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Analysis of direct CP violation in $B^- \to D^0 D_s^-, D^0 D^-$ decays

A.K. Giri[1], R. Mohanta[2], M.P. Khanna[1]

[1] Physics Department, Panjab University, Chandigarh-160 014, India
[2] School of Physics, University of Hyderabad, Gachibowli, Hyderabad-500 046, India

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Abstract. We investigate the possibility of observing direct CP violation in the decay modes $B^- \to D^0 D_s^-$ and $D^0 D^-$ within the standard model. Including the contributions arising from the tree, annihilation, QCD as well as electroweak penguins with both time- and space-like components, we find that the direct CP asymmetry in $B^- \to D^0 D_s^-$ is very small, $\sim 0.2\%$, but in $B^- \to D^0 D^-$ decay it can be as large as 4%. Approximately 10^7 charged B mesons are required to experimentally observe the CP asymmetry parameter for the latter case. Since this is easily accessible with the currently running B-factories, the decay mode $B^- \to D^0 D^-$ may be pursued to look for CP violation.

1 Introduction

CP violation is one of the least understood phenomena in particle physics [1–3], although it was observed in the K^0–$\bar{K}^0$ mixing system more than 35 years ago. In the standard model (SM), CP violation arises from a complex phase in the Cabibbo–Kobayashi–Maskawa quark mixing matrix [4]. Outside the kaon system, decays of B-mesons provide a rich ground for investigating CP violation [5,6]. Within the SM, the CP violation is often characterized by the so-called unitarity triangle [7]. By measuring CP violating rate asymmetries in B decays, one can extract α, β and γ, the three interior angles of the unitarity triangle. The sum of these three angles must be equal to 180° in the SM with three generations. At present we are at the beginning of the B-factory era in particle physics, which will provide us valuable insights in the phenomena of CP violation. One of the main programs of the presently running and the upcoming B-factories is to measure the size of CP violation in as many B decay modes as possible so as to establish the pattern of CP violation in various B decays. Among the most interesing B decay channels, the "gold plated" mode $B_d \to J/\psi K_s$, [8] allows the determination to be made of the angle β of the unitarity triangle of CKM matrix. Recent measurement of the CP asymmetry in the $B^0 \to J/\psi K^0$ and other related processes e.g. $\psi' K^0, \eta_c K^0$ etc. by the BELLE [9] and BaBar [10] detectors at the KEK and SLAC B-factories together with the earlier measurement of CDF [11] constitute the first significant signal of CP violation outside the neutral kaon system. The charmless rare B decays such as $B \to \pi\pi$, πK etc. are also potentially important for the study of CP violation, as widely discussed in the literature [12]. These decays in general proceed through two types of amplitudes: $b \to u$ tree amplitudes and $b \to s/d$ penguin amplitudes. The interference of these two amplitudes can give large CP violating asymmetries provided the strong FSI phase differences between these two amplitudes are not too small. Also a global analysis of the branching ratios and direct CP asymmetries in these decays can yield interesting information of the flavor sector of the standard model and at the same time provide a window to new physics. The CLEO as well as the BELLE collaborations have recently reported the observation of some rare two body decays of the type $B \to \pi K$ as well as upper bounds for the decay modes $B \to \pi\pi$ and $B \to K\bar{K}$ [13].

While the most promising proposal for observing CP violation in the B-system involves the mixing between neutral B-mesons [1], the decays of charged B-mesons are also of particular importance for establishing the detailed nature of CP violation. Since charged B-mesons cannot mix, a measurement of the CP violating observable in these decays would be a clear sign of "direct CP violation" which has been searched for in K-system for quite long with indefinite success. Only recently, such a kind of CP violating effect has been observed in the K-system by the NA48 [14] and KTeV [15] collaborations. For the bottom meson case usually the charmless rare B decay modes are preferred to study the direct CP violation, as these decay modes proceed with more than one Feynman diagrams, in particular the decay mode $B^+ \to \pi^+ K$, which can provide a direct CP asymmetry at the 20% level, if the strong rescattering phase difference is significantly large. Recently there has been significant progress in the theoretical understanding of the hadronic decays $B \to \pi K$, and methods have been developed to extract information on the CKM angle γ from the rate measurements for these processes [16]. In this paper we would like to look for some additional decay channels which could help us in establishing the presence of direct CP violation as quickly as possible. With this purpose we investigate the direct CP violating effects in the decays of charged B-mesons to two charmed mesons, i.e. $B^- \to D^0 D_s^-$ and $D^0 D^-$. It is worth emphasizing that these decay modes are flavor self-tagging processes which should be favored for experimental reconstustions. The decay mode $B^- \to D^0 D_s^-$ has already been observed experimentally with a branching ratio $(1.3 \pm 0.4)\%$ and the upper limit for the $B^- \to D^0 D^-$ channel is found to be $< 6.7 \times 10^{-3}$ [17]. These decay modes, which are described by quark level transitions like $b \to c\bar{c}q$ ($q = s/d$ for D_s^-/D^- in the final state), proceed through three distinct types of flavor topologies. These are the color allowed but Cabibbo suppressed tree, annihilation and the QCD as well as electroweak penguin diagrams. To get significant direct CP violation one would require two interfering amplitudes of comparable strengths, with different strong and weak phases. The weak phases arise from the superposition of various penguin contributions and the usual tree diagrams. The strong phases are generated by the perturbative penguin loops (hard final state interaction) [18] or final state interactions involving two different isospins. Since the decay modes we considered here are single isospin channels, i.e. the final states $D^0 D_s^-$ and $D^0 D^-$ are with isospin $I = 1/2$ and 1 respectively, the second type of FSI strong phase differences are absent for these channels. Therefore at first sight it appears that direct CP violating effects in these channels would be negligibly small as the tree contribution dominates over the other diagrams and thus have been overlooked in the literature. But detailed calculation shows that this is actually not so. In fact, the CP violating effects in the $B^- \to D^0 D^-$ channel can be as large as a few percent, which can be experimentally accessible in the first round of B-factories. The reason for the existence of such a significant CP violating parameter may be the fact that although the tree diagram for the $b \to c\bar{c}d$ transition is color allowed, it is doubly Cabibbo suppressed, and hence its magnitude is not very much larger than the penguin contributions. However, in contrast to the $B \to \pi K$ channel, the decay mode $B^- \to D^0 D^-$ cannot allow us to determine the angle γ from the corresponding CP violating effects. CP violating effects in the decays of the neutral B-meson into double charmed mesons have been extensively studied in [19–22], where it has been shown that these channels can be used as a method alternative to the $J/\psi K_s$ mode for the extraction of the angle β.

In our analysis, we use the standard theoretical framework to study the non-leptonic $B^- \to D^0 D_s^-(D^-)$ decay modes, which is based on the effective Hamiltonian approach in conjuction with the factorization hypothesis. The short distance QCD corrected Hamiltonian is calculated to next-to-leading order. The renormalization scheme and scale problems with the factorization approach for matrix elements can be circumvented by employing the scale and scheme independent effective Wilson coefficients. In the literature the contributions of space-like penguins are neglected assuming form factor suppression. But as pointed out in [23] the effect of space-like penguin amplitudes can be remarkably enhanced by the hadronic matrix elements involving $(V-A)(V+A)$ or $(S+P)(S-P)$ currents. Therefore we have included the space- and time-like contributions of both QCD and EW penguins, the annihilation contribution in addition to the dominant tree diagrams. Assuming the factorization approximation, the matrix elements of the tree and time-like penguin diagrams have been calculated in the BSW model [24], whereas for the evaluation of the matrix elements of the space and annihilation diagrams we have employed the Lepage and Brodsky model [25].

This paper is organized as follows. In Sect. 2 we briefly discuss the effective Hamiltonian, together with the quark level matrix elements and the numerical value of the Wilson coefficients in the effective Hamiltonian approach. Assuming the factorization approximation, the matrix elements of tree and time-like penguins are evaluated in the BSW model and for the space-like and annihilation diagrams we use the LB (Lepage and Brodsky) model. The determination of the CP violating asymmetry is presented in Sect. 3 and Sect. 4 contains our conclusion.

2 Framework

The effective Hamiltonian $\mathcal{H}_{\text{eff}}$ for the decay modes $B^- \to D^0 D_s^-$ and $D^0 D^-$ which are described by the quark level transitions $b \to c\bar{c}q$ (where $q = s$ for the former and d for the latter) have three classes of flavor topologies: the dominant tree, annihilation, and both QCD as well as electroweak penguins given by [5]

$$\mathcal{H}_{\text{eff}} = \frac{G_{\text{F}}}{\sqrt{2}} \Bigg\{ \lambda_u [c_1(\mu) O_1^u(\mu) + c_2(\mu) O_2^u(\mu)] + \lambda_c [c_1(\mu) O_1^c(\mu) + c_2(\mu) O_2^c(\mu)] + (\lambda_u + \lambda_c) \sum_{i=3}^{10} c_i(\mu) O_i(\mu) \Bigg\} + \text{h.c.}, \tag{1}$$

where $\lambda_u = V_{ub} V_{uq}^*$ and $\lambda_c = V_{cb} V_{cq}^*$, and $c_i(\mu)$ are the Wilson coefficients evaluated at the renormalization scale μ. The four fermion operators O_{1-10} are given by

$$\begin{aligned}
O_1^u &= (\bar{u}b)_{V-A} (\bar{q}u)_{V-A}, \\
O_2^u &= (\bar{u}_\alpha b_\beta)_{V-A} (\bar{q}_\beta u_\alpha)_{V-A}, \\
O_1^c &= (\bar{c}b)_{V-A} (\bar{q}c)_{V-A}, \\
O_2^c &= (\bar{c}_\alpha b_\beta)_{V-A} (\bar{q}_\beta c_\alpha)_{V-A}, \\
O_{3(5)} &= (\bar{q}b)_{V-A} \sum_{q'} (\bar{q}' q')_{V-A(V+A)}, \\
O_{4(6)} &= (\bar{q}_\alpha b_\beta)_{V-A} \sum_{q'} (\bar{q}'_\beta q'_\alpha)_{V-A(V+A)}, \\
O_{7(9)} &= \frac{3}{2} (\bar{q}b)_{V-A} \sum_{q'} e_{q'} (\bar{q}' q')_{V+A(V-A)}, \\
O_{8(10)} &= \frac{3}{2} (\bar{q}_\alpha b_\beta)_{V-A} \sum_{q'} e_{q'} (\bar{q}'_\beta q'_\alpha)_{V+A(V-A)},
\end{aligned} \tag{2}$$

where $O_{1,2}$ are the tree level current–current operators, $O_{3\text{–}6}$ the QCD penguin operators and $O_{7\text{–}10}$ the EW penguin operators. $(\bar{q}_1 q_2)_{(V\pm A)}$ denote the usual $(V \pm A)$ currents. The sum over q' runs over the quark fields that are active at the scale $\mu = O(m_b)$, i.e., $(q' \in u, d, s, c, b)$. The Wilson coefficients depend (in general) on the renormalization scheme and the scale μ at which they are evaluated. In the next-to-leading order their values are obtained in the naive dimensional regularization (NDR) scheme at $\mu = m_b(m_b)$ [26]: $c_1 = 1.082, c_2 = -0.185$, $c_3 = 0.014, c_4 = -0.035, c_5 = 0.009, c_6 = -0.041$, $c_{7/\alpha} = -0.002, c_{8/\alpha} = 0.054, c_{9/\alpha} = -1.292$ and $c_{10/\alpha} = 0.263$.

However, the physical matrix elements $\langle P_1 P_2 | \mathcal{H}_{\text{eff}} | B \rangle$ are obviously independent of both the scheme and the scale. Hence the dependence on the Wilson coefficients must be compensated by a commensurate calculation of the hadronic matrix elements in a non-perturbative framework such as lattice QCD. Presently, this is not a viable strategy as the calculation of the matrix elements $\langle P_1 P_2 | O_i | B \rangle$ is beyond the scope of the current lattice technology. However, perturbation theory comes to (partial) rescue; with the help of it one-loop matrix elements can be rewritten in terms of the operators and the effective Wilson coefficients c_i^{eff} which are scheme and scale independent:

$$\langle q q' \bar{q}' | \mathcal{H}_{\text{eff}} | b \rangle = \sum_{i,j} c_i^{\text{eff}}(\mu) \langle q q' \bar{q}' | O_j | b \rangle^{\text{tree}}. \tag{3}$$

The effective Wilson coefficients c_i^{eff} may be expressed by [27]

$$\begin{aligned}
c_1^{\text{eff}}|_{\mu=m_b} &= c_1(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{1i} c_i(\mu), \\
c_2^{\text{eff}}|_{\mu=m_b} &= c_2(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{2i} c_i(\mu), \\
c_3^{\text{eff}}|_{\mu=m_b} &= c_3(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{3i} c_i(\mu) \\
&\quad -\frac{\alpha_s}{24\pi}(C_t + C_p + C_g), \\
c_4^{\text{eff}}|_{\mu=m_b} &= c_4(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{4i} c_i(\mu) \\
&\quad +\frac{\alpha_s}{8\pi}(C_t + C_p + C_g), \\
c_5^{\text{eff}}|_{\mu=m_b} &= c_5(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{5i} c_i(\mu) \\
&\quad -\frac{\alpha_s}{24\pi}(C_t + C_p + C_g), \\
c_6^{\text{eff}}|_{\mu=m_b} &= c_6(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{6i} c_i(\mu) \\
&\quad +\frac{\alpha_s}{8\pi}(C_t + C_p + C_g), \\
c_7^{\text{eff}}|_{\mu=m_b} &= c_7(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{7i} c_i(\mu) \\
&\quad +\frac{\alpha}{8\pi} C_e, \\
c_8^{\text{eff}}|_{\mu=m_b} &= c_8(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{8i} c_i(\mu), \\
c_9^{\text{eff}}|_{\mu=m_b} &= c_9(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{9i} c_i(\mu) \\
&\quad +\frac{\alpha}{8\pi} C_e, \\
c_{10}^{\text{eff}}|_{\mu=m_b} &= c_{10}(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)\mathrm{T}} \ln\frac{m_b}{\mu} + \hat{r}^{\mathrm{T}}\right)_{10i} c_i(\mu),
\end{aligned} \tag{4}$$

where $\hat{r}^{\mathrm{T}}$ and $\gamma^{(0)\mathrm{T}}$ the transpose of the matrices $\hat{r}$ and $\gamma^{(0)}$, arise from the vertex corrections to the operators O_1–O_{10} derived in [28], which are explicitly given in [29].

The quantities C_t, C_p, C_e and C_g are arising from the peguin type diagrams of the operators $O_{1,2}$, the QCD penguin type diagrams of the operators O_3–O_6, the electroweak penguin type diagrams of $O_{1,2}$ and the tree level diagrams of the dipole operator O_g respectively, which are given in the NDR scheme (after $\overline{\text{MS}}$ renormalization) by

$$\begin{aligned}
C_t &= -\left(\frac{\lambda_u}{\lambda_t}\tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t}\tilde{G}(m_c)\right) c_1, \\
C_p &= [\tilde{G}(m_s) + \tilde{G}(m_b)] c_3 + \sum_{i=u,d,s,c,b} \tilde{G}(m_i)(c_4 + c_6), \\
C_g &= -\frac{2 m_b}{\sqrt{\langle k^2 \rangle}} c_g^{\text{eff}}, \\
c_g^{\text{eff}} &= -1.043, \\
C_e &= -\frac{8}{9}\left(\frac{\lambda_u}{\lambda_t}\tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t}\tilde{G}(m_c)\right)(c_1 + 3c_2), \\
\tilde{G}(m_q) &= \frac{2}{3} - G(m_q, k, \mu),
\end{aligned} \tag{5}$$

$$G(m, k, \mu) = -4\int_0^1 \mathrm{d}x x(1-x) \ln\left(\frac{m^2 - k^2 x(1-x)}{\mu^2}\right). \tag{6}$$

It should be noted that the quantities C_t, C_p C_e and C_g depend on the CKM matrix elements, the quark masses, the scale μ and k^2, the momentum transferred by the virtual particles appearing in the penguin diagrams. In the factorization approximation there is no model independent way to keep track of the k^2 dependence; the actual value of k^2 is model dependent. From simple kinematics [30] one expects k^2 to be typically in the range

$$\frac{m_b^2}{4} \le k^2 \le \frac{m_b^2}{2}. \tag{7}$$

Since the branching ratio and the CP asymmetry depend crucially on the parameter k^2, here we would like to take a specific value for it, based on the valence quark approximation instead of the conventionally used value $k^2 = m_b^2/2$. As discussed in [23] the averaged value of the squared momentum transfer for $B^-(b\bar{u}) \to D^0(c\bar{u}) D_q^-(q\bar{c})$ is given by

$$\langle k^2 \rangle = m_b^2 + m_q^2 - 2 m_b E_q, \tag{8}$$

Table 1. Numerical value of the effective Wilson coefficients $c_i^{\rm eff}$ for $b \to s$ and $b \to d$ transitions

	$b \to s$ Time-like	$b \to s$ Space-like	$b \to d$ Time-like	$b \to d$ Space-like
$c_1^{\rm eff}$	1.168	1.168	1.168	1.168
$c_2^{\rm eff}$	-0.366	-0.366	-0.366	-0.366
$c_3^{\rm eff}$	0.0225+i0.0044	$-(0.0096 + \mathrm{i}0.0003)$	0.0197+i 0.005	$-(0.0123 - \mathrm{i}0.0066)$
$c_4^{\rm eff}$	$-(0.0456 + \mathrm{i}0.0133)$	$(0.0505 + \mathrm{i}0.0009)$	$-(0.0373 + \mathrm{i}0.015)$	$(0.0586 - \mathrm{i}0.0199)$
$c_5^{\rm eff}$	0.0132+i0.0044	$-(0.0189 + \mathrm{i}0.0003)$	0.0104+i0.005	$-(0.0216 - \mathrm{i}0.0066)$
$c_6^{\rm eff}$	$-(0.0478 + \mathrm{i}0.0133)$	$(0.0483 + \mathrm{i}0.0009)$	$-(0.0395 + \mathrm{i}0.015)$	$(0.0564 - \mathrm{i}0.0199)$
$c_7^{\rm eff}/\alpha$	$-(0.0282 + \mathrm{i}0.0363)$	$-(0.0303 - \mathrm{i}0.0018)$	$-(0.0119 + \mathrm{i}0.0398)$	$-(0.0143 + \mathrm{i}0.0391)$
$c_8^{\rm eff}/\alpha$	0.055	0.055	0.055	0.055
$c_9^{\rm eff}/\alpha$	$-(1.4252 + \mathrm{i}0.0363)$	$-(1.4273 - \mathrm{i}0.0018)$	$-(1.4089 + \mathrm{i}0.0398)$	$-(1.4113 + \mathrm{i}0.0391)$
$c_{10}^{\rm eff}/\alpha$	0.481	0.481	0.481	0.481

where the energy of the quark q in the final D_q^--particle is determinable from

$$E_q + \sqrt{E_q^2 - m_q^2 + m_c^2} + \sqrt{4E_q^2 - 4m_q^2 + m_c^2} = m_b \quad (9)$$

for time-like penguin channels, or from

$$E_q + \sqrt{E_q^2 - m_q^2 + m_u^2} = m_b + m_u \quad (10)$$

for space-like penguin diagrams. m_b, m_q and m_c denote the masses of the decaying b-quark, daughter q-quark and the c-quark (created as a $c\bar{c}$ pair from the virtual gluon, photon or Z-particle in the penguin loop). For numerical calculations, we have taken the CKM matrix elements expressed in terms of the Wolfenstein parameters with values $A = 0.815$, $\lambda = \sin\theta_c = 0.2205$, $\rho = 0.175$ and $\eta = 0.37$ [29]. The choice of ρ and η correspond to the CKM triangle: $\alpha = 91°$, $\beta = 24°$ and $\gamma = 65°$. At the scale $\mu \sim m_b$, we use the current quark masses of [29] $m_u(m_b) = 3.2\,\mathrm{MeV}$, $m_d(m_b) = 6.4\,\mathrm{MeV}$, $m_s(m_b) = 90\,\mathrm{MeV}$, $m_c(m_b) = 0.95\,\mathrm{GeV}$ and $m_b(m_b) = 4.34\,\mathrm{GeV}$. With the specific value of k^2 obtained from (8)–(10), we obtain the values of the effective renormalization scheme and scale independent Wilson coefficients for the $b \to s$ and $b \to d$ transitions as given in Table 1.

Now we want to calculate the matrix element $\langle D_q^- D^0| O_i|B^-\rangle$ using the factorization approximation, where O_i are the four quark current operators listed above. In this approximation, the hadronic matrix elements of the four quark operators $(\bar{c}b)_{(V-A)}(\bar{q}c)_{(V-A)}$ split into the product of two matrix elements, $\langle D^0|(\bar{c}b)_{(V-A)}|B^-\rangle$ and $\langle D_q^-| (\bar{q}c)_{(V-A)}|0\rangle$, where a Fierz transformation has been used so that flavor quantum numbers of the currents match with those of the hadrons. Since Fierz rearranging yields operators which are in the color singlet–singlet and octet–octet forms, this procedure results, in general, in matrix elements which have the right flavor quantum numbers but involve both singlet–singlet and octet–octet current operators. However, there is no experimental information available for the octet–octet part. So in the factorization approximation, one discards the color octet–octet piece and compensates this by treating N_c, the numbers of colors as a free parameter, and its value is extracted from the data of two body non-leptonic decays.

The matrix elements of the $(V-A)(V+A)$ operators i.e. (O_6 and O_8) can be transformed into $(V-A)(V-A)$ form by using Fierz ordering and the Dirac equation, which are given by

$$\langle D_q^- D^0|O_6|B^-\rangle = R_q\langle D_q^- D^0|O_4|B^-\rangle, \quad (11)$$

with

$$R_q = \frac{2m_{D_q^-}}{(m_b - m_c)(m_q + m_c)}, \quad (12)$$

where the quark masses are the current quark masses. The same relation works for O_8.

Hence, one obtains the transition amplitude for $B^- \to D_s^- D^0$ and $D^- D^0$ as (the factor $G_{\rm F}/2^{1/2}$ is suppressed)

$$\begin{aligned} &A(B^- \to D_s^- D^0) \\ &= \lambda_u\Big\{(a_4 + a_{10} + (a_6 + a_8)R_s)X^{(BD^0,D_s^-)} \\ &\quad +(a_1 + a_4 + a_{10} + (a_6 + a_8)R_s')X^{(B,D^0D_s^-)}\Big\} \\ &\quad +\lambda_c\Big\{(a_1 + a_4 + a_{10} + (a_6 + a_8)R_s)X^{(BD^0,D_s^-)} \\ &\quad +(a_4 + a_{10} + (a_6 + a_8)R_s')X^{(B,D^0D_s^-)}\Big\}, \end{aligned} \quad (13)$$

$$\begin{aligned} &A(B^- \to D^-\bar{D}^0) \\ &= \lambda_u\Big\{(a_4 + a_{10} + (a_6 + a_8)R_d)X^{(BD^0,D^-)} \\ &\quad +(a_1 + a_4 + a_{10} + (a_6 + a_8)R_d')X^{(B,D^0D^-)}\Big\} \\ &\quad +\lambda_c\Big\{(a_1 + a_4 + a_{10} + (a_6 + a_8)R_d)X^{(BD^0,D^-)} \\ &\quad +(a_4 + a_{10} + (a_6 + a_8)R_d')X^{(B,D^0D^-)}\Big\}, \end{aligned} \quad (14)$$

where

$$\begin{aligned} X^{(BD^0,D_q^-)} &= \langle D_s^-|(\bar{q}c)|0\rangle\langle D^0|(\bar{c}b)|B\rangle, \\ X^{(B,D_q^-D^0)} &= \langle D^0D_q^-|(\bar{q}c)|0\rangle\ \langle 0|(\bar{u}b)|B\rangle. \end{aligned} \quad (15)$$

$X^{(BD^0, D_q^-)}$ denotes matrix elements of the tree and time-like penguins, whereas $X^{(B, D^0 D_q^-)}$ stand for the annihilation and space-like amplitudes. We have

$$R_{q'} = \frac{2m_B^2}{(m_q - m_u)(m_b + m_u)}, \tag{16}$$

which arises from the transformation of $(V-A)(V+A)$ operators into $(V-A)(V-A)$ form for space-like penguins. It should be noted that $\lambda_u = V_{ub}V_{us}^*$ for $B^- \to D^0 D_s^-$ whereas $\lambda_u = V_{ub}V_{ud}^*$ for $B^- \to D^0 D^-$ and similar expressions for λ_c. The coefficients $a_1, a_2 \cdots a_{10}$ are combinations of the effective Wilson coefficients given by

$$a_{2i-1} = c_{2i-1}^{\text{eff}} + \frac{1}{N_c^{\text{eff}}} c_{2i}^{\text{eff}}, \quad a_{2i} = c_{2i}^{\text{eff}} + \frac{1}{N_c^{\text{eff}}} c_{2i-1}^{\text{eff}}, \quad i = 1, 2 \cdots 5, \tag{17}$$

where N_c^{eff} is the effective number of colors treated as a free parameter in order to model the non-factorizable contributions to the matrix elements and its value can be extracted from the two body non-leptonic B decays. A recent analysis of $B \to D\pi$ data gives $N_c^{\text{eff}} \sim 2$ [31]. Therefore, in our analysis, we take two sets of values for N_c^{eff}, i.e., $N_c^{\text{eff}} = 2$ and $N_c^{\text{eff}} = 3$, which characterizes naive factorization.

The factorized hadronic matrix elements are evaluated using the BSW model [24], which are given by

$$X^{(BD^0, D_q^-)} = \mathrm{i} f_{D_q} F_0^{BD}(m_{D_q}^2)(m_B^2 - m_{D^0}^2). \tag{18}$$

The matrix element of the annihilation and space-like penguins are given by [23]

$$\begin{aligned} &\langle D^0 D_q^- | (\bar{q}u)(\bar{u}b) | B^- \rangle \\ &= \mathrm{i} f_B f_+^a(m_B^2) \left[m_{D_q}^2 - m_{D^0}^2 - \frac{m_{D_q} - m_{D^0}}{m_{D_q} + m_{D^0}} m_B^2 \right], \end{aligned} \tag{19}$$

where the value of the annihilation form factor is given by $f_+^a(m_B^2) = \mathrm{i}16\pi\alpha_s f_B^2 / m_B^2$ [25].

After obtaining the transition amplitude, the branching ratio is given by

$$\text{BR} = \frac{|\boldsymbol{p}|}{8\pi m_B^2} \frac{|A(B^- \to D^0 D_q^-)|^2}{\Gamma}, \tag{20}$$

where $|\boldsymbol{p}|$ is the momentum of the emitted particles and Γ is the total decay width.

Using (13)–(19) we obtain the transition amplitude (in units of $G_F/2^{1/2}$):

$$\begin{aligned} &A(B^- \to D^0 D_s^-) \\ &= \lambda_u(0.1898 - \mathrm{i}0.6483) + \lambda_c(0.1889 + \mathrm{i}4.418) \\ &\quad [\lambda_u(0.2019 - \mathrm{i}0.6817) + \lambda_c(0.201 + \mathrm{i}4.698)], \end{aligned} \tag{21}$$

$$\begin{aligned} &A(B^- \to D^0 D^-) \\ &= \lambda_u(0.2259 - \mathrm{i}0.5616) + \lambda_c(0.2259 + \mathrm{i}4.8185) \\ &\quad [\lambda_u(0.2393 - \mathrm{i}0.589) + \lambda_c(0.2393 + \mathrm{i}5.124)], \end{aligned} \tag{22}$$

where we have used the decay constants (in MeV) $f_{D_s} = 280$, $f_D = 300$ [17] and f_B=180 [32]. In the above equations, the upper values correspond to $N_c^{\text{eff}} = 2$ and the lower bracketed values to $N_c^{\text{eff}} = 3$.

3 CP violating asymmetry

For charged $B^\mp$ decays, the CP violating rate asymmetries in partial decay rates are defined by

$$a_{CP} = \frac{\Gamma(B^- \to f^-) - \Gamma(B^+ \to f^+)}{\Gamma(B^- \to f^-) + \Gamma(B^+ \to f^+)}. \tag{23}$$

As these decays are all self-tagging, the measurement of these CP violating asymmetries is essentially a counting experiment in well-defined final states. Their rate asymmetries require both weak and strong phase differences in the interfering amplitudes. The weak phase difference arises from the superposition of amplitudes from various tree (current–current) and penguin diagrams. The strong phases which are needed to obtain non-zero values for a_{CP}, are generated by absorptive parts in penguin diagrams (hard final state interactions).

For the B-meson decaying to a final state f and the charge conjugated $B^- \to f$ we may, without any loss of generality, write the transition amplitude as

$$A(f) = \lambda_u A_u \mathrm{e}^{\mathrm{i}\delta_u} + \lambda_c A_c \mathrm{e}^{\mathrm{i}\delta_c}, \tag{24}$$

$$\bar{A}(f) = \lambda_u^* A_u \mathrm{e}^{\mathrm{i}\delta_u} + \lambda_c^* A_c \mathrm{e}^{\mathrm{i}\delta_c}, \tag{25}$$

where $\lambda_i = V_{ib}V_{iq}^*$, A_u and A_c denote the contribution from tree and penguin operators proportional to the product of CKM matrix elements λ_u and λ_c respectively. The corresponding strong phases are denoted by δ_u and δ_c respectively. Thus the direct CP violating asymmetry is given by

$$\begin{aligned} a_{CP} &= \frac{-2\mathrm{Im}(\lambda_u \lambda_c^*)\mathrm{Im}(A_u A_c^*)}{|\lambda_u A_u|^2 + |\lambda_c A_c|^2 + 2\mathrm{Re}(\lambda_u \lambda_c^*)\mathrm{Re}(A_u A_c^*)} \\ &= \frac{2\sin\gamma \sin(\delta_u - \delta_c)}{\left|\frac{\lambda_u A_u}{\lambda_c A_c}\right| + \left|\frac{\lambda_c A_c}{\lambda_u A_u}\right| + 2\cos\gamma\cos(\delta_u - \delta_c)}, \end{aligned} \tag{26}$$

where the weak phases entering in the $b \to s/d$ transition is equal to $(-\gamma)$, as we are using the Wolfenstein approximation, in which λ_c has no weak phase and the phase of λ_u is $-\gamma$. The strong phase $(\delta_u - \delta_c)$ is caused by the final state interactions. The strong phases are given by

$$\sin(\delta_u - \delta_c) = \frac{1}{|A_u A_c|}(\mathrm{Im}A_u \mathrm{Re}A_c - \mathrm{Im}A_c \mathrm{Re}A_u), \tag{27}$$

$$\cos(\delta_u - \delta_c) = \frac{1}{|A_c A_u|}(\mathrm{Re}A_u \mathrm{Re}A_c + \mathrm{Im}A_u \mathrm{Im}A_c). \tag{28}$$

4 Conclusion

Using the next-to-leading order QCD corrected effective Hamiltonian, with the scale and scheme independent Wilson coefficients, we have systematically studied the two charm hadronic decay modes $B^- \to D^0 D_s^-$ and $D^0 D^-$ within the framework of the generalized factorization. The non-factorizable contributions are parameterized in terms of N_c^{eff}, the effective number of colors. For numerical calculations, we have used two different sets of values for these parameters:

Table 2. Branching ratio and CP asymmetry in % for $B^- \to D^0 D_s^-, D^0 D^-$ decay modes

Decay modes	Branching ratios (BR) $N_c^{\rm eff}=2$	$N_c^{\rm eff}=3$	Expt.	CP asymmetry $N_c^{\rm eff}=2$	$N_c^{\rm eff}=3$
$B^- \to D^0 D_s^-$	1.29×10^{-2}	1.46×10^{-2}	$(1.3 \pm 0.4) \times 10^{-2}$	0.18	0.18
$B^- \to D^0 D^-$	8.72×10^{-4}	9.86×10^{-4}	$< 6.7 \times 10^{-3}$	3.62	3.62

(1) $N_c^{\rm eff} = 2$,
(2) $N_c^{\rm eff} = 3$, which holds for naive factorization.

The existence of a direct CP violating rate asymmetry requires two interfering amplitudes having different CP non-conserving weak phases and CP conserving strong phases. The former may arise either from the standard model CKM matrix or from new physics, while the latter may arise from the absorptive part of a penguin diagram or from final state interaction effects of two different isospins. Since the channels we considered here are single isospin channels, the second class of strong phase differences do not arise for these channels. In our analysis, the weak phases are due to the CKM matrix and the strong phase differences arise due to the absorptive part of penguin diagrams. The branching ratio and the CP violating asymmetry parameter are estimated using (20) and (26) and are presented in Table 2.

From the results we have observed the following:

(i) The predicted branching ratio for the decay mode $B^- \to D^0 D_s^-$ agrees very well with the experimental value for $N_c^{\rm eff} = 2$, and the CP violating parameter for this mode is quite small.
(ii) The branching ratio for the decay mode $B^- \to D^0 D^-$ lies below the present experimental upper bound and the CP violating parameter for this mode is quite significant. The number of charged B-mesons required to observe this CP violating signal to three standard deviations is given by $N_B = 9/(\mathrm{BR} \times a_{CP}^2) \approx 7.9 \times 10^6$, which is easily accessible with the currently running B-factories.

Since we have obtained these results using model dependent calculations, in general they may suffer from some theoretical uncertainties. Let us briefly point out the possible sources of these uncertainties.

(1) Evaluation of matrix elements: We have used the generalized factorization approximation along with the BSW [24] and the Lepage and Brodsky [25] models to evaluate the transition matrix elements. The non-factorizable contributions are taken care of by considering the effective number of colors $N_c^{\rm eff} = 2$, obtained from the experimental data of $B \to D\pi$. Although the generalized factorization approximation and BSW model are rather successful in explaining the data on a number of exclusive B decays, there might be some amount of uncertainty introduced due to them.
(2) Running quark masses at the scale m_b: The current quark masses arise in the decay amplitudes because the equation of motion has been applied to the matrix elements obtained from the Fierz transformation of $(V-A)(V+A)$ penguin operators. Since the current quark masses are not known precisely, this will result in large uncertainties in the predicted results.
(3) There will be some amount of uncertainty due to the parameters (ρ, η, A, λ) of the CKM matrix.
(4) The magnitude of momentum transfer to the gluon/γ/Z in the penguin diagram: We have employed the valence quark approximation to fix the value of k^2 for calculating the value of the effective Wilson coefficients. The common argument is that while the CP violating rate asymmetry is quite sensitive to the value of k^2, this is not the case for decay rate.
(5) Strong FSI phases: In our analysis the strong phases are generated by the absorptive parts of the penguin diagrams. The value of the strong phase difference is quite sensitive to the CP violating rate asymmetry.

Although in general the theoretical predictions suffer from many hadronic uncertainties we conclude that there is a fair chance for observing direct CP violation asymmetry in the $B^- \to D^0 D^-$ decay channel. It has been emphasized in [19–22] that the neutral B-meson decay modes to two charmed mesons can be used to measure the unitarity angle β as an alternative to the gold plated mode $B \to J/\psi K$. We argue further here that the mode $B^- \to D^0 D^-$ can be used to quickly settle down the search for observing direct CP violation outside the kaon system, if the SM description of CP violation is correct, or else could provide us a clear indication of the presence of new physics. It should be noted here that the decay mode is flavor self-tagging and hence experimentally favorable. Furthermore, although the CP asymmetry is found to be rather small, i.e., at the level of 4%, the branching ratio is expected to be of order $\mathcal{O}(10^{-3})$, thereby making this channel accessible and interesting at the B-factories.

To summarize, since the modes we consider are direct decays and not time dependent, they may be observed in any experimental setting where a large number of B-mesons are produced. Apart from the SLAC and KEK asymmetric B-factories these include CLEO and hadronic B experiments such as HERA-b, BTeV, Collider Detectors at Fermilab (CDF), D0 and CERN LHC-b or high luminosity Z-factory. In particular, we emphasize that the decay mode $B^- \to D^0 D^-$ may be searched for in the first round of B-factory experiments (where it can be easily accessible) to observe direct CP violation or to provide us a hint for the presence of new physics.

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