A study on the rare radiative decay

\[ B_c \rightarrow D_s^*\gamma \]

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Abstract

We study the decay \( B_c \rightarrow D_s^*\gamma \). There are two mechanisms contributing to the process. One proceeds through the short distance \( b \rightarrow s\gamma \) transition and the other occurs through weak annihilation accompanied by a photon emission. The electromagnetic penguin contribution is estimated by perturbative QCD and found to be 4.68 \( \times 10^{-18}GeV \). In particular, we find the contribution of the weak annihilation is 6.25 \( \times 10^{-18}GeV \) which is in the same order as that of the electromagnetic penguin. The total decay rate \( \Gamma(B_c \rightarrow D_s^*\gamma) \) is predicted to be 1.45 \( \times 10^{-17}GeV \) and the branching ratio \( Br(B_c \rightarrow D_s^*\gamma) \) is predicted to be 2.98 \( \times 10^{-5} \) for \( \tau_{B_c} = 1.35ps \). The decays \( B_c \rightarrow D_s^*\gamma \) can be well studied at LHC in the near future.

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1 INTRODUCTION

Recently the physics of $B_c$ meson have got intensive attention[1,2]. It is believed that the $B_c$ meson is the next and the final member of B mesons. It provides unique opportunities to examine various heavy quark fragmentation models, heavy quark spin-flavor symmetry, different quarkonium bound state models, and properties of inclusive and exclusive decay channels. Being made of two heavy quarks of different flavors, $B_c$ weak decays also offer a rich source to measure CKM matrix elements of the standard model[2].

The interest of studying rare B decays lies on the fact that those decays, induced by the flavor-changing neutral currents $b \to s\gamma$, are usually controlled by the one-loop electromagnetic penguin diagrams. They play important role in testing loop effects in the standard model and in searching for physics beyond the standard model( so called new physics). Most recently the weak radiative decays of B meson and bottom baryons have been systematically studied in [3,4]. Hence a precise study of the weak radiative decays of $B_c$ needs to be undertaken. In the present paper, we will address $B_c$ radiative decay $B_c \to D^*_s\gamma$.

The subprocess $b \to s\gamma$, taken as a free decay, is usually treated as the only flavor-changing contribution to $B \to K^{*}\gamma$[5], so do $B_c \to D^*_s\gamma$. However, boundstate effects could seriously modify the results of the assumption. Boundstate effects include modifications from weak annihilation which involves no flavor-changing neutral currents at all. The effects of the weak annihilation mechanism are expected to be large in the $B_c$ decays. We shall address this point in detail below.

Unfortunately, the well known chiral-symmetry[6] and the heavy quark symmetry[7] cannot be applied straightforwardly to this process. Recently, a perturbative QCD(PQCD) analysis of the B meson exclusive decays seems give a good prediction[8]. As it is argued in ref.[9] that $B_c$ two body nonleptonic decay can be conveniently studied within the framework of PQCD suggested by Brodsky-Lepage[10] and then developed in [8].

The reason to use PQCD to analyze of $B_c$ radiative decay is as bellow. In the process $b \to s\gamma$, s quark obtains large momentum by recoiling. In order to form a bound state with the spectator $\tau$ quark, the most part of the momentum of s quark must be transferred to $\tau$ by a hard scattering process, since in the final bound state (i.e. $D^*_{s\tau}$) the heavy charm
should share the most part of the momentum of $D^*_s$. The hard scattering is suitable for PQCD calculation\cite{8,10}. The Feynmen diagrams we will calculate are given in Fig.1.

The paper is organized as follows. In section 2, we display our calculations. We present our numerical results in Section 3. Section 4 contains the discussions and conclusions.

## 2 CALCULATION

Exclusive processes at large momentum transfer are exploited by Brodsky-Lepage\cite{10} within PQCD starting with a Fock component expansion of the involved hadrons, where a twist expansion suggests that the contribution from the lowest order Fock component dominates the physical observable under consideration. An exclusive process then involves a perturbatively calculable hard scattering amplitude convoluted with nonperturbative soft physics wavefunctions of the initial and the final hadrons. Although these wavefunctions are incalculable from first principle, but they are universal for each hadron, \textit{i.e} they are factorized out from the hard scattering amplitude and hence are independent of the process involved.

The factorization scheme advocated by Brodsky-Lepage\cite{10} is employed, where the momenta of quarks are taken as some fractions $x$ of the total momentum of the meson weighted by a soft physical distribution functions $\phi_H(x)$. The peaking approximation is used for $\phi_H(x)$\cite{11}. The distribution amplitude of $B_c$ is

$$
\phi_{B_c}(x) = \frac{1}{2\sqrt{3}} f_{B_c} \delta(x - \frac{m_i}{m_1 + m_2}),
$$

where $m_i$ corresponds to $m_b, m_c$. The distribution function of $D^*_s$ is

$$
\phi_{D^*_s}(x) = \frac{1}{2\sqrt{3}} f_{D^*_s} \delta(x - \frac{\epsilon_{D^*_s}}{m_{D^*_s}}),
$$

where $\epsilon_{D^*_s}$ is defined as

$$
\epsilon_{D^*_s} = \frac{m_{D^*_s} - m_c}{m_{D^*_s}}.
$$

The spinor part of $B_c$ and the final meson $D^*_s$ are

$$
\begin{align*}
\frac{1}{\sqrt{2}} \gamma_5, & \quad \frac{1}{\sqrt{2}} \epsilon,
\end{align*}
$$

where $\epsilon$ is the polarization vector of $D^*_s$. 

3
2.1 Electromagnetic penguin contributions

Within the standard model, the process is governed by the electromagnetic penguin operators\cite{12}, for \( m_s \ll m_b \)

\[ H_{\text{eff}} = \frac{4G_F}{\sqrt{2}}V_{tb}V_{ts}^* C_7(\mu)O_7. \]  

(5)

Here we use the notation of Grinstein et al.\cite{12}

\[ O_7 = \frac{e}{16\pi^2}m_b\sigma^{\mu\nu}F_{\mu\nu} \frac{1 + \gamma_5}{2}b, \]  

(6)

we denote it by a blob in Fig.1.a.

The coefficient of \( O_7 \) includes the QCD corrections, at the scale of W mass,

\[ C_7(m_W) = \frac{x}{24(1-x)}[(1-x)(8x^2 + 5x - 7) + 6x(3x - 2)\ln x], \]  

(7)

where \( x = m_t^2/m_W^2 \), and when it runs down to low scale \( \mu = m_b \), it turns out to be

\[ C_7(m_b) = y^\frac{16}{4}\left[C_7(m_W) - \frac{58}{135}(y_{11} - 1) - \frac{29}{189}(y_{22} - 1)\right] \]  

(8)

where \( y = \alpha_s(m_b)/\alpha_s(m_W) \). We have neglected the mixing of \( O_7 \) with other operators which give small effects.

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

\[ M_a = \int_0^1 dx_1 dy_1 \phi_{D_2}^* \phi_{B_c} C \]

\[ \times \left\{ \frac{1}{4} \left( \text{Tr} \left[ \left( \frac{\phi^*}{m_{D_2}} \sigma_{\mu\nu} (1 + \gamma_5) k^\mu \eta^\nu (\phi - y_1 \frac{\phi}{m_{D_2}}) + m_b \gamma_\alpha (\phi + m_{B_c}) \gamma_5 \gamma^\alpha \right) \right] \frac{1}{D_1} \frac{1}{D_3} \right\} \]  

\[ + \text{Tr} \left[ \left( \frac{\phi^*}{m_{D_2}} \sigma_{\mu\nu} (1 + \gamma_5) k^\mu \eta^\nu (\phi - y_1 \frac{\phi}{m_{D_2}}) + m_b \gamma_\alpha (\phi + m_{B_c}) \gamma_5 \gamma^\alpha \right) \right] \frac{1}{D_2} \frac{1}{D_3} \right\} \]  

(9)

where \( \eta \) is the polarization vector of the photon, \( x_1 \) and \( y_1 \) are the momentum fractions shared by charm in \( B_c \) and \( D_{2s} \), respectively. The denominators are

\[ D_1 = (1 - y_1)(m_{B_c}^2 - m_{D_2}^2 y_1) - m_b^2 \]

\[ D_2 = (1 - x_1)(m_{D_2}^2 - m_{B_c}^2 x_1) \]  

(10)

\[ D_3 = (x_1 - y_1)(x_1 m_{B_c}^2 - m_{D_2}^2 y_1), \]

and the constant \( C \) is

\[ C = \frac{f_{B_c} f_{D_2^s}}{2\sqrt{3} 2\sqrt{3}} C_F C_7(m_b) e \frac{\alpha_s}{4\pi} 4G_F \sqrt{2} V_{tb} V_{ts}^* V_{s_b} m_b. \]  

(11)
The result is

\[ M_a = i \varepsilon_{\mu\nu\alpha\beta} \epsilon^{*\mu} \eta^{*\nu} q^\alpha k^\beta f_{1^{\text{peng}}}^{\gamma}((\eta^* \cdot \epsilon^* p \cdot k - \eta^* \cdot p^* \epsilon^* k)) f_{2^{\text{peng}}}^{\gamma}, \]  
\] (12)

with the formfactors

\[ f_{1^{\text{peng}}}^{\gamma} = f_{2^{\text{peng}}}^{\gamma} = C \int_0^1 dx_1 dy_1 \delta(x_1 - \frac{m_c}{M_{B_c}^*})\delta(y_1 - 1 + \epsilon_D) 4 \times \left\{ \left[ m_{B_c}(1 - y_1)(m_{B_c} - 2m_D^*) - m_b(2m_{B_c} - m_D^*) \right] \frac{1}{D_1 D_3} - m_{B_c} m_{D_c^*} (1 - x_1) \frac{1}{D_2 D_3} \right\}, \] (13)

2.2 The weak annihilation contribution

Being made up of two different heavy flavors, \( B_c \) meson is also the unique place to probe the weak annihilation mechanism.

Using the formalism developed by Cheng et al. [3], the amplitude of Fig.1.b are found to be

\[ M_b = i \varepsilon_{\mu\nu\alpha\beta} \epsilon^{*\mu} \eta^{*\nu} q^\alpha k^\beta f_{1^{\text{anni}}}^\gamma((\eta^* \cdot \epsilon^* p \cdot k - \eta^* \cdot p^* \epsilon^* k)) f_{2^{\text{anni}}}^\gamma, \]  
\] (14)

with

\[ f_{1^{\text{anni}}}^\gamma = \kappa a_1 \left[ \left( \frac{e_s}{m_s} + \frac{e_c}{m_c} \right) \frac{m_{D^*_c}}{m_{B_c}} + \left( \frac{e_c}{m_c} + \frac{e_b}{m_b} \right) \frac{m_{D^*_c} m_{B_c}}{m_{B_c}^2 - m_{D^*_c}^2} \right], \]

\[ f_{2^{\text{anni}}}^\gamma = -\kappa a_1 \left[ \left( \frac{e_s}{m_s} - \frac{e_c}{m_c} \right) \frac{m_{D^*_c}}{m_{B_c}} + \left( \frac{e_c}{m_c} - \frac{e_b}{m_b} \right) \frac{m_{D^*_c} m_{B_c}}{m_{B_c}^2 - m_{D^*_c}^2} \right], \]  
\] (15)

where \( \kappa = e G_V V_{cb} V_{cs} f_{B_c} f_{D^*_c}/\sqrt{2} \), and the \( m_i \) is the constituent quark mass. The parameter \( a_1 \), corresponding to the one appearing in the nonleptonic \( B \) decays, includes the QCD corrections to the four Fermion operator \((\bar{\tau}b)_{V-A}(\bar{\tau}c)_{V-A}\). In nonleptonic \( B \) decays, \( a_1 \) is extracted to be 1.01 from the recent CLEO data on \( B \rightarrow D^{(*)} \pi (\rho) \) and \( B \rightarrow J/\psi K^* \) by Cheng et al. [3]. Here we take \( a_1 = 1 \) in our numerical results.

3 Numerical Results

In order to have a numerical estimation we adopt the following parameters.

1. Meson mass and the constituent quark mass

\[ M_{D^*_c} = 2.11 \text{GeV}, \quad m_b = 4.7 \text{GeV}, \quad m_c = 1.6 \text{GeV}, \quad m_s = 510 \text{MeV} \]
from the Particle Data Group[14]. $m_{B_c}$ has been estimated to be about 6.27GeV in the literature[15,16].

2. Decay constants of $B_c$ and $D_s^*$. Up to now $f_{D_s}$ has been reported by three groups

$$f_{D_s} = \begin{cases} 
232 \pm 45 \pm 20 \pm 48 \text{MeV}, & \text{S.Aoki et al.,[17]} \\
344 \pm 37 \pm 52 \pm 42 \text{MeV}, & \text{D.Acosta et al.,[18]} \\
430 \pm 1^{150} \pm 40 \pm 48 \text{MeV}, & \text{J.Z.Bai et al.,[19]} 
\end{cases}$$

where the first errors are statistical, the second are systematic, and the third are uncertainties involved in extracting the branching fraction $B(D_s^+ \to \mu^+\nu_\mu)$. From heavy quark symmetry, $f_{D_s^*} = f_{D_s}$, we take $f_{D_s^*} = 344\text{MeV}$.

The pseudoscalar decay constant $f_{B_c}$ is related to the ground-state $c\bar{b}$ wave function at the origin by the Van Royen-Weisskopf formula[20] modified for color

$$f_{B_c} = \sqrt{\frac{12}{m_{B_c}}} |\Psi_{100}(0)| = \sqrt{\frac{3}{\pi m_{B_c}}} | R_{10}(0)|.$$ 

Using the recent results for $R_{10}$ derived by Quigg[21]. One finds

$$f_{B_c} = \begin{cases} 
500\text{MeV}, & \text{(Buchmuller - Tye potential[22])} \\
510\text{MeV}, & \text{(Power - law potential[23])} \\
479\text{MeV}, & \text{(Logarithmic potential[24])} \\
696\text{MeV}, & \text{(Cornell potential[25])} 
\end{cases}$$

we will take $f_{B_c} = 500\text{MeV}$ for our numerical estimation.

3. CKM elements and strong couplings

For quark mixing matrix elements, we will take $V_{cb} = 0.040[26], |V_{ts}| = V_{cb}, |V_{cs}| = 0.9745$ and $V_{tb} = 1$. The QCD coupling constant $\alpha_s(\mu)$ at any renormalization scale $\mu$ can be calculated from $\alpha_s(M_Z) = 0.117$ via

$$\alpha_s(\mu) = \frac{\alpha_s(M_Z)}{1 - \frac{\alpha_s(M_Z)}{2\pi} \beta_0 \frac{\alpha_s(M_Z)}{2\pi} \ln(M_Z/\mu)}$$

and $\beta = 11 - \frac{2}{3}n_f$. We get $C_7(M_W) = -0.1953$ and $C_7(M_b) = -0.2939$.

With the values given above for various quantities, we can compute the form factors, the decay rates and the branching ratios.

The results are

$$f_1^{open} = 1.170 \times 10^{-9}, f_2^{open} = f_1^{open}, f_1^{anni} = 1.724 \times 10^{-9}, f_2^{anni} = -8.273 \times 10^{-10},$$
in units of \( GeV^{-1} \).

The ratios of long- and short distance contributions to the form factors \( f_1 \) and \( f_2 \) are

\[
\frac{f_{anni}}{f_{1\, penguin}} = 1.474, \quad \frac{f_{anni}}{f_{2\, penguin}} = -0.710,
\]

which are independent of \( f_{B_c} \) and \( f_{D_s^*} \), and expected to be reliable.

From the amplitude formula eq(12)(14), we get

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

We finally obtain the decay rates

\[
\Gamma(B_c \to D_s^*\gamma) = \begin{cases} 
4.68 \times 10^{-18} \text{GeV}, & \text{only penguin} \\
6.25 \times 10^{-18} \text{GeV}, & \text{only annihilation} \\
1.45 \times 10^{-17} \text{GeV}, & \text{penguin + annihilation}
\end{cases}
\]

The lifetime of \( B_c \) is predicted to be \( \tau_{B_c} = 1.35 \pm 0.15 \text{ps} \) by Quigg[27], but other theorists report smaller values \( \tau_{B_c} = 0.4 \text{ps} \sim 0.7 \text{ps}[27] \). We estimate the branching ratio \( Br(B_c \to D_s^*\gamma) \) as a function of \( \tau_{B_c} \). The results are given in Table 1.

| \( Br(B_c \to D_s^*\gamma) \) | \( \tau_{B_c} = 0.4 \text{ps} \) | \( \tau_{B_c} = 0.7 \text{ps} \) | \( \tau_{B_c} = 1.0 \text{ps} \) | \( \tau_{B_c} = 1.35 \text{ps} \) |
|-----------------|-----------------|-----------------|-----------------|-----------------|
| \( Br^{penguin} \) | 2.85 \times 10^{-6} | 4.98 \times 10^{-6} | 7.11 \times 10^{-6} | 9.61 \times 10^{-6} |
| \( Br^{anni} \) | 3.80 \times 10^{-6} | 6.65 \times 10^{-6} | 9.50 \times 10^{-6} | 1.28 \times 10^{-5} |
| \( Br^{total} \) | 8.83 \times 10^{-6} | 1.54 \times 10^{-5} | 2.21 \times 10^{-5} | 2.98 \times 10^{-5} |

Table 1

### 4 DISCUSSION AND CONCLUSION

We have studied two mechanism contributing to the process \( B_c \to D_s^*\gamma \). The short distance one(Fig.1.a) induced by electro-magnetic penguin has been estimated within PQCD framework which is suitable for the process involved large energy release and hard scattering. The distribution functions \( \phi_H(x) \) coincide with the usual nonrelativistic Coulomb wave functions of \( \bar{c}\bar{b} \) and \( c\bar{s} \) systems, and expected to work well in the present case.

An unique nature of \( B_c \) decays is that ”the spectator” is heavy. Assume that the velocities of the heavy spectator are equal to \( v_1 \) of \( B_c \) (before scattering) and to \( v_2 \) of
the final meson (after scattering). In two body $B_c$ decays, $v_1 \cdot v_2$ is always large. For $B_c \to D^*_s\gamma$, $v_1 \cdot v_2 = 1.7$. Certainly, there is no phase-space for the propagators appearing in Fig.1.a to be on-shell, so the imaginary part of $M_a$ is absent. This is different from the situations in ref.[8]. Here we recall that the momentum square of the hard scattering exchanged by gluon is about $3.6GeV^2$, which is large enough for PQCD analysis. The hard scattering process cannot be included conveniently in the soft hadronic process described by the wavefunction of the final bound state, that is an important reason why we cannot apply the commonly used models with spectator ansatz, for example, HQET[7], BSW, ISGW[28] models, to the two body $B_c$ decays. That is also one of the reasons why the commonly used models cannot reliably predict the processes of $B \to$ two light hadrons or $\gamma+$light hadron, since the large mass scale of the decaying $B$ meson and the nearly massless final states imply that the transition amplitude is governed by hard process. Hence we conclude that the spectator ansatz is broken down in such kind of process and it is misleading to apply straightforwardly the commonly used models with spectator ansatz to such processes discussed above. Another competing mechanism is the weak annihilation. We find that it is equally important as the former ones. The situation is different from that of the radiative weak $B^\pm$ decays which are overwhelmingly dominated by electromagnetic penguin contributions. The results stem from at least two reasons:

1. The compact size of $B_c$ meson enhances the importance of the annihilation decays.

2. In $B_c \to D^*_s\gamma$ case, the CKM factor of weak annihilation contribution is $|V_{cs}V_{cb}|$, but in $B^\pm \to K^*\gamma$ case, the CKM part is $|V_{ub}V_{us}|$ which is much smaller than the former.

We have neglected the contribution of the vector meson dominance (VMD). Using the methods developed in [29], combined with the PQCD calculation of the form factors of the process $B_c \to D^*_sJ/\Psi(\Psi')$ (following the conversion $J/\Psi(\Psi') \to \gamma$), we find that the VMD contribution is negligibly small, due to the small $J/\Psi(\Psi') \to \gamma$ coupling $e/g_{J/\Psi(\Psi')\gamma} \approx 0.025(0.016)$.

Finally, we want to say a few words about the possibility of observing the interesting process at Tevatron and at the CERN Large Hadron Collider (LHC). The numbers of $B_c$ produced at Tevatron and LHC are estimated to be [30] $16000$ (for $25Pb^{-1}$ integrated luminosities with cuts of $P_T(B_c) > 10GeV, |y(B_c)| < 1$) and $2.1 \times 10^8$ (for $100fb^{-1}$ integrated luminosities with cuts of $P_T(B_c) > 20GeV, |y(B_c)| < 2.5$), respectively. Using the
numbers and $Br(B_c \rightarrow D_s^* \gamma)$, we find that the decay channel is unobservable at Tevatron, but more than one thousand events of interest will be produced at LHC, so it can be well studied at LHC in the near future.

5 Acknowledgment

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Table and Figure Captions

Table.1. $Br(B_c \rightarrow D_s^* \gamma)$ is displayed for various $B_c$ lifetime. The second row is for the contribution of electromagnetic penguin. The third row is for the contribution of annihilation. The final line is for the total results. Fig.1.a. The leading penguin graphs contribute to $B_c \rightarrow D_s^* \gamma$.

Fig.1.b. The leading W-annihilations to $B_c \rightarrow D_s^* \gamma$. Contributions due to photon emission from other quarks are denoted by ellipses.
Fig. 1