mainly induced by the flavor-changing $b \to s\gamma$ neutral currents [3], the bound state effects in the decay $B_c \to D_s^* \gamma$ may be rather large. Bound state effects include modifications from weak annihilation, which involve no neutral flavor-changing currents at all. The effects of the weak annihilation mechanism are expected to be rather large due to the large CKM amplitude. We will address this point in detail below.

Unfortunately, the well-known chiral symmetry [4] and heavy quark symmetry [5] can not be applied to this process. Recently, a perturbative QCD (PQCD) analysis of B meson decays seems to give a good prediction [6]. As argued in Ref. [7], the two-body nonleptonic decay of B_c meson can be studied conveniently within the framework of PQCD, as suggested by Brodsky and Lepage [8] and then developed in Ref. [6]. Here, we preview the reliability of a PQCD analysis of B_c radiative decay: in the process $b \to s\gamma$, the s quark obtains large momentum by recoiling, in order to form a bound state with the spectator $\bar{c}$ quark; most of the momentum of the s quark must be transferred to $\bar{c}$ by a hard scattering process. PQCD [6,8] can be used in the calculation for the hard scattering process because the heavy charm usually shares most of the momentum of the final state (i.e., D_s^*). The relevant Feynman diagrams are given in Fig. 1.

Like in $B \to K^* \gamma$, the subprocess $b \to s\gamma$ in $B_c \to D_s^* \gamma$ is usually controlled by the one-loop electromagnetic penguin diagrams [Fig. 1(a)]. It plays an important role in testing loop effects in the SM and in searching for physics beyond

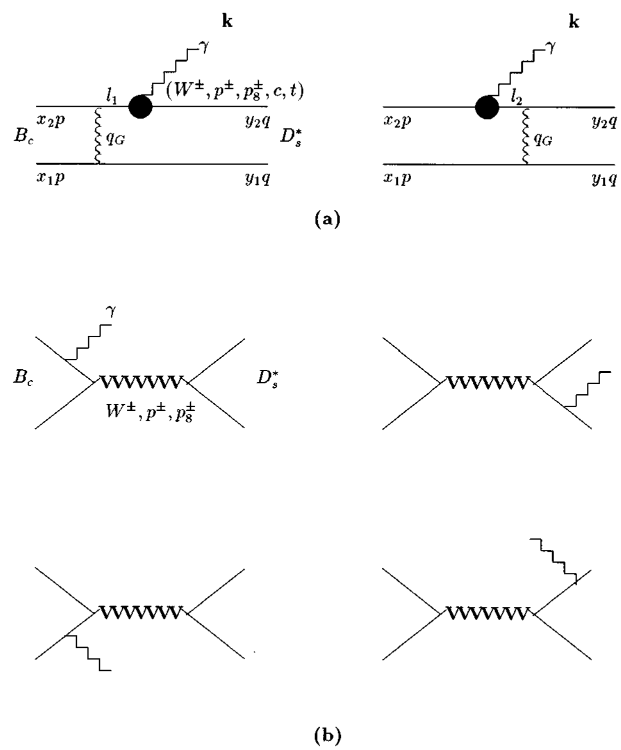

FIG. 1. (a) shows the Feynman diagrams, which contribute to the decay $B_c \to D_s^* \gamma$ through the short distance $b \to s\gamma$ mechanism. The blob represents the electromagnetic penguin operators contributing to $b \to s\gamma$. $x_2 p$ and $x_1 p$ are momenta of b and c quarks in the B_c meson, respectively. $y_2 q$ and $y_1 q$ are momenta of s and c quarks in the D_s^* meson, respectively. (b) represents the Feynman diagrams, which contribute to the decay $B_c \to D_s^* \gamma$ through the weak annihilation.

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the SM (so-called new physics).

Most recently, the contribution of the electromagnetic penguin interaction to the branching ratio $B(b\to s\gamma)$ from pseudo Goldstone bosons (PGB's) in the one-generation technicolor model (OGTM) has been estimated in Ref. [9]. However, we know that there are some problems [such as flavor-changing neutral currents (FCNC's), the large positive contributions to the parameters] in most conventional technicolor (TC) models. Walking technicolor (WTC) has been advocated as a solution to the problem of large flavor-changing neutral current interactions in extended technicolor (ETC) theories of quark and lepton mass generation [10]. Furthermore, the electroweak parameter S in WTC models is smaller than that in the simple QCD-like ETC models and consequently, its deviation from the SM value may fall within current experimental bounds [11]. To explain the large hierarchy of the quark masses, multiscale WTC models (MWTCM's) are further proposed [12].

However, as discussed in Ref. [13], the correction of PGB's in MWTCM's to the $Z\to bb$ branching ratio (R_b) is too large when compared with recent data from the CERN e^+e^- collider LEP. In this paper we calculate the contribution to the branching ratio $B_c\to D_s^*\gamma$ from the PGB's in the MWTCM and find that such a contribution is too large when compared with the CLEO constraint for the inclusive decay $b\to s\gamma$. In general, there are two mechanisms which contribute to the decay $B_c\to D_s^*\gamma$: one proceeds through the short distance $b\to s\gamma$ transition, while the other proceeds through weak annihilation accompanied by photon emission. On the other hand, if top color [14] is further introduced to the multiscale walking technicolor model, the modification from the PGB's in the top color-assisted multiscale walking technicolor model (TAMWTCM) to $B_c\to D_s^*\gamma$ is strongly suppressed, and, therefore, can be consistent with the recent CLEO data for the branching ratio $B(b\to s\gamma)$ [15].

This paper is organized as follows. In Sec. II, we display our calculations in the SM and MWTCM and present the numerical results. Section III contains the discussion.

II. CALCULATION

Using the factorization scheme [8] within PQCD, the momenta of the quarks are taken as some fractions x of the total momentum of the meson weighted by a soft physics distribution functions $\Phi_H(x)$. The peaking approximation is used for $\Phi_H(x)$ [16]; the distribution amplitudes of B_c and D_s^* are

$$\Phi_{B_c}(x)=\frac{f_{B_c}\delta(x-\epsilon_{B_c})}{2\sqrt{3}}, \tag{1a}$$

$$\Phi_{D_s^*}(x)=\frac{f_{D_s^*}\delta(x-\epsilon_{D_s^*})}{2\sqrt{3}}, \tag{1b}$$

where $f_{B_c},f_{D_s^*}$ are decay constants of B_c and D_s^*, respectively, and

$$\epsilon_{B_c}=m_c/m_{B_c}, \tag{1c}$$

$$\epsilon_{D_s^*}=\frac{M_{D_s^*}-m_c}{m_{D_s^*}}. \tag{1d}$$

The spinor parts of the B_c and D_s^* wave functions are

$$\frac{(\not p+m_{B_c})\gamma_5}{\sqrt{2}}, \tag{2a}$$

$$\frac{(\not p-m_{D_s^*})\not\epsilon}{\sqrt{2}}, \tag{2b}$$

where ϵ is the polarization vector of D_s^*.

A. Electromagnetic penguin contribution

The short distance electromagnetic penguin process is governed by the electromagnetic penguin operators [1]. At the weak scale $\mu=m_b$, the effective Hamiltonian for $b\to s\gamma$ transition is

$$H_{\text{eff}}=\frac{4G_F}{\sqrt{2}}V_{tb}V_{ts}^*C_7(m_b)O_7, \tag{3}$$

where

$$O_7=\frac{em_b\overline{s\sigma}_{\mu\nu}F^{\mu\nu}(1+\gamma_5)b}{32\pi^2} \tag{4a}$$

and which is denoted by a blob in Fig. 1(a). The corresponding coefficient of O_7 has the form

$$C_7(m_b)=\rho^{-16/23}\left[C_7(m_W)+\frac{8}{3}(\rho^{2/23}-1)C_8(m_W)\right] + C_2(m_W)\sum_{i=1}^{8}h_i\rho^{-a_i} \tag{4b}$$

with

$$\rho=\alpha_s(m_b)/\alpha_s(m_W),\quad C_2(m_W)=-1, \tag{4c}$$

$$h_i=\left(\frac{626126}{272277},-\frac{56281}{51730},-\frac{3}{7},-\frac{1}{14},-0.6494,-0.0380,-0.0186,-0.0057\right),$$

$$a_i=\left(\frac{14}{23},\frac{16}{23},\frac{6}{23},-\frac{12}{23},0.4086,-0.4230,-0.8994,0.1456\right), \tag{4d}$$

and $C_7(m_W)=\frac{1}{2}A(x)$ and $C_8(m_W)=\frac{1}{2}C(x)$ in the standard model with $x=(m_t/m_W)^2$. The functions $A(x)$ and $C(x)$ arise from graphs with W boson exchange.

In MWTCM, the relevant Feynman rules are the same as Ref. [17]:

$$[p^+ - u_i - d_j] = i \frac{1}{\sqrt{6} F_Q} V_{u_i d_j} [m_{u_i}(1-\gamma_5) - m_{d_j}(1+\gamma_5)], \tag{5a}$$

$$[p_8^+ - u_i - d_j] = i \frac{V_{u_i d_j}}{F_Q} \lambda^a [m_{u_i}(1-\gamma_5) - m_{d_j}(1+\gamma_5)], \tag{5b}$$

where $u=(u,c,t)$, $d=(d,s,b)$, and $V_{u_i d_j}$ is the element of CKM matrix, and finally F_Q is the decay constant of technipions composed of Q in MWTCM.

By explicit calculations, one can get [9]

$$C_7(m_W) = \frac{1}{2} A(x) + \frac{1}{3\sqrt{2} G_F F_Q^2} [B(y) + 8B(z)], \tag{5c}$$

$$C_8(m_W) = \frac{1}{2} C(x) + \frac{1}{3\sqrt{2} G_F F_Q^2} \{D(y) + [8D(z) + E(Z)]\}, \tag{5d}$$

where $y=(m_t/m_{p^\pm})^2$ and $z=(m_t/m_{\frac{\pm}{8}})^2$. The functions B, D, and E arise from diagrams with color singlet and color octet charged PGB's of MWTCM, and the explicit expressions for relevant functions are

$$A(x) = -\frac{x}{12(1-x)^4} [(1-x)(8x^2+5x-7) + 6x(3x-2)\ln x], \tag{6a}$$

$$B(x) = \frac{x}{72(1-x)^4} [(1-x)(22x^2-53x+25) + 6(3x^2-8x+4)\ln x], \tag{6b}$$

$$C(x) = -\frac{x}{4(1-x)^4} [(1-x)(x^2-5x-2) - 6x \ln x], \tag{6c}$$

$$D(x) = \frac{x}{24(1-x)^4} [(1-x)(5x^2-19x+20) - 6(x-2)\ln x], \tag{6d}$$

$$E(x) = -\frac{x}{8(1-x)^4} [(1-x)(12x^2-15x-5) + 18x(x-2)\ln x]. \tag{6e}$$

Now we write down the amplitude of Fig. 1(a) as

$$M_a = \int_0^1 dx_1 dy_1 \Phi_{D_s^*}(y_1) \Phi_{B_c}(x_1) \frac{-iG_F}{\sqrt{2}} V_{tb} V_{ts}^* C_7(m_b) m_b e \frac{\alpha_s(m_b)}{2\pi} C_F \Big\{ T_r[(\not{q} - m_{D_s}^*) \not{\epsilon} \sigma_{\mu\nu}(1+\gamma_5) k^\nu \eta^\mu (\not{p} - y_1 \not{q} + m_b) \times \gamma_\alpha (\not{p} + m_{B_c}) \gamma_5 \gamma^\alpha] \frac{1}{D_1 D_3} + Tr[(\not{q} - m_{D_s^*}) \not{\epsilon} \gamma_\alpha (\not{q} - x_1 \not{p}) \sigma_{\mu\nu}(1+\gamma_5) k^\nu \eta^\mu (\not{p} + m_{B_c}) \gamma_5 \gamma^\alpha] \frac{1}{D_2 D_3} \Big\}, \tag{7}$$

where η is the polarization vector of photon, x_1, y_1 are the momentum fractions shared by charms in B_c and D_s^*, respectively. The functions D_1, D_2, and D_3 in Eq. (7) are the forms of

$$D_1 = (1-y_1)(m_{B_c}^2 - m_{D_s^*}^2 y_1) - m_b^2, \tag{8a}$$

$$D_2 = (1-x_1)(m_{D_s^*}^2 - m_{B_c}^2 x_1), \tag{8b}$$

$$D_3 = (x_1 - y_1)(x_1 m_{B_c}^2 - y_1 m_{D_s^*}^2). \tag{8c}$$

After explicit calculations, the amplitude M_a can be written as

$$M_a = i\varepsilon_{\mu\nu\alpha\beta} \eta_\mu k^\nu \epsilon^\alpha p^\beta f_1^{\rm peng} + \eta^\mu [\epsilon_\mu (m_{B_c}^2 - m_{D_s}^2) - (p+q)_\mu (\epsilon \cdot k)] f_2^{\rm peng} \tag{9}$$

with form factors

$$f_1^{\rm peng} = 2 f_2^{\rm peng} = C \int_0^1 dx_1 dy_1 \delta(x_1 - \epsilon_{B_c}) \delta(y_1 - \epsilon_{D_s^*}) \times \Big\{ [m_{B_c}(1-y_1)(m_{B_c} - 2m_{D_s^*}) - m_b(2m_{B_c} - m_{D_s^*})] \times \frac{1}{D_1 D_3} - m_{B_c} m_{D_s^*}(1-x_1) \frac{1}{D_2 D_3} \Big\}, \tag{10a}$$

where

$$C = \frac{e m_b f_{B_c} f_{D_s^*} C_7(m_b) C_F \alpha_s(m_b) G_F V_{tb} V_{ts}^*}{12\pi\sqrt{2}}. \tag{10b}$$

B. The weak annihilation contribution

As mentioned in Sec. I, B_c meson is also the unique probe of the weak annihilation mechanism.

In the SM, using the formalism developed by Cheng *et al.* [18], the amplitude of annihilation diagrams [see Fig. 1(b)] is

TABLE I. Form factors in the SM, MWTCM, and TAMWTCM, $f^{\rm peng}$ and $f^{\rm anni}$ represent form factors for electromagnetic penguin and weak annihilation process, respectively.

f_i	SM	MWTCM	TAMWTCM
$f_1^{\rm peng}$	-3.05×10^{-10}	$(0.50–1.13)\times10^{-8}$	$(0.44–1.67)\times10^{-9}$
$f_2^{\rm peng}$	-1.57×10^{-10}	$(2.50–5.65)\times10^{-9}$	$(2.22–8.38)\times10^{-10}$
$f_1^{\rm anni}$	7.10×10^{-10}	$(6.75–7.02)\times10^{-10}$	$(6.75–7.02)\times10^{-10}$
$f_2^{\rm anni}$	-1.70×10^{-10}	$(-1.66$ to $-1.53)\times10^{-10}$	$(-1.66$ to $-1.53)\times10^{-10}$

$$M_b^{(W)}=i\varepsilon_{\mu\nu\alpha\beta}\eta^\mu k^\nu \epsilon^\alpha p^\beta f^{\rm anni}_{1(W)}+\eta^\mu[\epsilon_\mu(m^2_{B_c}-m^2_{D_s})-(p+q)_\mu(\epsilon\cdot k)]f^{\rm anni}_{2(W)} \tag{11}$$

with

$$f^{\rm anni}_{1(W)}=2\zeta\left[\left(\frac{e_s}{m_s}+\frac{e_c}{m_c}\right)\frac{m_{D_s^*}}{m_{B_c}}+\left(\frac{e_c}{m_c}+\frac{e_b}{m_b}\right)\right]\frac{m_{D_s^*}m_{B_c}}{m^2_{B_c}-m^2_{D_s^*}}, \tag{12a}$$

$$f^{\rm anni}_{2(W)}=-\zeta\left[\left(\frac{e_s}{m_s}-\frac{e_c}{m_c}\right)\frac{m_{D_s^*}}{m_{B_c}}+\left(\frac{e_c}{m_c}-\frac{e_b}{m_b}\right)\right]\frac{m_{D_s^*}m_{B_c}}{m^2_{B_c}-m^2_{D_s^*}}, \tag{12b}$$

where

$$\zeta=ea_2\frac{G_F}{\sqrt{2}}V_{cb}V^*_{cs}f_{B_c}f_{D_s^*},\quad a_2 \text{ is a parameter.} \tag{12c}$$

In MWTCM, using the Feynman rules in Eq. (5a), Eq. (5b), and the methods in Ref. [18], we can write down the amplitude of charged PGB's annihilation diagrams [see Fig. 1(b)]:

$$M_b^{(p)}=i\varepsilon_{\mu\nu\alpha\beta}\eta^\mu k^\nu \epsilon^\alpha p^\beta f^{\rm anni}_{1(p)}+\eta^\mu[\epsilon_\mu(m^2_{B_c}-m^2_{D_s^*})-(p+q)_\mu(\epsilon\cdot k)]f^{\rm anni}_{2(p)} \tag{13}$$

with

$$f^{\rm anni}_{1(p)}=-\zeta'\left[\left(\frac{e_s}{m_s}+\frac{e_c}{m_c}\right)\frac{m_s-m_c}{m_{B_c}}+\left(\frac{e_b}{m_b}+\frac{e_c}{m_c}\right)\frac{m_b-m_c}{m_{B_c}}\right]\frac{m_{B_c}m_{D_s^*}}{m^2_{B_c}-m^2_{D_s^*}}, \tag{14a}$$

$$f^{\rm anni}_{2(p)}=\frac{1}{2}\zeta'\left[\left(\frac{e_s}{m_s}+\frac{e_c}{m_c}\right)\frac{m_{D_s^*}}{m_{B_c}}+\left(\frac{e_b}{m_b}+\frac{e_c}{m_c}\right)\right]\frac{m_{D_s^*}m_{B_c}}{m^2_{B_c}-m^2_{D_s^*}}, \tag{14b}$$

$$\zeta'=ea_2\left[\frac{2C_F}{m^2_{p_8^\pm}}+\frac{1}{12m^2_{p^\pm}}\right]\frac{V_{cb}V^*_{cs}}{F_Q^2}f_{B_c}f_{D_s^*}(m^2_{B_c}+m^2_{D_s^*}). \tag{14c}$$

The total annihilation amplitude [Fig. 1(b)] in the MWTCM is consequently the form of

$$M_b=M_b^{(W)}+M_b^{(p)}=i\varepsilon_{\mu\nu\alpha\beta}\eta^\mu k^\nu \epsilon^\alpha p^\beta f^{\rm anni}_1+\eta^\mu[\epsilon_\mu(m^2_{B_c}-m^2_{D_s^*})-(p+q)_\mu(\epsilon\cdot k)]f^{\rm anni}_2 \tag{15}$$

with

$$f^{\rm anni}_1=f^{\rm anni}_{1(W)}+f^{\rm anni}_{1(p)}, \tag{16a}$$

$$f^{\rm anni}_2=f^{\rm anni}_{2(W)}+f^{\rm anni}_{2(p)}. \tag{16b}$$

C. Numerical results

We will use the following values for various quantities as input in our calculation.

(i) Decay constants for pseudoscalar B_c and vector meson D_s^*,

$$f_{D_s^*}=f_{D_s}=344\ \text{MeV}$$

from the reports by three groups [19] and

$$f_{B_c}=500\ \text{MeV}$$

from the results in Ref. [20].

(ii) Meson mass and the constituent quark mass,

$$M_{D_s^*}=2.11\ \text{GeV},\quad m_b=4.7\ \text{GeV},$$

$$m_c=1.6\ \text{GeV},\quad m_s=0.51\ \text{GeV}$$

from the Particle Data Group [21], and

$$m_{B_c}=6.27\ \text{GeV}$$

as estimated in Ref. [22]. We also use $m_{B_c}\approx m_b+m_c$, $m_{D_s^*}\approx m_s+m_c$ in our calculation.

(iii) The parameter a_2 appearing in nonleptenic B decays was extracted recently from the CLEO data [23] on $B\to D^*\pi(\rho)$ and $B\to J/\psi K^*$ by Cheng *et al.* [18]. Here, we take

$$a_2=\tfrac{1}{2}(c_- - c_+)=0.21.$$

(iv) For CKM elements [21], we use

$$V_{cb}=0.04,\quad |V_{ts}|=V_{cb},\quad |V_{cs}|=0.9745,\quad V_{tb}=0.9991.$$

(v) The QCD coupling constant $\alpha_s(\mu)$ at any renormalization scale can be calculated from $\alpha_s(m_Z)=0.117$ via

$$\alpha_s(\mu)=\frac{\alpha_s(m_Z)}{1-(11-\frac{2}{3}n_f)[\alpha_s(m_Z)/2\pi]\ln(m_Z/\mu)},$$

TABLE II. The decay rates in the SM, MWTCM, and TAMWTCM. The $\Gamma^{\rm peng}$, $\Gamma^{\rm anni}$, and $\Gamma^{\rm total}$ represent $\Gamma(B_c \to D_s^* \gamma)$ through penguin, annihilation, and penguin+annihilation diagrams, respectively.

$\Gamma(B_c \to D_s^* \gamma)$	SM	MWTCM	TAMWTCM
$\Gamma^{\rm peng}$ (GeV)	3.18×10^{-19}	$(0.86–4.36)\times10^{-16}$	$(0.67–9.55)\times10^{-18}$
$\Gamma^{\rm anni}$ (GeV)	1.06×10^{-18}	$(0.94–1.03)\times10^{-18}$	$(0.94–1.03)\times10^{-18}$
$\Gamma^{\rm total}$ (GeV)	9.92×10^{-19}	$(0.93–4.52)\times10^{-16}$	$(0.23–1.26)\times10^{-17}$

and we obtain

$$\alpha(m_b)=0.203, \quad \alpha_s(m_W)=0.119.$$

(vi) For the masses of $m_{p^\pm}$ and $m_{p_8^\pm}$ in MWTCM, Ref. [12] has presented a constraint on them; here we take

$$m_{p^\pm}=(100-250)\ \text{GeV},$$

$$m_{p_8^\pm}=(300-600)\ \text{GeV}.$$

(vii) The decay constant F_Q satisfies the constraint [12]

$$F_\pi=\sqrt{F_\psi^2+3F_Q^2+N_L F_L^2}=246\ \text{GeV}.$$

It is found in Ref. [12] that $F_Q=F_L=20-40$ GeV. We will take

$$F_Q=40\ \text{GeV}$$

in our calculation.

We give the long- and short-distance contributions to the form factors f_1 and f_2 in the SM and MWTCM in Table I, so do the decay width in Table II using the amplitude formula:

$$\Gamma(B_c \to D_s^* \gamma)=\frac{(m_{B_c}^2-M_{D_s^*}^2)^3}{32\pi m_{B_c}^3}(f_1^2+4f_2^2).$$

The lifetime of B_c was given in Ref. [24]. In this paper we use

$$\tau_{B_c}=(0.4\ \text{ps}–1.35\ \text{ps})$$

to estimate the branching ratio $B(B_c \to D_s^* \gamma)$, which is a function of τ_{B_c}. The results are given in Table III.

III. DISCUSSION

We have studied two kinds of contributions to the process $B_c \to D_s^* \gamma$. For the short-distance one [as illustrated in Fig. 1(a)] induced by electromagnetic penguin diagrams, the momentum square of the hard scattering being exchanged by gluon is 3.6 GeV^2, which is large enough for PQCD analyzing. The hard scattering process cannot be included conveniently in the soft hadronic process described by the wave function of the final bound state. That is one important reason why we cannot apply the commonly used spectator model [25] to the two-body B_c decays. There is no phase space for the propagators appearing in Fig. 1(a) to go on shell, so the imaginary part of M_a is absent, unlike the situation in Ref. [6]. Another competitive mechanism is the weak annihilation. In SM, we find that it is more important than the former one. This situation is different from that of the radiative weak $B^\pm$ decays, which is overwhelmingly dominated by electromagnetic penguin diagram. The results stem from two reasons: one is that the compact size of B_c meson enhances the importance of annihilation decays; the other comes from the Cabibbo allowance. In $B_c \to D_s^* \gamma$ process, the CKM amplitude of weak annihilation is $|V_{cb}V_{cs}^*|$, but in $B_\pm \to K^\pm \gamma$ process the CKM part is $|V_{ub}V_{us}^*|$, which is much smaller than $|V_{cb}V_{cs}|$.

In addition, we find that the contribution from PGB's in MWTCM to the short distance process $b \to s\gamma$ is too large due to the smallness of the decay constant F_Q in this model. In contrast, the contribution from PGB's through the weak annihilation process is negligibly small. In general, the modification from PGB's in MWTCM is too large to be consistent with the recent CLEO data on the branching ratio $B(b \to s\gamma)$.

In view of the above situation, we consider the TAMWTCM. The motivation of introducing top-color to MWTCM is the following: in the original MWTCM, it is very difficult to generate the top quark mass as large as that measured in the Collider Detector at Fermilab (CDF) and D0 experiments [26], even with ''strong'' ETC [27]. Thus, top-color interactions for the third-generation quarks seem to be required at an energy scale of about 1 TeV [28]. In the TAMWTCM, top color is still a walking theory to avoid the large FCNC [14]. As in other top color-assisted technicolor theories, the electroweak symmetry breaking is driven mainly by technicolor interactions, which are strong near 1 TeV. The ETC interactions give contributions to all quark and lepton masses, while the large mass of the top quark is mainly generated by the top-color interactions introduced to the third-generation quarks. From Ref. [28], we can reasonably get the ETC-generated part of the top quark mass $m_t'=66k$ GeV

TABLE III. The branching ratios $(B_c \to D_s^* \gamma)$. The $B_{\rm total}^{\rm SM}$, $B_{\rm total}^{\rm MWTCM}$, and $B_{\rm total}^{\rm TAMWTCM}$ represent the branching ratio $(B_c \to D_s^* \gamma)$ in the SM, MWTCM, and TAMWTCM, respectively.

τ_{B_c}	0.4 ps	1.0 ps	1.35 ps
$B_{\rm total}^{\rm SM}$	6.03×10^{-7}	1.51×10^{-6}	2.04×10^{-6}
$B_{\rm total}^{\rm MWTCM}$	$(0.57–2.75)\times10^{-4}$	$(1.42–6.86)\times10^{-4}$	$(1.91–9.26)\times10^{-4}$
$B_{\rm total}^{\rm TAMWTCM}$	$(1.37–7.66)\times10^{-6}$	$(0.34–1.91)\times10^{-5}$	$(0.46–2.58)\times10^{-5}$

with $k \sim 1$–10^{-1}. To compare with the original MWTCM, we here take $m_t' = 35$ GeV as the input parameter in our calculation. (i.e., in the above calculations, $m_t = 174$ GeV is replaced by $m_t' = 35$ GeV, the other calculations are the same as the original MWTCM). The corresponding results obtained in the framework of TAMWTCM are also listed in Tables I–III. From the results in these tables, we can see that the modifications from PGB's in top color assisted MWTCM to $B_c \to D_s^* \gamma$ are strongly suppressed relative to that in the original MWTCM. The branching ratio $B(B_c \to D_s^* \gamma)$ in TAMWTCM is, therefore, consistent with the recent CLEO constraint on the branching ratio $B(b \to s\gamma)$ for a certain range of the parameters.

In this paper, we neglected the contribution of the vector meson dominance (VMD) [29] due to the smallness of $J/\psi(\psi')$-γ coupling.

Finally, we estimate the possibility of observing the interesting process of $B_c \to D_s^* \gamma$ at Tevatron and at the CERN Large Hadron Collider (LHC). The number of B_c at Tevatron and at LHC have estimated to be [30] 16 000 [for 25 Pb^{-1} integrated luminosities with cuts of $P_T(B_c) > 10$ GeV, $y(B_c) < 1$] and 2.1×10^8 (for 100 fb^{-1} integrated luminosities with cuts of $P_T(B_c) > 20$ GeV, $y(B_c) < 2.5$), respectively. By comparing the above predicted number of B_c events with the branching ratio $B^{\mathrm{SM}}_{\mathrm{total}}(B_c \to D_s^* \gamma)$ as given in Table III, one can understand that although this channel is unobservable at the Fermilab Tevatron, but more than one thousand events of interest will be produced at LHC, so it can be well studied at LHC in the future. Furthermore, it is easy to see that the branching ratio $B(B_c \to D_s^* \gamma)$ in the TAMWTCM is roughly one order higher than that in the SM. Therefore, if one finds a clear surplus of B_c events in LHC experiments than that expected in the SM, one may interpret it as a signal of new physics.

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