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Understanding the nature of $D_s(2317)$ and $D_s(2460)$ through nonleptonic B decays

Alakabha Datta, P.J. O'Donnell

Department of Physics, University of Toronto, Toronto, ON M5S 1A7, Canada

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Abstract

We consider the nonleptonic B decays $B \to D^{(*)} D_s(2317)$ and $B \to D^{(*)} D_s(2460)$, involving the newly discovered $D_s(2317)$ and the $D_s(2460)$ states. We find that experiments indicate disagreement with model calculations of their properties and/or breakdown of the factorization assumption for these decays. We point out that decays involving B_s mesons where the D_s resonances can be produced via the weak decay of the b quark can provide further information about the nature of these newly discovered states. We also propose a model to calculate the two body nonleptonic decays $B \to D^{(*)} D_s(2317)(D_s(2460))$, if the $D_s(2317)$ and $D_s(2460)$ are interpreted as DK and $D^* K$ molecules.

1. Introduction

There has been recent observations of an unexpectedly light narrow resonance in $D_s^+ \pi^0$ with a mass of 2317 MeV$/c^2$ by the BaBar Collaboration [1], together with another second narrow resonance in $D_s \pi^0 \gamma$ with a mass 2460 MeV$/c^2$ [2].

The smaller than expected masses and narrow widths of these states have led, among other explanations [3], to a multiquark–antiquark or a DK molecule interpretation of these states [4], or to an interpretation as p-wave states where the light degrees of freedom are in an angular momentum state $j_q = \frac{1}{2}$ [5], or even some combination of these [6]. There are also conflicting lattice interpretations of these states [7]. The mass difference between the $D_s(2317)$ and the well established lightest charm-strange meson, D_s, is $\Delta M = 350$ MeV$/c^2$. This is less than the kaon mass, thus kinematically forbidding the decay $D_s(2317) \to D_{u,d} + K$. The possible resonance at 2460 MeV$/c^2$ also has such a mass difference when taken with the lighter D^* state. The interpretation of these states as bound $D^{(*)} K$ molecules just below the $D^{(*)} K$ threshold is particularly interesting in the light of the recent discovery of a narrow resonance in the decay $J/\psi \to \gamma p \bar{p}$ [8] which has been interpreted as a zero baryon number, "deuteron-like singlet $^1S_0$" bound state of p and $\bar{p}$ [9].

In the heavy quark theory, the ground state heavy meson involving a heavy and a light quark has the light degrees of freedom in a spin-parity state $j_q^P = \frac{1}{2}^-$, corresponding to the usual pseudoscalar-vector meson doublet with $J^P = (0^-, 1^-)$. The first excited states involves a p-wave excitation, in which the light degrees of freedom have $j_q^P = \frac{1}{2}^+$ or $\frac{3}{2}^+$. This leads

E-mail address: odonnell@physics.utoronto.ca (P.J. O'Donnell).

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to two heavy doublets, the first giving $J^P = (0^+, 1^+)$ and the latter a heavy doublet with $J^P = (1^+, 2^+)$. Heavy quark symmetry rules out any pseudoscalar coupling of this doublet to the ground state at lowest order in the chiral expansion [10] and so these states are expected to be narrow. Recent Belle analysis of $B^- \to D^{(+*)}\pi^+\pi^-$ decays [11] indicate the presence of the 1^+ state in this multiplet at a mass of $M_{D_1^0} = (2421.4 \pm 2.0 \pm 0.4 \pm 0.8)\ \mathrm{MeV}/c^2$ with a width of $\Gamma_{D_1^0} = (23.7 \pm 2.7 \pm 0.2 \pm 4.0)$ MeV. The other state in the doublet (2^+) is also found with a mass of $M_{D_2^0} = (2461.6 \pm 2.1 \pm 0.5 \pm 3.3)\ \mathrm{MeV}/c^2$ with a width of $\Gamma_{D_1^0} = (45.6 \pm 4.4 \pm 6.5 \pm 1.6)$ MeV. In the D_s system the counterpart states to these are naively expected to be a 100 MeV heavier because of the strange quark mass and so these states can probably be identified with $D_{s1}(2536)$ and $D_{sJ}(2573)$ [12]. This is in line with the experimental observations that in the ground state the D_s mesons are about a 100 MeV heavier than their nonstrange counterparts.

The other excited doublet has $J^P = (0^+, 1^+)$. These states are expected to decay rapidly through s-wave pion emission in the $D_{u,d}$ system and by kaon emission in the D_s system and have large widths [13]. Observation of the 1^+ state in the D system was reported by CLEO [14] some time ago. The recent Belle analysis of $B^- \to D^{(+*)}\pi^+\pi^-$ decays [11] also find evidence for the states in this doublet at $M_{D_0^{*0}}(0^+) = (2308 \pm 17 \pm 15 \pm 20)\ \mathrm{MeV}/c^2$ with a width of $\Gamma_{D_1^0} = (276 \pm 21 \pm 18 \pm 60)$ MeV. The other state in the doublet is also found with a mass of $M_{D_1^{*0}}(1^+) = (2427 \pm 26 \pm 20 \pm 15)\ \mathrm{MeV}/c^2$ with a width of $\Gamma_{D_1^0} = (384^{+107}_{-75} \pm 24 \pm 70)$ MeV. Note that these states are broad as expected from theory. Naively then, we should expect the D_s counterparts of these states at $M_{D_s}(0^+) \approx 2408$ and $M_{D_s}(1^+) \approx 2527$. These numbers are consistent with quark model estimates [15] and we expect these states to be broad. The recently observed D_s resonances have masses below these expectations and are very narrow, decaying through isospin violating transitions to $D_s^{(*)}\pi$ final states. This has generated speculations that these states may not be p-wave excited states but rather something exotic like $D^{(*)}K$ molecules.

While the spectroscopy of these newly discovered states can provide clues to their structure, decays involving these states can yield further clues to their exact nature. We first look at nonleptonic B decays involving the p-wave D_s resonant states which we will denote by D_{s0}, corresponding to the p-wave, $j_q = \frac{1}{2}$, 0^+ state, and D_{s1}^* corresponding to the p-wave, $j_q = \frac{1}{2}$, 1^+ state. In B factories that do not produce the B_s mesons the D_s p-wave states cannot be directly produced via the weak current involving the b quark but they can only be produced through the $\bar{s}c$ current in the weak decay effective Hamiltonian. It was suggested in Refs. [16,17] that these theoretically expected broad states may be discovered through the three body decays $B \to D^{(*)}D^{(*)}K$ decays, where $D^{(*)}$ refer to D or D^*, if these states are above the $D^{(*)}K$ threshold. These three body decays can also be used to measure both $\sin 2\beta$ and $\cos 2\beta$ [16,18,19]. In hadron B factories the D_s resonant states can be produced directly from the weak decay of the b quark in the B_s meson.

In this Letter we concentrate on nonleptonic decays of the type $B \to D^{(*)}D_s(2317)$ and $B \to D^{(*)}D_s(2460)$, which are accessible at current B factories, and we also study nonleptonic decays of the types $B_s \to D_s(2317)M$ and $B_s \to D_s(2460)M$, where M is the meson formed by the emitted W. These latter decays can be studied at hadron B factories. Our purpose here is to explore what additional information about the structure and the properties of the new D_s states can be obtained from these nonleptonic decays.

2. Nonleptonic decay

Let us first assume that we can identify the newly discovered states $D_s(2317)$ with D_{s0} and $D_s(2460)$ with D_{s1}^*. In the Standard Model (SM) the amplitudes for $B \to D^{(*)}D_{s0}(D_{s1}^*)$, are generated by the following effective Hamiltonian [20]:

$$H_{\mathrm{eff}}^q = \frac{G_F}{\sqrt{2}}\Bigg[V_{fb}V_{fq}^*\big(c_1 O_{1f}^q + c_2 O_{2f}^q\big) - \sum_{i=3}^{10}\big(V_{ub}V_{uq}^* c_i^u + V_{cb}V_{cq}^* c_i^c + V_{tb}V_{tq}^* c_i^t\big)O_i^q\Bigg] + \mathrm{H.C.}, \qquad (1)$$

where the superscript u, c, t indicates the internal quark, f can be u or c quark, q can be either a d or a s quark depending on whether the decay is a $\Delta S = 0$ or $\Delta S = -1$ process. The operators O_i^q are defined as [21]

$$
\begin{aligned}
O_{1f}^q &= \bar{q}_\alpha \gamma_\mu L f_\beta \bar{f}_\beta \gamma^\mu L b_\alpha, \\
O_{2f}^q &= \bar{q} \gamma_\mu L f \bar{f} \gamma^\mu L b, \\
O_{3,5}^q &= \bar{q} \gamma_\mu L b \bar{q}' \gamma_\mu L(R) q', \\
O_{4,6}^q &= \bar{q}_\alpha \gamma_\mu L b_\beta \bar{q}'_\beta \gamma_\mu L(R) q'_\alpha, \\
O_{7,9}^q &= \frac{3}{2} \bar{q} \gamma_\mu L b e_{q'} \bar{q}' \gamma^\mu R(L) q', \\
O_{8,10}^q &= \frac{3}{2} \bar{q}_\alpha \gamma_\mu L b_\beta e_{q'} \bar{q}'_\beta \gamma_\mu R(L) q'_\alpha,
\end{aligned} \tag{2}
$$

where $R(L) = 1 \pm \gamma_5$, and q' is summed over all flavors except t. $O_{1f,2f}$ are the current–current operators that represent tree level processes. $O_{3–6}$ are the strong gluon induced penguin operators, and operators $O_{7–10}$ are due to γ and Z exchange (electroweak penguins), and "box" diagrams at loop level. The values of the Wilson coefficients can be found in Ref. [20].

In the factorization assumption the amplitude for $B \to D^{(*)} D_{s0}(D_{s1}^*)$, can now be written as

$$M = M_1 + M_2, \tag{3}$$

where

$$
\begin{aligned}
M_1 &= \frac{G_F}{\sqrt{2}} X_1 \langle D_{s0}(D_{s1}^*) | \bar{s} \gamma_\mu (1 - \gamma^5) c | 0 \rangle \\
&\quad \times \langle D^{(*)} | \bar{c} \gamma^\mu (1 - \gamma^5) b | B \rangle,
\end{aligned} \tag{4}
$$

$$
\begin{aligned}
M_2 &= \frac{G_F}{\sqrt{2}} X_2 \langle D_{s0}(D_{s1}^*) | \bar{s} (1 + \gamma^5) c | 0 \rangle \\
&\quad \times \langle D^{(*)} | \bar{c} (1 - \gamma^5) b | B \rangle,
\end{aligned} \tag{5}
$$

where

$$
\begin{aligned}
X_1 &= V_c \left(\frac{c_1}{N_c} + c_2 \right) + \frac{B_3}{N_c} + B_4 + \frac{B_9}{N_c} + B_{10}, \\
X_2 &= -2 \left(\frac{1}{N_c} B_5 + B_6 + \frac{1}{N_c} B_7 + B_8 \right).
\end{aligned} \tag{6}
$$

We have defined

$$B_i = - \sum_{q=u,c,t} c_i^q V_q \tag{7}$$

with

$$V_q = V_{qs}^* V_{qb}. \tag{8}$$

In the above equations N_c represents the number of colors. To simplify matters we neglect the small penguin contributions and so as a first approximation we will neglect M_2. The currents involving the heavy b and c quarks, $J_D^\mu = \langle D | \bar{c} \gamma^\mu (1 - \gamma_5) b | B(p) \rangle$ and $J_{D^*}^\mu = \langle D^*(\epsilon_1) | \bar{c} \gamma^\mu (1 - \gamma_5) b | B(p) \rangle$ can be expressed in terms of form factors [22]. In the heavy quark limit the various form factors are related to a universal Isgur–Wise function $\xi(v \cdot v_1)$ where v and v_1 are the four velocities of the B and the $D^{(*)}$ mesons. One can therefore write,

$$J_D^\mu = \sqrt{m_B} \sqrt{m_D} \xi(v \cdot v_1) \left[v^\mu + v_1^\mu \right] \tag{9}$$

and

$$
\begin{aligned}
J_{D^*}^\mu &= \sqrt{m} \sqrt{m_1} \xi(v \cdot v_1) \\
&\quad \times \left[-i \varepsilon^{\mu\nu\alpha\beta} \epsilon_{1\nu}^* v_\alpha v_{1\beta} + v_1^\mu \epsilon_1^* \cdot v \right. \\
&\quad \left. - \epsilon_1^{*\mu} (v \cdot v_1 + 1) \right].
\end{aligned} \tag{10}
$$

The matrix elements $\langle D_{s0} | \bar{s} \gamma_\mu (1 - \gamma^5) c | 0 \rangle$ and $\langle D_{s1}^* | \bar{s} \gamma_\mu (1 - \gamma^5) c | 0 \rangle$ are written in terms of the decay constants that are defined as

$$
\begin{aligned}
&\langle D_{s0}(P) | \bar{s} \gamma_\mu (1 - \gamma^5) c | 0 \rangle = i f_{D_{s0}} P_\mu, \\
&\langle D_{s1}^*(P, \varepsilon_2) | \bar{s} \gamma_\mu (1 - \gamma^5) c | 0 \rangle = M_{D_{s1}^*} f_{D_{s1}^*} \varepsilon_{2\mu}^*.
\end{aligned} \tag{11}
$$

We can now define the following ratios

$$
\begin{aligned}
R_{D0} &= \frac{\mathrm{BR}[B \to D D_{s0}]}{\mathrm{BR}[B \to D D_s]}, \\
R_{D^*0} &= \frac{\mathrm{BR}[B \to D^* D_{s0}]}{\mathrm{BR}[B \to D^* D_s]}, \\
R_{D1} &= \frac{\mathrm{BR}[B \to D D_{s1}^*]}{\mathrm{BR}[B \to D D_s^*]}, \\
R_{D^*1} &= \frac{\mathrm{BR}[B \to D^* D_{s1}^*]}{\mathrm{BR}[B \to D^* D_s^*]}.
\end{aligned} \tag{12}
$$

Let us focus on the ratio R_{D0} which within factorization and the heavy quark limit can be written as

$$R_{D0} = \left| \frac{f_{D_{s0}}}{f_{D_s}} \right|^2, \tag{13}$$

where we have neglected phase space (and other) effects that are subleading in the heavy quark expansion. Similarly we have

$$R_{D1} = \left| \frac{f_{D_{s1}^*}}{f_{D_s^*}} \right|^2. \tag{14}$$

Now in the heavy quark limit $f_{D_{s0}} = f_{D^*_{s1}}$ and $f_{D_s} = f_{D^*_s}$ and so one would predict $R_{D0} \approx R_{D1}$. There have been various estimates of the decay constant $f_{D_{s0}}$ in quark models [23] and in QCD sum rule calculations (see Refs. [17,19] and references therein); these typically find the p-wave, $j_q = \frac{1}{2}$ states to have the similar decay constants as the ground state mesons. We therefore expect $f_{D_{s0}} \sim f_{D_s}$ giving in addition to the heavy quark predictions

$$R_{D0} \approx R_{D1} \approx 1. \tag{15}$$

Experimentally Belle measures [11]

$$\begin{aligned} &\mathrm{BR}[B \to DD_s(2317)]\,\mathrm{BR}\big[D_s(2317) \to D_s\pi^0\big] \\ &\quad = \big(9.9^{+2.8}_{-2.5} \pm 3.0\big) \times 10^{-4}. \end{aligned} \tag{16}$$

The dominant decay of the $D_s(2317)$ is expected to be through the $D_s\pi$ mode [24,25] and so

$$\mathrm{BR}[D \to DD_s(2317)] \approx 10^{-3}. \tag{17}$$

Now using the measured branching ratio [12]

$$\begin{aligned} &\mathrm{BR}\big[B^+ \to \overline{D}^0 D_s^+\big] = (1.3 \pm 0.4) \times 10^{-2}, \\ &\mathrm{BR}\big[\overline{B}_d \to D^- D_s^+\big] = (8 \pm 3) \times 10^{-3} \end{aligned} \tag{18}$$

one obtains a combined branching ratio

$$\mathrm{BR}[B \to DD_s] \approx 10^{-2}. \tag{19}$$

This leads to $R_{D0} \approx \frac{1}{10}$ (or, $f_{D_{s0}} \sim \frac{1}{3} f_{D_s}$) which is a factor 10 smaller then theoretical expectations. There are a few possible explanations that can be put forward to explain this discrepancy between experiment and theoretical expectation and we will consider them now.

It is possible that the estimate of the decay constants of the p-wave, $j_q = \frac{1}{2}$ states in the various models are incorrect just like the mass predictions of these states are incorrect. This would require a major revision of model calculations that predict the properties of these states. From the experimental data we have seen that $f_{D_{s0}} \sim \frac{1}{3} f_{D_s}$ which gives, using $f_{D_{s0}} = f_{D^*_{s1}}$,

$$R_{D1} \approx \frac{1}{10}. \tag{20}$$

To check this we note that experimentally Belle measures [11]

$$\begin{aligned} &\mathrm{BR}[B \to DD_s(2460)]\,\mathrm{BR}\big[D_s(2460)] \to D_s\pi^0\big] \\ &\quad = \big(25.8^{+7.0}_{-6.0} \pm 7.7\big) \times 10^{-4}, \\ &\mathrm{BR}[B \to DD_s(2460)]\,\mathrm{BR}\big[D_s(2460)] \to D_s\gamma\big] \\ &\quad = \big(5.3^{+1.4}_{-1.3} \pm 1.6\big) \times 10^{-4}. \end{aligned} \tag{21}$$

Taking the central values we find

$$\mathrm{BR}\big[B^+ \to \overline{D}^0 D_s(2460)\big] \leqslant 31.1 \times 10^{-4}. \tag{22}$$

Using the measured branching ratio [12]

$$\begin{aligned} &\mathrm{BR}\big[B^+ \to \overline{D}^0 D_s^{+*}\big] = (9 \pm 4) \times 10^{-3}, \\ &\mathrm{BR}\big[\overline{B}_d \to D^- D_s^{+*}\big] = (1.0 \pm 0.5) \times 10^{-2} \end{aligned} \tag{23}$$

one can obtain, using the measured central values

$$\mathrm{BR}\big[B \to DD_s^*\big] \approx 10^{-2}. \tag{24}$$

This then leads to $R_{D1} \approx \frac{1}{3}$ which is in disagreement with Eqs. (15) and (20).

One might argue that factorization is not applicable to $B \to D^{(*)}D^{(*)}$ decays. However recent analysis in Ref. [26] find that factorization works well for these decays. Moreover, the quantities in Eq. (12) are ratios of nonleptonic decay amplitudes and so nonfactorizable effects may cancel. So what one really requires is significantly different nonfactorizable corrections between decays with the p-wave states in the final state and decays with the ground state mesons in the final state. It is possible that the discrepancies between experiments and theory may arise from a combination of incorrect model prediction of p-wave state properties and nonfactorizable effects.

3. Nonleptonic decays involving B_s decays

Another test of the nature of the newly discovered D_s states that does not rely on factorization or heavy quark symmetry involves the B_s mesons. As we indicated earlier, with the B_s meson, the p-wave D_s states can be produced via the weak current involving the b quark. We can now consider decays of the type $B_s \to D_s(2317)(D_s(2460))M$ where $M = \pi, \rho, K$, etc. With the identification of $D_s(2317)(D_s(2460))$ as the p-wave states these decays are the same as $B_s \to D_{s0}(D^*_{s1})M$.

One can now construct the ratios

$$T_{D_s(2317)} = \frac{\mathrm{BR}[B_s \to D_s(2317)M]}{\mathrm{BR}[B_d \to D_{d0}M]},$$
$$T_{D_s(2460)} = \frac{\mathrm{BR}[B_s \to D_s(2460)M]}{\mathrm{BR}[B_d \to D^*_{d1}M]},$$
$$T_{D_s} = \frac{\mathrm{BR}[B_s \to D_s M]}{\mathrm{BR}[B_d \to D_d M]},$$
$$T_{D_s^*} = \frac{\mathrm{BR}[B_s \to D_s^* M]}{\mathrm{BR}[B_d \to D_d^* M]}. \tag{25}$$

Now in the $SU(3)$ limit all the ratios are unity. Moreover, the ratio of ratios $r_0 = T_{D_s(2317)}/T_{D_s}$ and $r_1 = T_{D_s(2460)}/T_{D_s^*}$ are expected to have smaller flavour symmetry violations and hence smaller deviations from unity, as $SU(3)$ breaking effects in the ratios may cancel [27]. Hence any large deviation of $T_{D_s(2317)}$ and $T_{D_s(2460)}$ from unity would be inconsistent with the $j_q = \frac{1}{2}$ p-wave interpretation of the new D_s states. Note that the further assumption of factorization leads to $T_{D_s(2317)} \approx T_{D_s(2460)}$ and $T_{D_s} \approx T_{D_s^*}$ in the heavy quark limit.

As indicated earlier, among various other suggestions for the nature of the new D_s states is the idea that these states may be $D^{(*)}K$ molecules. There are no serious models of such meson molecules that one can use to calculate nonleptonic decays involving these states. Here we will attempt a rough qualitative estimate of nonleptonic decay rates assuming that the $D_s(2317)$ and $D_s(2460)$ states are really a DK molecule and a D^*K molecule, respectively. Consider the nonleptonic decay $B \to DD_s(2317)$. We assume that the decay proceeds through two stages: the first stage is the decay $B \to DDK$, followed by the state $D(p_2)K(p_K)$ forming the molecule $D_s(2317)$ with the probability given by $f(p_2, p_K)$ so that

$$d\Gamma\big(B \to DD_s(2317)\big) = \frac{1}{(2\pi)^3}\frac{1}{8M_B}\big|A\big(B \to D(p_1)D(p_2)K(p_K)\big)\big|^2 \times f(p_2, p_K)dE_K dE_2. \tag{26}$$

Without a model for $f(p_2, p_K)$ we cannot make predictions but nonetheless it is useful to define the average probability function $\bar{f}$ as

$$\bar{f} = \left[\int \big|A\big(B \to D(p_1)D(p_2)K(p_K)\big)\big|^2 \times f(p_2, p_K)\, dE_K\, dE_2\right] \times \left[\int \big|A\big(B \to D(p_1)D(p_2)K(p_K)\big)\big|^2 \times dE_K\, dE_2\right]^{-1}. \tag{27}$$

Hence we have

$$\mathrm{BR}\big(B^+ \to \overline{D}^{(0*)}D_s(2317)^+\big) = \mathrm{BR}\big(B^+ \to \overline{D}^{(0*)}D^+K^0\big) \times \bar{f},$$
$$\mathrm{BR}\big(B^0 \to D^{(-*)}D_s(2317)^+\big) = \mathrm{BR}\big(B^0 \to D^{(-*)}D^0K^+\big) \times \bar{f}. \tag{28}$$

We can define a similar function f^* and the average $\bar{f}^*$ for nonleptonic decays involving the $D_s(2460)$ and so

$$\mathrm{BR}\big(B^+ \to \overline{D}^{(0*)}D_s(2460)^+\big) = \mathrm{BR}\big(B^+ \to \overline{D}^{(0*)}D^{+*}K^0\big) \times \bar{f}^*,$$
$$\mathrm{BR}\big(B^0 \to D^{(-*)}D_s(2460)^+\big) = \mathrm{BR}\big(B^0 \to D^{(-*)}D^{0*}K^+\big) \times \bar{f}^*. \tag{29}$$

We can consider the ratios

$$Z^+_{\mathrm{res}} = \frac{\mathrm{BR}(B^+ \to \overline{D}^0 D_s(2460)^+)}{\mathrm{BR}(B^+ \to \overline{D}^0 D_s(2317)^+)},$$
$$Z^{+*}_{\mathrm{res}} = \frac{\mathrm{BR}(B^+ \to \overline{D}^{0*} D_s(2460)^+)}{\mathrm{BR}(B^+ \to \overline{D}^{0*} D_s(2317)^+)},$$
$$Z^0_{\mathrm{res}} = \frac{\mathrm{BR}(B^0 \to D^- D_s(2460)^+)}{\mathrm{BR}(B^0 \to \overline{D}^- D_s(2317)^+)},$$
$$Z^{0*}_{\mathrm{res}} = \frac{\mathrm{BR}(B^0 \to D^{-*} D_s(2460)^+)}{\mathrm{BR}(B^0 \to \overline{D}^{-*} D_s(2317)^+)},$$
$$Z^+_{3\text{-body}} = \frac{\mathrm{BR}(B^+ \to \overline{D}^0 D^{+*}K^0)}{\mathrm{BR}(B^+ \to \overline{D}^0 D^+K^0)},$$
$$Z^{+*}_{3\text{-body}} = \frac{\mathrm{BR}(B^+ \to \overline{D}^{0*} D^{+*}K^0)}{\mathrm{BR}(B^+ \to \overline{D}^{0*} D^+K^0)},$$
$$Z^0_{3\text{-body}} = \frac{\mathrm{BR}(B^0 \to D^- D^{0*}K^+)}{\mathrm{BR}(B^0 \to \overline{D}^- D^0K^+)},$$
$$Z^{0*}_{3\text{-body}} = \frac{\mathrm{BR}(B^0 \to D^{-*} D^{0*}K^+)}{\mathrm{BR}(B^0 \to \overline{D}^{-*} D^0K^+)}, \tag{30}$$

which are related as

$$Z^{+}_{\text{res}} = Z^{+}_{\text{3-body}} \frac{\bar{f}^{*}}{\bar{f}},$$
$$Z^{+*}_{\text{res}} = Z^{+*}_{\text{3-body}} \frac{\bar{f}^{*}}{\bar{f}},$$
$$Z^{0}_{\text{res}} = Z^{0}_{\text{3-body}} \frac{\bar{f}^{*}}{\bar{f}},$$
$$Z^{0*}_{\text{res}} = Z^{0*}_{\text{3-body}} \frac{\bar{f}^{*}}{\bar{f}}. \tag{31}$$

Using the measured three-body branching ratios [28]

$$\begin{aligned}
&\mathrm{BR}\big(B^{+} \to \overline{D}^{0} D^{+} K^{0}\big) \\
&\quad = (0.18 \pm 0.07 \pm 0.04) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{0} \to D^{-} D^{0} K^{+}\big) \\
&\quad = (0.17 \pm 0.03 \pm 0.03) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{+} \to \overline{D}^{0*} D^{+} K^{0}\big) \\
&\quad = \big(0.41^{+0.15}_{-0.14} \pm 0.08\big) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{0} \to D^{-*} D^{0} K^{+}\big) \\
&\quad = \big(0.31^{+0.04}_{-0.03} \pm 0.04\big) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{+} \to \overline{D}^{0} D^{+*} K^{0}\big) \\
&\quad = \big(0.52^{+0.10}_{-0.09} \pm 0.07\big) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{0} \to D^{-} D^{*0} K^{+}\big) \\
&\quad = (0.46 \pm 0.07 \pm 0.07) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{+} \to \overline{D}^{0*} D^{+*} K^{0}\big) \\
&\quad = \big(0.78^{+0.26}_{-0.21} \pm 0.14\big) \times 10^{-2}, \\
&\mathrm{BR}\big(B^{0} \to D^{-*} D^{0*} K^{+}\big) \\
&\quad = \big(1.18 \pm 0.10 \pm 0.17\big) \times 10^{-2}
\end{aligned} \tag{32}$$

which are proportional to either $1 - \bar{f}$ or $1 - \bar{f}^{*}$ and assuming $\bar{f}^{*} \approx \bar{f}^{*}$ allows one to obtain, with the central values of the measurements,

$$Z^{+}_{\text{3-body}} = 2.89 \tag{33}$$

which can be compared to $Z^{+}_{\text{res}} = 3.14$ from Eqs. (17) and (22). If fact the prediction $Z^{+}_{\text{res}} \sim 3$, $Z^{+*}_{\text{res}} \sim 3$, $Z^{0}_{\text{res}} \sim 3$ and $Z^{0*}_{\text{res}} \sim 3$ are consistent within the errors for the three-body branching ratios in $Z^{+}_{\text{3-body}}$, $Z^{+*}_{\text{3-body}}$, $Z^{0}_{\text{3-body}}$ and $Z^{0*}_{\text{3-body}}$. We also obtain $\bar{f} \approx \bar{f}^{*} \approx 0.3$, from Eqs. (17) and (32) which indicates that a sizeable fraction of the $D^{(*)}K$ state form molecules.

Finally we can extend this model also to the case where the D_s resonance is produced via the weak current containing the b quark in B_s decays. Consider the decays $B_s \to D_s(2317)M$ where M is the emitted meson. The form factor for $B \to D_s(2317)$ transition can then be related to $B \to DK$ transition. In other words, we can write

$$\begin{aligned}
&\mathrm{BR}[B_s \to D_s(2317)M] \\
&\quad = \mathrm{BR}[B_s \to DKM]\bar{f}',
\end{aligned} \tag{34}$$

where

$$\begin{aligned}
\bar{f}' = {} & \bigg[\int \big|A\big(B \to D(p_2)K(p_K)M(p_1)\big)\big|^2 \\
& \quad \times f(p_2, p_K)\, dE_K\, dE_2\bigg] \\
& \quad \times \bigg[\int \big|A\big(B \to D(p_2)K(p_K)M(p_1)\big)\big|^2 \\
& \quad \times dE_K\, dE_2\bigg]^{-1}.
\end{aligned} \tag{35}$$

We can similarly define $\bar{f}^{*\prime}$ as

$$\begin{aligned}
&\mathrm{BR}[B_s \to D_s(2460)M] \\
&\quad = \mathrm{BR}[B_s \to D^{*}KM]\bar{f}^{*\prime}.
\end{aligned} \tag{36}$$

Note that the ratios $T_{D_s(2317)}$ and $T_{D_s(2460)}$ (Eq. (25)) in the molecular model are no longer equal to unity in the $SU(3)$ limit since that depended on the identification of these states as p-wave states. Therefore the measurement of these ratios can provide useful information on the nature of the $D_s(2317)$ and the $D_s(2460)$ states.

4. Summary and conclusions

In summary, in this work, we have considered the nonleptonic B decays $B \to D^{(*)}D_s(2317)(D_s(2460))$, involving the newly discovered $D_s(2317)$ and the $D_s(2460)$ states. We have discussed the implication of the measured nonleptonic decays for the properties and the nature of these states. If these states are the p-wave multiplet with the light degrees of freedom in the $j_q = \frac{1}{2}$ state, then we find that experiments indicate disagreement with model calculation of their properties and/or breakdown of the factorization assumption. We have suggested further tests involving nonleptonic B_s meson decays, that do not assume factorization but assumes $SU(3)$ flavour symmetry, that can further shed light on the true nature of these newly discovered states. Finally, we have also proposed a model to calculate the two body nonleptonic decays $B \to D^{(*)} D_s(2317)(D_s(2460))$, assuming that the $D_s(2317)$ and $D_s(2460)$ are DK and $D^* K$ molecules. The model relates these two body nonleptonic decays to the three-body B decays of the type $B \to D^{(*)} D^{(*)} K$.

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