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The rare annihilation decays $\bar{B}^0_{s,d} \to J/\psi\gamma$

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Abstract. We investigate the physics potential of the annihilation decays $\bar{B}^0_{s,d} \to J/\psi\,\gamma$ in the standard model and beyond. In a naive factorization approach, the branching ratios are estimated to be $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 1.395 \times 10^{-6}$ and $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 5.398 \times 10^{-8}$. In the framework of QCD factorization, we compute the non-factorizable corrections and get $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 5.795 \times 10^{-8}$, $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 2.435 \times 10^{-9}$. Future measurements of these decays would be useful for testing the factorization frameworks. The smallness of these decays in the SM makes them sensitive probes of new physics. As an example, we will consider the possible admixture of the $(V + A)$ charge current to the standard $(V - A)$ current. This admixture will give a significant contribution to the decays.

1 Introduction

Decays of B mesons to final states containing charmonium constitute a very sensitive laboratory for the study of electro-weak interactions, as well as the dynamics of the strong interactions. The semi-inclusive J/ψ productions in B decays have always raised tough challenges for understanding its large rate [1, 2], which requires a large contribution beyond the color singlet model for charmonium production at the m_b scale [3–7]. Currently its large production rate could be understood in the framework of non-relativistic QCD (NRQCD) effective field theory [8]. However, the momentum spectrum of J/ψ, especially the excess of slow J/ψ meson, is still hard to theoretically explain; it may reveal interesting phenomena of the possible intrinsic charm component of the B [9], the decay $B \to J/\psi$ baryon anti-baryon [10] and the production of the $s\bar{d}g$ hybrid [11]. The exclusive decays $B \to J/\psi K^{(*)}$ have also attracted very extensive theoretical and experimental studies, which involve more complicated strong dynamics. Recent studies [12,13] have shown that it is hard to account for its large rates and polarizations theoretically.

In this paper, we present a study of the radiative annihilation decays $\bar{B}^0_{s,d} \to J/\psi\gamma$, which are much rarer than $B \to J/\psi K^{(*)}$; however, they involve simpler hadronic dynamics. In a naive factorization approach, these decays involve form factors which are similar to those of the radiative leptonic decays [14–17]. It is shown that the form factors could be described simply in terms of a convolution of the B meson distribution function with a perturbative kernel [15]. Beyond the naive factorization approach, non-factorization contributions only arise from one loop vertex QCD corrections which can be calculated properly in the framework of QCD factorization [18]. We find $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 1.395 \times 10^{-6}$ and $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 5.398 \times 10^{-8}$ in the naive factorization approach. With the QCD factorization approach, these branching ratios are reduced to $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 5.795 \times 10^{-8}$ and $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 2.435 \times 10^{-9}$, because the effective coefficient a_2 gets smaller than in naive factorization. Such effects have also been found in two-body non-leptonic B decays [18–21], where the vertex type one loop QCD corrections make $|a_2|$ much smaller than in the naive one, i.e., $a_2 = C_2 + C_1/N_c$. Compared with two-body non-leptonic B decays, the troublesome hard spectator scattering contribution is absent and only the well defined vertex type non-factorizable corrections are encountered. To this extent, the decays $\bar{B}^0_{s,d} \to J/\psi\gamma$ could be used to test the factorization schemes. On the other hand, these decays could serve as probes for new physics activities at the low energy scale. As an example, we take these decays as probes of the chirality of weak currents induced in the decays $b \to c\bar{c}s(d)$.

This paper is organized as follows. In Sect. 2, we present our study of $\bar{B}^0_{s,d} \to J/\psi\gamma$ in the standard model (SM). In Sect. 3, we will calculate the effect of the possible admixture of the $(V + A)$ current $g_{\rm R}(\bar{q}_1 q_2)_{V+A}$ to the standard $(V - A)$ current $g_{\rm L}(\bar{q}_1 q_2)_{V-A}$ in the decays $\bar{B}^0_{s,d} \to J/\psi\,\gamma$. This admixture could lead to enhancement of the decays branching ratios. The effects of the small value of $g_{\rm R}/g_{\rm L}$ could show up being signals of the activity of new physics.

2 $\bar{B}^0_{s,d} \to J/\psi\,\gamma$ in the SM

We start our study from the effective Hamiltonian relevant to $\bar{B}^0_q \to J/\psi\,\gamma$ decays in the standard model (SM) [22]:

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$$\mathcal{H}_{\text{eff}} = \frac{G_{\text{F}}}{\sqrt{2}} \left\{ V_{cb}V^*_{cq}\left[C_1(\mu)\mathcal{O}^c_1(\mu) + C_2(\mu)\mathcal{O}^c_2(\mu)\right] - V_{tb}V^*_{tq}\sum_{i=3}^{10} C_i(\mu)\mathcal{O}_i(\mu) \right\}. \quad (1)$$

Here $q = s, d$ and C_i $(i = 1, \dots, 10)$ are the Wilson coefficients at next-to-leading order evaluated at the renormalization scale μ. The effective operators can be expressed explicitly as follows [22]:

$$\begin{aligned}
\mathcal{O}^c_1 &= (\bar{c}_\alpha b_\alpha)_{V-A} \otimes (\bar{q}_\beta c_\beta)_{V-A},\\
\mathcal{O}^c_2 &= (\bar{c}_\alpha b_\beta)_{V-A} \otimes (\bar{q}_\beta c_\alpha)_{V-A},\\
\mathcal{O}_3 &= (\bar{q}_\alpha b_\alpha)_{V-A} \otimes (\bar{c}_\beta c_\beta)_{V-A},\\
\mathcal{O}_4 &= (\bar{q}_\alpha b_\beta)_{V-A} \otimes (\bar{c}_\beta c_\alpha)_{V-A},\\
\mathcal{O}_5 &= (\bar{q}_\alpha b_\alpha)_{V-A} \otimes (\bar{c}_\beta c_\beta)_{V+A},\\
\mathcal{O}_6 &= (\bar{q}_\alpha b_\beta)_{V-A} \otimes (\bar{c}_\beta c_\alpha)_{V+A}, \qquad (2)\\
\mathcal{O}_7 &= \frac{3}{2}(\bar{q}_\alpha b_\alpha)_{V-A} \otimes e_c(\bar{c}_\beta c_\beta)_{V+A},\\
\mathcal{O}_8 &= \frac{3}{2}(\bar{q}_\alpha b_\beta)_{V-A} \otimes e_c(\bar{c}_\beta c_\alpha)_{V+A},\\
\mathcal{O}_9 &= \frac{3}{2}(\bar{q}_\alpha b_\alpha)_{V-A} \otimes e_c(\bar{c}_\beta c_\beta)_{V-A},\\
\mathcal{O}_{10} &= \frac{3}{2}(\bar{q}_\alpha b_\beta)_{V-A} \otimes e_c(\bar{c}_\beta c_\alpha)_{V-A},
\end{aligned}$$

where α and β are the SU(3) color indices.

Under naive factorization, the $\bar{B}^0_q \to J/\psi\ \gamma$ decays are represented by Fig. 1. The dominant mechanism is the radiation of the photon from the light quark in the B meson. Generally the amplitude is suppressed by one order of Λ_{QCD}/m_b because the J/ψ meson must be transversely polarized and the B meson is heavy. Radiation of the photon from the remaining three quark lines is further suppressed additionally by a power of $(\Lambda_{\text{QCD}}/m_b)$, which will be neglected in this paper.

In the heavy quark limit, the decay amplitude of $\bar{B}^0_q \to J/\psi\ \gamma$ to leading order is

$$A(\bar{B}^0_q \to J/\psi\gamma) = \frac{G_{\text{F}}}{\sqrt{2}}V_{cb}V^*_{cq}\bar{a}_q\sqrt{4\pi\alpha_e}f_{J/\psi}M_{J/\psi}F_V \times \left\{-\epsilon_{\mu\nu\rho\sigma}\eta^\mu_\perp\varepsilon^\nu_\perp v^\rho q^\sigma + \mathrm{i}\left[(\varepsilon_\perp\cdot\eta_\perp)(v\cdot q) - (\eta_\perp\cdot q)(\varepsilon_\perp\cdot v)\right]\right\}, \quad (3)$$

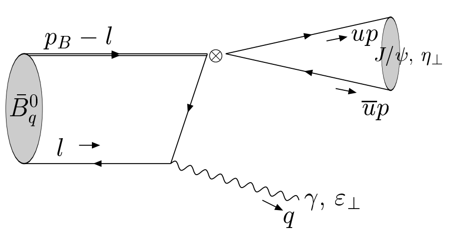


Fig. 1. The Feynman diagram for the leading contribution to $\bar{B}^0_q \to J/\psi\ \gamma$ decays. The photon radiated from the other quarks are suppressed by a power of $\mathcal{O}(\Lambda_{\text{QCD}}/m_b)$

where the approximation $V_{tb}V^*_{tq} \approx -V_{cb}V^*_{cq}$ has been made, and the parameter $\bar{a}_q$ is defined by

$$\bar{a}_q = a_2 + a_3 + a_5 + a_7 + a_9, \quad (4)$$

with $a_{2i} = C_{2i} + \frac{1}{N_c}C_{2i-1}$ and $a_{2i-1} = C_{2i-1} + \frac{1}{N_c}C_{2i}$. $\varepsilon_\perp$ and $\eta_\perp$ are the transverse polarization vectors of photon and J/ψ meson, respectively. The form factor F_V is defined by [14–17]

$$\langle\gamma(\varepsilon_\perp, q)|(\bar{q}b)_{V-A}|\bar{B}^0_q\rangle = \sqrt{4\pi\alpha_e} \times \left[-F_V\epsilon_{\mu\nu\rho\sigma}\varepsilon^\nu_\perp v^\rho q^\sigma + \mathrm{i}F_A(\varepsilon_{\perp\mu}v\cdot q - q_\mu\varepsilon_\perp\cdot v)\right]. \quad (5)$$

At leading order of $\mathcal{O}(1/m_b)$, the two form factors are given by

$$F_A = F_V = \frac{Q_q f_B M_B}{2\sqrt{2}E_\gamma\lambda_B}, \quad (6)$$

and λ_B is the first inverse moment of the B meson's distribution amplitude:

$$\frac{1}{\lambda_B} = \int_0^\infty \mathrm{d}l_+ \frac{\phi^B_1(l_+)}{l_+}. \quad (7)$$

Now we can write down the helicity amplitude

$$\begin{aligned}
\mathcal{M}_{++} &= -\frac{G_{\text{F}}}{\sqrt{2}}V_{cb}V^*_{cq}\bar{a}_q\sqrt{4\pi\alpha_e}F_V f_{J/\psi}M_{J/\psi}(2\mathrm{i}E_\gamma),\\
\mathcal{M}_{--} &= 0. \qquad (8)
\end{aligned}$$

From the helicity amplitude in (8), the branching ratio reads

$$\mathcal{B}(\bar{B}^0_q \to J/\psi\gamma) = \frac{\tau_{B_q}|P_c|}{8\pi M^2_B}\left(|\mathcal{M}_{--}|^2 + |\mathcal{M}_{++}|^2\right); \quad (9)$$

here P_c is the CM momentum of J/Ψ, and τ_{B_q} is the lifetime of the $\bar{B}^0_q$ meson.

For numerical analysis, we use the following input parameters [23]:

$$\begin{aligned}
M_{B_s} &= 5.370\,\text{GeV}, & \tau_{B_s} &= 1.461\,\text{ps},\\
m_b &= 4.8\,\text{GeV}, & V_{cb} &= 0.0412,\\
M_{B_d} &= 5.279\,\text{GeV}, & \tau_{B_d} &= 1.542\,\text{ps},\\
m_c &= 1.47\,\text{GeV}, & V_{cd} &= 0.224,\\
M_{J/\psi} &= 3.097\,\text{GeV}, & \lambda_B &= 0.35\,\text{GeV} \quad [15,18],\\
N_c &= 3, & V_{cs} &= 0.996,
\end{aligned}$$

and the decay constants $f_{B_s} = 210\,\text{MeV}$, $f_{B_d} = 180\,\text{MeV}$, $f_{J/\psi} = 405\,\text{MeV}$, and the Wilson coefficients at the $\mu = m_b$ scale: $C_1 = 1.082, C_2 = -0.185, C_3 = 0.014, C_4 = -0.035, C_5 = 0.009, C_6 = -0.041, C_7 = -\frac{0.002}{137}, C_8 = \frac{0.054}{137}, C_9 = -\frac{1.292}{137}, C_{10} = -\frac{0.263}{137}$ [22].

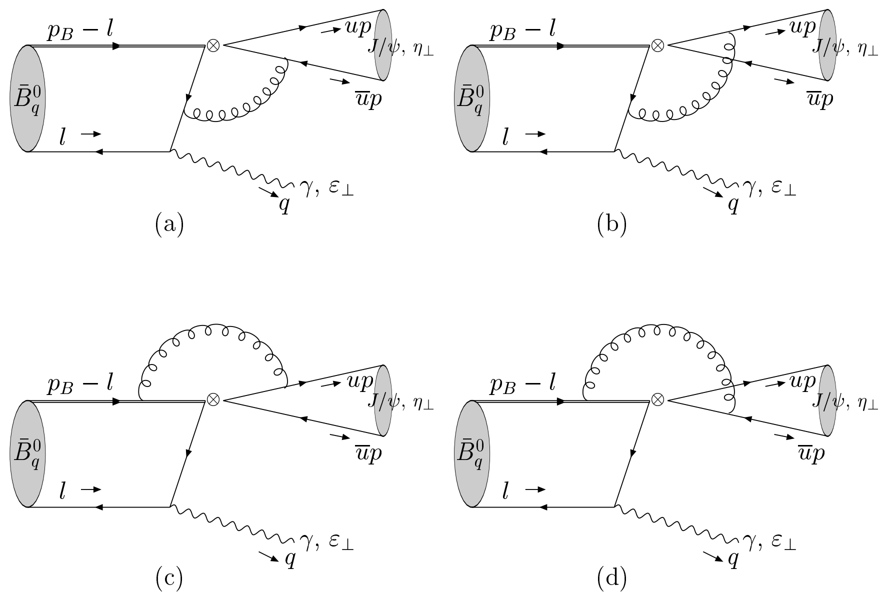


Fig. 2a–d. Non-factorizable contribution at order α_s. $\otimes$ denotes the insertions of the color-octet operators $\mathcal{O}^c_1, \mathcal{O}_4, \mathcal{O}_6, \mathcal{O}_8, \mathcal{O}_{10}$ in $\bar{B}^0_q \to J/\psi\ \gamma$. Other diagrams with the photon radiating from the remaining three quark lines are suppressed and neglected

In the naive factorization, we get the branching ratios

$$\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 1.395 \times 10^{-6}, \tag{10}$$

$$\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 5.398 \times 10^{-8}. \tag{11}$$

In the above calculations, non-factorizable contributions are neglected. However, the non-factorizable contributions may be important. These radiative corrections at order α_s can be obtained by calculating the amplitudes in Fig. 2. The QCD factorization approach advocated recently in [18] allows us to compute the non-factorizable corrections in the heavy quark limit.

In our calculation, we take the momentum of the B meson $P^\mu_B = M_B v^\mu$ and the photon flying along the $n_- = (1,0,0,-1)$ direction, where the four-velocity $v = (1,0,0,0)$ satisfies $v^2 = 1$. In the heavy quark limit, the B meson's light-cone projection operator can be written as [19, 24]

$$M^B_{\alpha\beta} = \frac{\mathrm{i}}{4N_c} f_B M_B \left\{(1+\not v)\gamma_5\left[\Phi^B_1(\rho) + \not n_- \Phi^B_2(\rho)\right]\right\}_{\alpha\beta}, \tag{12}$$

where ρ is the momentum fraction carried by the spectator quark of the B meson and the normalization conditions are

$$\int_0^1 \mathrm{d}\rho \Phi^B_1(\rho) = 1, \qquad \int_0^1 \mathrm{d}\rho \Phi^B_2(\rho) = 0. \tag{13}$$

For the J/ψ meson, we take

$$M^{J/\psi}_{\perp\rho\sigma} = -\frac{f_{J/\psi}}{4N_c}\left[\not\eta_\perp(\not P_{J/\psi} + M_{J/\psi})\right]_{\rho\sigma}\Phi^{J/\psi}(u). \tag{14}$$

Because the charm quark is heavy, the wave function $\Phi^{J/\psi}(u)$ is a symmetric function under $u \to 1-u$ and should be sharply peaked around $u = 1/2$ [11].

The calculation of the non-factorizable contributions depicted in Fig. 2. is straightforward. Including the contributions, the amplitudes for $\bar{B}^0_q \to J/\psi\ \gamma$ will be

$$\begin{aligned} A(\bar{B}^0_q \to J/\psi\gamma) &= \frac{G_F}{\sqrt{2}} V_{cb} V^*_{cq} \bar{a}'_q \sqrt{4\pi\alpha_e} f_{J/\psi} M_{J/\psi} F_V \\ &\times \{-\epsilon_{\mu\nu\rho\sigma}\eta^\mu_\perp \varepsilon^\nu_\perp v^\rho q^\sigma + \mathrm{i}\left[(\varepsilon_\perp\cdot\eta_\perp)(v\cdot q)\right. \\ &\left. - (\eta_\perp\cdot q)(\varepsilon_\perp\cdot v)\right]\}, \end{aligned} \tag{15}$$

and

$$\bar{a}'_q = a'_2 + a'_3 + a'_5 + a'_7 + a'_9. \tag{16}$$

The $\mathcal{O}(\alpha_s)$ corrections are summarized in a'_i which are calculated to be

$$\begin{aligned} a'_2 &= C_2 + \frac{C_1}{N_c} + \frac{\alpha_s}{4\pi}\frac{C_F}{N_c} C_1 F, \\ a'_3 &= C_3 + \frac{C_4}{N_c} + \frac{\alpha_s}{4\pi}\frac{C_F}{N_c} C_4 F, \\ a'_5 &= C_5 + \frac{C_6}{N_c} - \frac{\alpha_s}{4\pi}\frac{C_F}{N_c} C_6 F, \\ a'_7 &= C_7 + \frac{C_8}{N_c} - \frac{\alpha_s}{4\pi}\frac{C_F}{N_c} C_8 F, \\ a'_9 &= C_9 + \frac{C_{10}}{N_c} + \frac{\alpha_s}{4\pi}\frac{C_F}{N_c} C_{10} F, \end{aligned} \tag{17}$$

where $C_F = (N_c^2 - 1)/(2N_c)$, the α_s terms are the non-factorizable contributions which come from one gluon exchange between the two currents of the color-octet operators $\mathcal{O}^c_1, \mathcal{O}_4, \mathcal{O}_6, \mathcal{O}_8, \mathcal{O}_{10}$, as shown by Fig. 2; we have

$$\begin{aligned} F &= -16 - 12\ln\frac{\mu}{M_B} - 18\mathrm{i}\pi \\ &+ \int_0^1 \mathrm{d}u \Phi^{J/\psi}(u)\left\{\left(\frac{5-6u}{1-u}\right)\ln(u)\right. \\ &+ \frac{1-z}{1-z+uz}\ln(1-z) \\ &+ \left(\frac{1}{1-z+uz} - \frac{5}{1-uz}\right) uz\ln(uz) \end{aligned}$$

$$+\left(\frac{2}{1-u}+\frac{10}{1-uz}+\frac{1-z}{1-z+uz}\right)\mathrm{i}\pi\bigg\}, \quad (18)$$

where $z=\frac{M^2_{J/\psi}}{M^2_B}$. We have neglected the difference between m_b and M_B which is a sub-leading effect in the heavy quark limit.

In the calculation, the $\overline{\mathrm{MS}}$ renormalization scheme is used. We neglect the small effect of box diagrams and also neglected l_+^2 terms entering in the loop calculation which are higher twist effects. Therefore, the integral involving $\Phi_2^B(\rho)$ is absent and the remaining integrals could be related to the form factor F_V.

To order of α_s corrections, the helicity amplitudes are

$$\mathcal{M}_{++}=-\frac{G_F}{\sqrt{2}}V_{cb}V^*_{cq}\bar{a}'_q\sqrt{4\pi\alpha_e}F_V f_{J/\psi}M_{J/\psi}(2\mathrm{i}E_\gamma), \quad (19)$$

$$\mathcal{M}_{--}=0. \quad (20)$$

For $\Phi^{J/\psi}(u)=6u(1-u)$, we obtain

$$\mathcal{B}(\bar{B}^0_s\to J/\psi\gamma)=5.795\times10^{-8}, \quad (21)$$

$$\mathcal{B}(\bar{B}^0_d\to J/\psi\gamma)=2.435\times10^{-9}. \quad (22)$$

Because the shape of the J/ψ wave function is unknown, it is worth considering other possibilities besides the asymptotic form. The delta-function form $\Phi^{J/\psi}(u)=\delta\left(u-\frac{1}{2}\right)$ as used in [25] appeals to the naive expectation of the wave function in the non-relativistic limit. Using this wave function, the results are

$$\mathcal{B}(\bar{B}^0_s\to J/\psi\gamma)=9.772\times10^{-8}, \quad (23)$$

$$\mathcal{B}(\bar{B}^0_d\to J/\psi\gamma)=3.464\times10^{-9}. \quad (24)$$

The $\bar{B}^0_s\to J/\psi\gamma$ decays may be measured at Tevatron and LHC in the future. The decays $\bar{B}^0_d\to J/\psi\gamma$ could be studied in the planned Super B factories at KEK and SLAC succeeding Belle and BaBar. Potentially they could be enhanced by new physics and the enhancement might be measured at those facilities.

We note that the two different distribution functions of J/Ψ give quite different numerical results. The asymptotic form $\Phi^{J/\Psi}(u)=6u(1-u)$ is the relativistic limit one which is the analog to that of the π meson, while the delta-function like $\Phi^{J/\Psi}(u)=\delta\left(u-\frac{1}{2}\right)$ is the non-relativistic limit approach. In principle the distribution function of J/Ψ could be available from lattice simulations of QCD in the near future. Once the distribution function is available, theoretical predictions in this paper should be refined. We expect that the refinement would be made before the running of the CERN Large Hadron Collider (LHC) and Super B factories.

3 The admixture of $(V+A)$ current in $\bar{B}^0_{s,d}\to J/\psi\gamma$

B decays are known to be governed by weak couplings and small mixing matrix elements. These decays are therefore very sensitive to new kinds of interactions and in particular to right-handed couplings [26]. It is conventionally assumed that the B decays proceed via the pure $(V-A)$ current in the SM. However, it turns out to be surprisingly difficult to exclude the possibility that the dominant B decays occur via a $(V+A)$ coupling [27]. The $(V+A)$ coupling has been studied in [28–30]. In the annihilation decays $\bar{B}^0_{s,d}\to J/\psi\,\gamma$, there is a possible admixture of the $(V+A)$ charged current $g_R(\bar{q}_1q_2)_{V+A}$ to the standard $(V-A)$ current $g_L(\bar{q}_1q_2)_{V-A}$. A small value of g_R/g_L is not ruled out so far and can be sought for as one possible sign of new physics. In what follows, we will examine the effect of a possible admixture.

Assuming the admixture of $(V+A)$ charge currents $(b\to c)$ and $(c\to q)$ to the SM $(V-A)$ currents, the effective four-fermion interaction operators for $\bar{B}^0_q\to J/\psi\,\gamma$ here can be written as

$$\begin{aligned}\mathcal{O}^c_1&=[(\bar{c}_\alpha b_\alpha)_{V-A}+\xi(\bar{c}_\alpha b_\alpha)_{V+A}]\otimes\\&\quad[(\bar{q}_\beta c_\beta)_{V-A}+\xi'(\bar{q}_\beta c_\beta)_{V+A}],\\ \mathcal{O}^c_2&=[(\bar{c}_\alpha b_\beta)_{V-A}+\xi(\bar{c}_\alpha b_\beta)_{V+A}]\otimes\\&\quad[(\bar{q}_\beta c_\alpha)_{V-A}+\xi'(\bar{q}_\beta c_\alpha)_{V+A}],\end{aligned} \quad (25)$$

with $\xi=g_R/g_L$ for the current $(b\to c)$, and $\xi'=g'_R/g'_L$ for the current $(c\to q)$. Here we use the approximation $\tilde{a}_q=a_2$ and $\tilde{a}'_q=a'_2$.

The amplitude for $\bar{B}^0_q\to J/\psi\gamma$ is

$$\begin{aligned}A(\bar{B}^0_q\to J/\psi\gamma)&=\frac{G_F}{\sqrt{2}}V_{cb}V^*_{cq}\sqrt{4\pi\alpha_e}f_{J/\psi}M_{J/\psi}F_V\\&\times\{(1+\xi')a_2[-(1+\xi)\epsilon_{\mu\nu\rho\sigma}\eta^\mu_\perp\varepsilon^\nu_\perp v^\rho q^\sigma\\&\quad+\mathrm{i}(1-\xi)(\varepsilon_\perp\cdot\eta_\perp)(v\cdot q)]\\&+(1-\xi')\frac{\alpha_s}{4\pi}\frac{C_F}{N_C}C_1F[(\xi-1)\epsilon_{\mu\nu\rho\sigma}\eta^\mu_\perp\varepsilon^\nu_\perp v^\rho q^\sigma\\&\quad+\mathrm{i}(1+\xi)(\varepsilon_\perp\cdot\eta_\perp)(v\cdot q)]\}.\end{aligned} \quad (26)$$

The branching ratios can be read

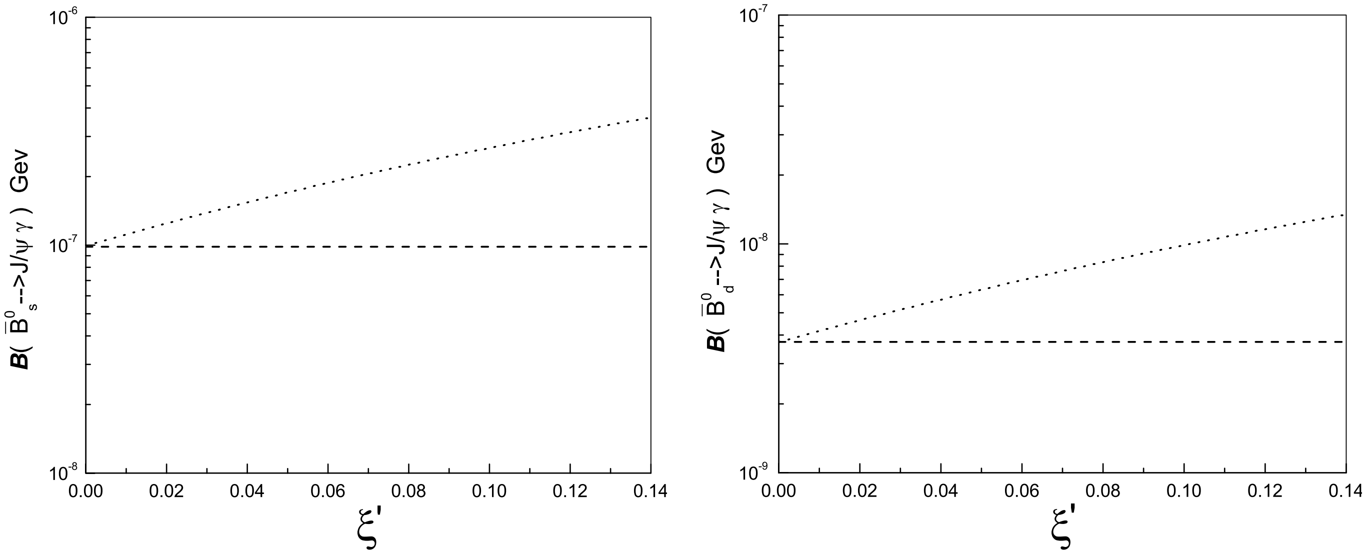
$$\begin{aligned}&\mathcal{B}(\bar{B}^0_q\to J/\psi\gamma)\\&=\frac{\tau_{B_q}|P_c|}{8\pi M_B^2}\frac{G_F^2}{2}\left|V_{cb}V^*_{cq}\right|^2 4\pi\alpha_e f^2_{J/\psi}M^2_{J/\psi}F_V^2 4E_\gamma^2\\&\quad\times\left\{\left|(1+\xi')a_2+(1-\xi')\frac{\alpha_s}{4\pi}\frac{C_F}{N_C}C_1F\right|^2\right.\\&\quad\left.+\xi^2\left|(1+\xi')a_2-(1-\xi')\frac{\alpha_s}{4\pi}\frac{C_F}{N_C}C_1F\right|^2\right\}.\end{aligned} \quad (27)$$

Here the factor ξ^2 stems from $\xi(\bar{c}b)_{V+A}$. Because ξ is small, in what follows we will neglect ξ^2.

Using $\Phi^{J/\psi}(u)=6u(1-u)$, the branching ratios are

$$\begin{aligned}&\mathcal{B}(\bar{B}^0_s\to J/\psi\gamma)\\&=(0.19+2.21\xi'+9.94\xi'^2)\times5.23\times10^{-7},\end{aligned} \quad (28)$$

$$\begin{aligned}&\mathcal{B}(\bar{B}^0_d\to J/\psi\gamma)\\&=(0.19+2.10\xi'+10.06\xi'^2)\times1.98\times10^{-8}.\end{aligned} \quad (29)$$



a **b**

Fig. 3a,b. $\mathcal{B}(\bar{B}^0_{s,d} \to J/\psi\ \gamma)$ as a function of $\xi' = g'_{\mathrm{R}}/g'_{\mathrm{L}}$. The dash lines are the SM predictions and the dot lines are the results of including a small admixture of the $(V+A)$ quark current

We can see that the branching ratios of these decays are very sensitive to the possible presence of the admixture of right-hand current. Comparing with the SM predictions, these decays are enhanced. The numerical results are displayed as an illustration in Fig. 3.

4 Conclusion

We have studied the radiative annihilation decays $\bar{B}^0_{s,d} \to J/\psi\ \gamma$ within the framework of the QCD factorization approach. Physically, the factorization method is applicable because the transverse size of J/ψ is small in the heavy quark limit. We have shown that the non-factorizable radiative corrections at order α_{s} change the magnitude significantly compared to the leading-order result corresponding to the naive factorization. In naive factorization, we find $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 1.395 \times 10^{-6}$ and $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 5.398 \times 10^{-8}$, which are much larger than $\mathcal{B}(\bar{B}^0_s \to J/\psi\gamma) = 5.795 \times 10^{-8}$ and $\mathcal{B}(\bar{B}^0_d \to J/\psi\gamma) = 2.435 \times 10^{-9}$ in the framework of QCD factorization at the order of α_{s}. It is interesting to note that these decays involve simpler hadronic dynamics than two-body B non-leptonic decays. Experimental measurements would be very useful for understanding the mechanics of J/ψ productions and testing the factorization frameworks. Simultaneously these decays also are the background for the interesting decays $\bar{B}^0_{s,d} \to \mu^+\mu^-\gamma$. On the other hand, these decays may be sensitive to new physics. As an illustration, we have investigated the effects of the admixture of right-hand currents. We find that these decays are sensitive to admixture of the right-handed $(c \to s, d)$ current and the effect of the admixture of the right-handed $(b \to c)$ current is negligibly small. Experimentally these decays could be studied at CERN LHC and the planned Super high luminosity B factories at KEK and SLAC.

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