Search of $D^{*}_{sJ}$ mesons in $B$ meson decays

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We propose that the search of the $B \to D^{*}_{sJ}M$ decays, $M = D$, $\pi$, and $K$, can discriminate the different theoretical postulations for the nature of the recently observed $D^{*}_{sJ}$ mesons. The ratio of the branching ratios $B(B \to D^{*}_{sJ}M)/B(B \to D^{*}_{sJ}M) \approx 0.1$ supports that the $D^{*}_{sJ}$ mesons are quark-antiquark (multi-quark) bound states. The Belle measurement of the $B \to D^{*}_{sJ}D$ branching ratios seems to indicate an unconventional picture.

BaBar collaboration observed a narrow state with $J^{P} = 0^{+}$, denoted by $D^{*}_{sJ}(2317)$, from the $D^{+}_s\pi^0$ invariant mass distribution [1], whose mass was determined to be 2317.6 ± 1.3 MeV, and whose width is consistent with the experimental resolution, being less than 10 MeV. This observation has been confirmed by CLEO, and another new state $D^{*}_{sJ}(2463)$ with $J^{P} = 1^{+}$ was found in the $D^{*}_{sJ}\pi^0$ channel with the mass splitting 351.6 ± 1.7 ± 1.0 MeV from the ordinary vector meson $D^{*}_{s}$ and with the width being less than 7 MeV [2]. It is then an urgent subject to understand the nature of these newly observed states, and many theoretical speculations have appeared in the literature. In this paper we shall propose an experimental strategy, which can make a substantial contribution to this subject.

The measured masses and widths of the new states do not match the predictions from typical potential models. For example, the mass and width of the scalar $D^{*}_{sJ}(2317)$ meson were expected to be around 2.48 GeV and 160 MeV [3], respectively. It has been shown that the masses and widths of the $D_s$ system can not be explained simultaneously in the potential model [4]. To resolve the discrepancy, either the theoretical models need to be modified, or the new mesons are unconventional bound states. For the former, a unitarized quark model has been adopted, which includes the coupling of the scalar meson to an OZI-allowed two-meson channel [5]. A low-mass scalar $D_s$ meson as a quark-antiquark state could be obtained. For the latter, the $D^{*}_{sJ}(2317)$ meson has been interpreted as a $DK$ molecule [6], a $D\pi$ molecule [7], a four-quark state [8], and a mixing of the conventional state and the four-quark state [9]. However, it was argued that the charm-strange, and even bottom-strange, four-quark states could not be bound [10]. A lattice study in the static limit, which predicts a larger mass for the scalar $D_s$ meson as a quark-antiquark state, supports the multi-quark postulation [11]. The larger scalar mass in the quark-antiquark picture has been confirmed by a sum-rule analysis [12].

Considering the above series of claims and counter claims, it is worthwhile to look for alternative theoretical and experimental viewpoints, which may help to clarify the controversy. For example, it has been claimed that the existence of a new $I = 0$ “$D\bar{D}$ bound state” with a mass less than 3660 MeV would support the four-quark picture [13]. Whether the $D^{*}_{sJ}$ meson radiative transition is consistent with the branching ratios of the conventional $D^{*}_{s0}$ and $D^{*}_{s1}$ mesons also serves the purpose [14]. In this work we propose that the search of the $B \to D^{*}_{sJ}M$ decays, $M = D$, $\pi$ and $K$, can discriminate the different theoretical postulations for the $D^{*}_{sJ}$ content. In the quark-antiquark picture the $B \to D^{*}_{sJ}M$ branching ratios are expected to be of the same order of magnitude as the $B \to D^{*}_{sJ}M$ ones, since the $D^{*}_{sJ}$ meson decay constants should be close to those of the conventional $D^{*}_{sJ}$ mesons as required by chiral symmetry [14]. We shall assume that the chiral symmetry is a good symmetry in our analysis. In the unconventional picture the corresponding decay amplitudes involve additional hard scattering the four valence quarks of the $D^{*}_{sJ}$ mesons participate. The branching ratios are then at least suppressed by the coupling constant and by inverse powers of heavy meson masses, such that they are smaller than the $B \to D^{*}_{sJ}M$ ones by a factor of 10.

The $B_d(P_1) \to D^{*+}_{sJ}(P_2)\pi^- (P_3)$ decay occurs through the effective Hamiltonian,

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{ub} V_{cb} \left[ C_1(\mu) O_1(\mu) + C_2(\mu) O_2(\mu) \right],$$

(1)

with the four-fermion operators $O_1 = (\bar{s}_i c_j)(\bar{u}_j b_i)$ and $O_2 = (\bar{s}_i c_j)(\bar{u}_j b_i)$, $\bar{q}_i j_q \equiv q_j \gamma_\mu (1 - \gamma_5) q_i$ and $i$ and $j$ being the color indices, the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements $V_{ij}$, and the Wilson coefficients $C_{1,2}(\mu)$. We choose a frame, in which the $B$ meson is at rest and the pion momentum $P_3$ is in the minus direction in the light-cone coordinates. The two-body decay rate is expressed as $\Gamma = |A|^2/(16\pi m_B)$, $m_B$ being the $B$ meson mass and $A$ the decay amplitude.

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In the quark-antiquark picture the above decay contains a color-allowed amplitude, which is written, in the factorization assumption (FA), as
\[ A = i \frac{G_F}{\sqrt{2}} V_{ub} V_{cs} (m_B^2 - m_{sJ}^2) f_{D_{sJ}} F_0^{B\pi}(m_{sJ}^2) a_1, \]
with the $D_{sJ}$ meson decay constant $f_{D_{sJ}}$, the $D_{sJ}$ meson (pion) mass $m_{D_{sJ}}$, and $a_1 = C_2 + C_4/N_c$, $N_c = 3$ being the number of colors. Employing the inputs $G_F = 1.16639 \times 10^{-5}$ GeV$^{-2}$, $|V_{ub}| = 0.003$, $|V_{cs}| = 0.976$, $m_B = 5.28$ GeV, $\tau_{B^0} = 1.542 \times 10^{-12}$s, $m_{D_{sJ}} = 2.32$ GeV, and $f_{D_{sJ}} = 0.24$ GeV, $F_0^{B\pi}(m_{sJ}^2) = 0.33$ from the light-cone-sum-rule results, and $a_1 = 1.1$ for a wide range of the renormalization scale $\mu$, we have the branching ratio,
\[ B(B_d \rightarrow D_{sJ}^{*+}\pi^-) = 3.0 \times 10^{-5}, \]
close to the Belle and BaBar measurements, $B(B_d \rightarrow D_{sJ}^{*+}\pi^-) = (2.4^{+1.0}_{-0.8} \pm 0.7, 4.6^{+1.2}_{-1.1} \pm 1.3) \times 10^{-5}$. Because of $m_{D_{sJ}(2317)} \approx m_{D_{sJ}(2463)}$, the result in Eq. 3 holds for both the $D_{sJ}^{*}(2317)$ and $D_{sJ}^{*}(2463)$ mesons.

If the $D_{sJ}$ meson is a four-quark bound state, the $B_d$ meson decays into $D_{sJ}^{*+}\pi^-$ through the diagrams Figs. 1(a)-1(d), in which all its four valence quarks participate hard scattering. An extra hard gluon is then necessary for producing the $u\bar{u}$ quark pair, and more virtual lines appear. For the type of Figs. 1(e) and 1(f), the exchanged gluon, being of collinear origin with the momentum in the plus direction, should be absorbed into the two-parton $D_{sJ}^{*+}$ meson distribution amplitude. That is, Figs. 1(e) and 1(f) contribute to the analysis in the quark-antiquark picture.

There are two color configurations,
\[ \frac{1}{N_c^2} c^h s^b n^c \bar{n}^c, \frac{1}{12} f^{ab_1 c_1} f^{abc_2 c_2} c^b_1 s^c_1 n^b_2 \bar{n}^c_2, \]
where both the $c\bar{s}$ and $n\bar{n} = (u\bar{u} + d\bar{d})/\sqrt{2}$ pairs are in the color-singlet 1 and color-octet 8 states, respectively. The average over colors has been made explicit. The Wilson coefficients associated with each diagram from the 11 and 88 configurations are listed in Table I. It is found that the Wilson coefficient for Fig. 1(a) from 11 is largest. The contributions from Figs. 1(b) and 1(c), besides a pair cancellation, are down by a small ratio $C_1/a_1 \sim -0.2$. As shown later, the amplitude corresponding to Fig. 1(d), where the hard gluon attaches the $b$ quark, is suppressed by a power of $\Lambda_{QCD}/m_{D_{sJ}} \sim 0.1$, though it is not down by a Wilson coefficient. Hence, we can safely drop Figs. 1(b)-1(f), and consider only Fig. 1(a) from the 11 color configuration.
A quantitative analysis of Fig. 1(a) requires the knowledge of the four-parton $D_{sJ}^*$ meson distribution amplitude. Before this information is available, we make a simple estimation also in FA. Insert the Fierz identity,

$$1_{ij}1_{ik} = \frac{1}{8}(1 - \gamma_5)_{ik}(1 - \gamma_5)_{lj} + \frac{1}{8}(1 + \gamma_5)_{ik}(1 + \gamma_5)_{lj} + \frac{1}{8}[\gamma_{\nu}(1 - \gamma_5)]_{ik}[(1 - \gamma_5)\gamma_{\nu}]_{lj} + \frac{1}{8}[\gamma_{\nu}(1 + \gamma_5)]_{ik}[(1 + \gamma_5)\gamma_{\nu}]_{lj},$$

into Fig. 1(a) to factorize the fermion flows. The first term, inserted in the way indicated by the lower dashed line, gives the factorization of the $B \to \pi$ form factor from the full amplitude. The insertion of the third term indicated by the upper dashed line then leads to a nonvanishing hard kernel and to the matrix element $\langle D_{sJ}^* | \bar{c} \gamma_{\mu}(1 - \gamma_5)s\bar{u} \gamma_{\nu}(1 - \gamma_5) | u \rangle = 0$, which defines the $D_{sJ}^*$ meson decay constant. There exists another factorization of fermion flows with the fourth (first) term in Eq. 5 inserted at the lower (upper) dashed line. However, this factorization introduces the matrix element $\langle D_{sJ}^* | \bar{c} \gamma_{\mu}(1 - \gamma_5)s\bar{u}(1 - \gamma_5) | u \rangle$, which is suppressed by a power of $m_{D_{sJ}^*}/m_B$ compared to the previous one.

We derive the $B_d \to D_{sJ}^* \pi^-$ decay amplitude in FA in the four-quark picture,

$$A = \frac{G_F}{\sqrt{2}} V_{ub}^* V_{cs} \langle D_{sJ}^* | \bar{c} \gamma_{\mu}(1 - \gamma_5)s\bar{u} | u \rangle \langle \pi^- | \bar{b} \gamma_{\nu}(1 - \gamma_5)u | B_d \rangle a_1 \hat{H}^\nu,$$

with the hard kernel,

$$\hat{H}^\nu = \frac{g^2}{32\sqrt{2} \Lambda_{QCD}} \frac{C_F}{N_c} tr[(1 - \gamma_5) J_\alpha \gamma_{\beta}(1 - \gamma_5) \gamma_{\nu} \gamma_{\tau} \gamma_5],$$

where $l_u$ and $l_g$ are the momenta carried by the internal $u$ quark and gluon, respectively, and the denominator $\sqrt{2}$ comes from the definition of $n\bar{u}$. To be precise, $\hat{H}^\nu$ should be expressed as a convolution of Eq. 7 with the four-parton distribution amplitude over the momentum fractions of the valence $\bar{s}$, $u$ and $\bar{u}$ quarks. For the purpose of estimation, we regard that these valence quarks carry the fixed momentum fractions of $O(\Lambda_{QCD}/m_{D_{sJ}^*})$ [21]. Therefore, the components of $l_u$ and $l_g$ have the orders of magnitude,

$$l_u \sim l_g \sim \frac{m_B}{\sqrt{2}} \left( \frac{\Lambda_{QCD}}{m_{D_{sJ}^*}}, \frac{1}{2}, \mathbf{0}_T \right),$$

where the valence $\bar{u}$ quark in the pion has been assumed to take half of the pion momentum. The virtual $\bar{b}$ quark momentum in Fig. 1(d) has the components $l_b \sim (m_B/\sqrt{2})(1, 1/2, \mathbf{0}_T)$, such that Fig. 1(d) is power-suppressed by $l_u^2/l_b^2 \sim \Lambda_{QCD}/m_{D_{sJ}^*}$ compared to Fig. 1(a) as stated before.

Next step is to evaluate the matrix element,

$$\langle D_{sJ}^*(P_2) | \bar{c} \gamma_{\mu}(1 - \gamma_5)s\bar{u} \gamma_{\nu}(1 - \gamma_5) | u \rangle = \frac{B}{m_{D_{sJ}^*}} P_{2\mu} P_{2\nu} f_{D_{sJ}^*},$$

which has been parametrized in terms of a dimensional constant $B$. Under the heavy quark symmetry, this matrix element should be close to $\langle D^0 | \bar{c} \gamma_{\mu}(1 - \gamma_5)u\bar{u} \gamma_{\nu}(1 - \gamma_5) | u \rangle$. The equation of motion for the heavy $c$ quark with the momentum $P_c \approx P_2$ and the relation $P_{2\mu}^2 = m_B^2$ lead to $\langle D^0 | \bar{c}(1 - \gamma_5)u\bar{u} \gamma_{\nu}(1 - \gamma_5) | u \rangle = iBP_{2\nu} f_D$ with the $D^0$ meson decay constant $f_D$. The Fierz transformation of the four-quark operator and FA of the matrix element give $\langle D^0 | \bar{c} \gamma_{\nu}(1 - \gamma_5)u\bar{u}(1 - \gamma_5) | u \rangle \approx \langle D^0 | \bar{c} \gamma_{\nu}(1 - \gamma_5)u | u \rangle \langle u | \bar{u}(1 - \gamma_5)u \rangle$. Substituting the definition of $f_D$, we derive

$$B \approx \langle 0 | \bar{u}(1 - \gamma_5)u | 0 \rangle = \langle 0 | \bar{u}u | 0 \rangle \approx -0.24 \text{ GeV}^3,$$

where the standard value of the quark condensate has been adopted.

| configuration | (a) | (b) | (c) | (d) |
|--------------|-----|-----|-----|-----|
| 11           | $a_1/N_c$ | $C_1/N_c$ | $C_1/N_c$ | $a_1/N_c$ |
| 88           | $C_1/N_c$ | $C_1/N_c$ | $C_1/N_c$ | $C_1(1/N_c - 1)$ |
Equation (6) then becomes

\[ A = i \frac{G_F}{\sqrt{2}} V_{ub} V_{cs} (m_B^2 - m_s^2) f_{D_{sj}^*} f_{D_J} B(\pi) a_1 R, \]

with the ratio,

\[ R = \frac{g^2 C_F}{32 \sqrt{2} N_c} \frac{\text{tr}(1 - \gamma_5) I_{u_\gamma\alpha}(1 - \gamma_5) g_2 \gamma^3 I_{u_\gamma\alpha}(1 - \gamma_5) g_2 \gamma^3]}{m_{D_{sj}^*}^2 \alpha_s^2(\mu)} \times \left( \frac{m_{D_{sj}^*}}{m_B} \right)^2 \approx 0.275, \]

for the inputs \( \alpha_s/\pi = 0.2 \) in \( b \to c \) transitions and \( \Lambda_{QCD} \approx 0.3 \text{ GeV}, \) \( l_s^2 \) and \( l_g^2 \) from Eq. (5) have been inserted. It is easy to see that the decay amplitude in the four-quark picture is down by the coupling constant \( \alpha_s \), by the color number \( 1/N_c \), and by the powers \( (m_{D_{sj}^*}/m_B)^2 \). We conclude that the \( B_d \to D_{s,j}^+ \pi^- \) branching ratio in the four-quark picture should be smaller than that in the quark-antiquark one by a suppression factor,

\[ \frac{B(4)(B_d \to D_{s,j}^+ \pi^-)}{B(3)(B_d \to D_{s,j}^+ \pi^-)} \approx R^2 \approx 0.08. \]

If the \( B_d \to D_{s,j}^+ \pi^- \) branching ratios are observed at the \( 10^{-5} \) level as in Eq. (3), the \( D_{s,j}^+ \) meson is likely to be a conventional quark-antiquark state. If it is observed with the \( 10^{-6} \) (around \( 2.4 \times 10^{-6} \)) branching ratio, the four-quark picture is preferred.

There is already a hint from the \( B \to D_{s,j}^* D \) decays, to which our analysis can be generalized straightforwardly simply by substituting the \( B \to D \) form factor for the \( B \to \pi \) form factor. The \( B \to D_{s,j}^* D \) branching ratios have been measured by Belle recently:

\[ B(B^+ \to D_{s,j}^+(2317)\bar{D}^0) \times B(D_{s,j}^+(2317) \to D_s^+ \pi^0) = 8.1^{+3.0}_{-2.7} \times 10^{-4}, \]
\[ B(B^+ \to D_{s,j}^+(2463)\bar{D}^0) \times B(D_{s,j}^+(2463) \to D_s^+ \pi^0) = 11.9^{+6.4}_{-4.9} \times 10^{-4}, \]
\[ B(B^+ \to D_{s,j}^+(2463)\bar{D}^0) \times B(D_{s,j}^+(2463) \to D_s^+ \gamma) = 5.6^{+1.6}_{-1.5} \times 10^{-4}. \]

The first data together with \( B(B^+ \to D_s^+ \bar{D}^0) = 1.3 \pm 0.4 \% \) imply

\[ \frac{B(B^+ \to D_{s,j}^+(2317)\bar{D}^0)}{B(B^+ \to D_s^+ \bar{D}^0)} \approx 0.06, \]

and the four-quark content of the \( D_{s,j}^* \) meson. The latter two data, assuming that the \( D_{s,j}^+(2463) \) decays only through the channels \( D_s^+ \pi^0 \) and \( D_s \gamma \), lead to \( B(B^+ \to D_{s,j}^+(2463)\bar{D}^0) \approx 0.18 \% \) and the ratio,

\[ \frac{B(B^+ \to D_{s,j}^+(2463)\bar{D}^0)}{B(B^+ \to D_s^+ \bar{D}^0)} \approx 0.14, \]

which also gives a similar indication. It is unlikely that the dramatically different branching ratios in Eq. (15) is due to the different decay constants \( f_{D_{sj}^*} \) and \( f_{D_s} \) from the viewpoint of heavy quark symmetry.

At last, we discuss another ideal mode for the purpose, the \( B_d \to D_{s,j}^* K^{(*)} \) decay, which occurs through the operators \( O_1 = (d_i u_j) (c_i b_j) \) and \( O_2 = (d_i u_j) (c_j b_i) \) with the product of the CKM matrix element \( V_{ud} V_{cs}^* \). Since this mode involves only the annihilation topologies, FA does not apply. Hence, we estimate its branching ratio in the quark-antiquark picture using the perturbative QCD (PQCD) approach, in which a transition matrix element is expressed as the convolution of hard kernels of the valence quarks with hadron distribution amplitudes. The derivation of the factorization formulas at leading power in \( 1/m_B \) and leading order in \( \alpha_s \) follow that for the \( B \to D^{(*)} \pi(\rho, \omega) \) decays in (20). We shall present the explicit expressions elsewhere. Adopting the \( D_{s,j}^* \) meson distribution amplitudes the same as those for the \( D^{(*)} \) meson in (20) (the \( B, K \) and \( K^* \) meson distribution amplitudes have been known from the literature), we obtain the branching ratios listed in Table (11). The predictions for the \( B_d \to D_{s,j}^* K^+ \) mode are close the Belle and BaBar measurements (17, 18, 22). The estimation for the \( B_d \to D_{s,j}^* K^{(*)} \) branching ratios in the four-quark picture is similar, and the results are also smaller than those in the quark-antiquark picture by a factor about 0.08. For example, the \( B_d \to D_{s,j}^+(2317)K^+ \) branching ratio is expected to be around \( 4.2 \times 10^{-6} \).
TABLE II: $B_d \rightarrow D_{sJ}^{(*)+} K^{(*)}$ branching ratios (in units of $10^{-5}$) in the quark-antiquark picture from the PQCD approach.

| $B_d \rightarrow D_{sJ}^{(*)+}(2317)K^+$ | $B_d \rightarrow D_{sJ}^{(*)+}(2317)K^{(*)+}$ | $B_d \rightarrow D_{sJ}^{(*)+}(2463)K^+$ |
|---------------------------------|---------------------------------|---------------------------------|
| 5.35                            | 7.79                            | 7.01                            |

Our estimation given above applies to other non-quark-antiquark models for the $D_{sJ}^*$ content, such as a molecule, up to order of magnitude. One of the differences is that the 88 color configuration is excluded, which is not essential anyway.

In summary, a measurement of the $B \rightarrow D_{sJ}^* M$ branching ratios, $M = D, \pi$ and $K$, can provide more information on the nature of the new $D_{sJ}^*$ mesons. If the $D_{sJ}^*$ mesons are multi-quark bound states, they will be more difficult to be produced than the conventional $D_{sJ}^{(*)}$ mesons in exclusive $B$ meson decays: the branching ratios will be one order of magnitude smaller. The suppression is a combined effect of $\alpha_s, 1/N_c$ and $(m_{D_{sJ}}/m_B)^2$, which arise from the additional hard scattering the four valence quarks of the $D_{sJ}^*$ mesons participate. More precise data are necessary for drawing a conclusion, though the recently measured $B \rightarrow D_{sJ}^* D$ branching ratios have indicated an unconventional picture.

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