Looking for D_{sJ}^{*} mesons in B meson decays

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We propose that the search of the $B\to D_{sJ}^{*}M$ decays, $M=D$, π, and K, can discriminate the different theoretical postulations for the nature of the recently observed D_{sJ}^{*} mesons. The ratio of the branching ratios $B(B\to D_{sJ}^{*}M)/B(B\to D_{s}^{(*)}M)\approx 1$ (0.1) supports that the D_{sJ}^{*} mesons are quark-antiquark (multiquark) bound states. The Belle measurement of the $B\to D_{sJ}^{*}D$ branching ratios seems to indicate an unconventional picture.

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The BaBar Collaboration observed a narrow state with $J^{P}=0^{+}$, denoted as $D_{sJ}^{*}(2317)$, 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 $D_{s}^{*+}\pi^{0}$ channel with the mass splitting $351.6\pm 1.7\pm 1.0$ 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 that 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 cannot 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 Okubo-Zweia-Iizukai (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 $D_{s}\pi$ 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 multiquark 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 counterclaims, 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, $M=D$, π, 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 $B\to D_{s}^{(*)}M$ ones, since the D_{sJ}^{*} meson decay constants should be close to those of the conventional $D_{s}^{(*)}$ mesons as required by chiral symmetry [15]. 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 to which 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 $B\to D_{s}^{(*)}M$ 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,

$$H_{\rm eff}=\frac{G_{F}}{\sqrt{2}}V_{ub}V_{cs}^{*}[C_{1}(\mu)\mathcal{O}_{1}(\mu)+C_{2}(\mu)\mathcal{O}_{2}(\mu)], \quad (1)$$

with the four-fermion operators $\mathcal{O}_{1}=(\bar{s}_{i}c_{j})(\bar{u}_{j}b_{i})$ and $\mathcal{O}_{2}=(\bar{s}_{i}c_{i})(\bar{u}_{j}b_{j})$, $(\bar{q}_{i}q_{j})\equiv\bar{q}_{i}\gamma_{\mu}(1-\gamma_{5})q_{j}$ (i and j being the color indices), the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements V's, 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.

In the quark-antiquark picture, the above decay contains a color-allowed amplitude, which is written, in the factorization assumption (FA), as

$$A=i\frac{G_{F}}{\sqrt{2}}V_{ub}^{*}V_{cs}(m_{B}^{2}-m_{\pi}^{2})f_{D_{sJ}^{*}}F_{0}^{B\pi}(m_{D_{sJ}^{*}}^{2})a_{1}, \quad (2)$$

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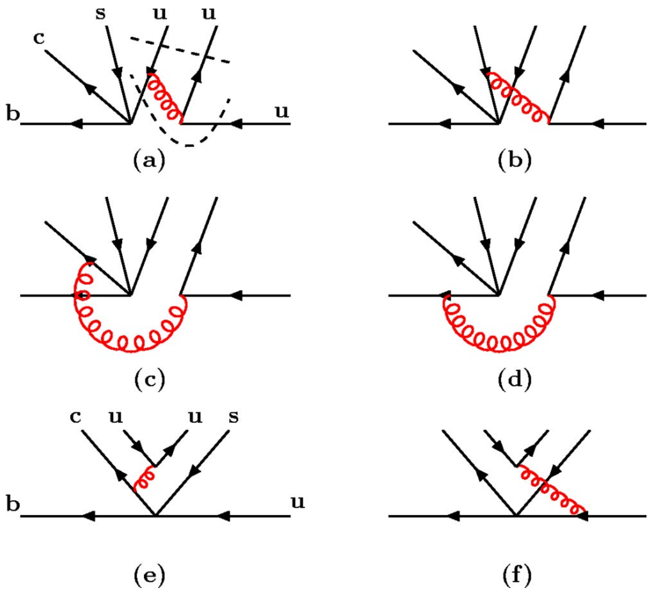

FIG. 1. (a)-(d) Diagrams contributing to the $B_d \to D_{sJ}^{*+}\pi^-$ decay in the four-quark picture for the D_{sJ}^* content. (e) and (f) Diagrams contributing in the quark-antiquark picture.

with the D_{sJ}^* meson decay constant $f_{D_{sJ}^*}$, the D_{sJ}^* meson (pion) mass $m_{D_{sJ}^*}$ (m_π), and $a_1 = C_2 + C_1/N_c$, $N_c = 3$ being the number of colors. Employing the inputs $G_F = 1.166\,39 \times 10^{-5}$ 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, $F_0^{B\pi}(m_{D_{sJ}^*}^2) = 0.33$ from the light-cone-sum-rule results [16], and $a_1 = 1.1$ for a wide range of the renormalization scale μ, we have the branching ratio,

$$B(B_d \to D_{sJ}^{*+}\pi^-) = 3.0 \times 10^{-5}, \tag{3}$$

close to the Belle and BaBar measurements [17,18], $B(B_d \to D_s^+\pi^-) = (2.4^{+1.0}_{-0.8} \pm 0.7, 4.6^{+1.2}_{-1.1} \pm 1.3) \times 10^{-5}$. 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 in 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,

$$\frac{1}{N_c^2} c^b \bar{s}^b n^c \bar{n}^c, \;\; \frac{1}{12} f^{ab_1c_1} f^{ab_2c_2} c^{b_1}\bar{s}^{c_1} n^{b_2}\bar{n}^{c_2}, \tag{4}$$

where both the $c\bar{s}$ and $n\bar{n} \equiv (u\bar{u} + d\bar{d})/\sqrt{2}$ pairs are in the color-singlet 1 and color-octet 8 states, respectively [19]. 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 [20], 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 $\bar{b}$ quark, is suppressed by a power of $\Lambda_{\rm QCD}/m_{D_{sJ}^*} \sim 0.1$, 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.

TABLE I. Wilson coefficients associated with the diagrams in Fig. 1.

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	$C_1(1/N_c - 1)$

A quantitative analysis of Fig. 1 requires 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,

$$\begin{aligned} 1_{ij}1_{lk} = & \frac{1}{8}(1-\gamma_5)_{ik}(1-\gamma_5)_{lj} + \frac{1}{8}(1+\gamma_5)_{ik}(1+\gamma_5)_{lj} \\ & + \frac{1}{8}[\gamma_\nu(1-\gamma_5)]_{ik}[(1-\gamma_5)\gamma^\nu]_{lj} \\ & + \frac{1}{8}[\gamma_\nu(1+\gamma_5)]_{ik}[(1+\gamma_5)\gamma^\nu]_{lj} \\ & + \frac{1}{8}(\sigma^{\nu\lambda})_{ik}(\sigma_{\nu\lambda})_{lj}, \end{aligned} \tag{5}$$

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 $\langle D_{sJ}^{*+}|\bar{c}\gamma_\mu(1-\gamma_5)s\bar{u}\gamma_\nu(1-\gamma_5)u|0\rangle$, 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 $\langle D_{sJ}^{*+}|\bar{c}\gamma_\mu(1-\gamma_5)s\bar{u}(1-\gamma_5)u|0\rangle$, 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,

$$\begin{aligned} A = & \frac{G_F}{\sqrt{2}} V_{ub}^* V_{cs} \langle D_{sJ}^{*+}|\bar{c}\gamma_\mu(1-\gamma_5)s\bar{u}\gamma_\nu(1-\gamma_5)u|0\rangle \\ & \times \langle \pi^-|\bar{b}\gamma^\mu(1-\gamma_5)u|B_d\rangle a_1 H^\nu, \end{aligned} \tag{6}$$

with the hard kernel,

$$H^\nu = \frac{g^2}{32\sqrt{2}} \frac{C_F}{N_c} \frac{\mathrm{tr}[(1-\gamma_5) l\!\!/_u \gamma_\beta (1-\gamma_5)\gamma^\nu \gamma^\beta]}{l_u^2 l_g^2}, \quad (7)$$

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 $n\bar{n}$. To be precise, H^ν 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 $O(\Lambda_{\mathrm{QCD}}/m_{D_{sJ}^*})$ [21]. Therefore, the components of l_u and l_g have the orders of magnitude

$$l_u \sim l_g \sim \frac{m_B}{\sqrt{2}} \left(\frac{\Lambda_{\mathrm{QCD}}}{m_{D_{sJ}^*}}, \frac{1}{2}, \mathbf{0}_T \right), \quad (8)$$

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 $l_u^2/l_b^2 \sim \Lambda_{\mathrm{QCD}}/m_{D_{sJ}^*}$ compared to Fig. 1(a) as stated before.

The next step is to evaluate the matrix element,

$$\langle D_{sJ}^{*+}(P_2)|\bar{c}\gamma_\mu(1-\gamma_5)s\bar{u}\gamma_\nu(1-\gamma_5)u|0\rangle = i\frac{B}{m_{D_{sJ}^*}} P_{2\mu}P_{2\nu} f_{D_{sJ}^*}, \quad (9)$$

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

$$B \approx \langle 0|\bar{u}(1-\gamma_5)u|0\rangle = \langle 0|\bar{u}u|0\rangle \approx -0.24 \ \mathrm{GeV}^3, \quad (10)$$

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

Equation (6) then becomes

$$A = i\frac{G_F}{\sqrt{2}} V_{ub}^* V_{cs}(m_B^2 - m_\pi^2) f_{D_{sJ}^*} F_0^{B\pi}(m_{D_{sJ}^*}^2) a_1 R, \quad (11)$$

with the ratio

$$R = \frac{g^2 C_F}{32\sqrt{2}N_c} B \frac{\mathrm{tr}[(1-\gamma_5) l\!\!/_u \gamma_\beta (1-\gamma_5) P\!\!\!/_2 \gamma^\beta]}{m_{D_{sJ}^*} l_u^2 l_g^2},$$

$$= -\sqrt{2}\pi^2 \left(\frac{\alpha_s}{\pi}\right) \frac{C_F}{N_c} \frac{B}{m_{D_{sJ}^*}\Lambda_{\mathrm{QCD}}^2} \left(\frac{m_{D_{sJ}^*}}{m_B}\right)^2 \approx 0.275 \quad (12)$$

for the inputs $\alpha_s/\pi = 0.2$ in $b \to c$ transitions [21] and $\Lambda_{\mathrm{QCD}} \approx 0.3$ GeV. l_u^2 and l_g^2 from Eq. (8) have been inserted. It is easy to see that the decay amplitude in the four-quark picture is down by the coupling constant α_s, by the color number $1/N_c$, and by the powers $(m_{D_{sJ}^*}/m_B)^2$. We conclude that the $B_d \to D_{sJ}^{*+}\pi^-$ branching ratio in the four-quark picture should be smaller than that in the quark-antiquark one by a suppression factor,

$$\frac{B^{(4)}(B_d \to D_{sJ}^{*+}\pi^-)}{B^{(2)}(B_d \to D_{sJ}^{*+}\pi^-)} = R^2 \approx 0.08. \quad (13)$$

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

There is already a hint from the $B \to D_{sJ}^* D$ 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 $B \to D_{sJ}^* D$ branching ratios have been measured by Belle recently [22],

$$B[B^+ \to D_{sJ}^{*+}(2317)\bar{D}^0] \times B[D_{sJ}^{*+}(2317) \to D_s^+\pi^0] = (8.1^{+3.0}_{-2.7} \pm 2.4)\times 10^{-4},$$

$$B[B^+ \to D_{sJ}^{*+}(2463)\bar{D}^0] \times B[D_{sJ}^{*+}(2463) \to D_s^{*+}\pi^0] = (11.9^{+6.1}_{-4.9} \pm 3.6)\times 10^{-4},$$

$$B[B^+ \to D_{sJ}^{*+}(2463)\bar{D}^0] \times B[D_{sJ}^{*+}(2463) \to D_s^+\gamma] = (5.6^{+1.6}_{-1.5} \pm 1.7)\times 10^{-4}. \quad (14)$$

The first data together with $B(B^+ \to D_s^+ \bar{D}^0) = (1.3 \pm 0.4)\%$ [23] imply

$$\frac{B[B^+ \to D_{sJ}^{*+}(2317)\bar{D}^0]}{B(B^+ \to D_s^+\bar{D}^0)} \approx 0.06 \quad (15)$$

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

TABLE II. $B_d \to D_{sJ}^{*-} K^{(*)+}$ branching ratios (in units of 10^{-5}) in the quark-antiquark picture from the PQCD approach.

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

$$\frac{B[B^+ \to D_{sJ}^{*+}(2463)\overline{D}^0]}{B(B^+ \to D_s^+ \overline{D}^0)} \approx 0.14, \tag{16}$$

which also gives a similar indication. It is unlikely that the dramatically different branching ratio in Eq. (15) is due to the different decay constants $f_{D_{sJ}^*}$ and f_{D_s} [24] from the viewpoint of heavy quark symmetry.

Finally, we discuss another ideal mode for our purposes, the $B_d \to D_{sJ}^{*-} K^{(*)+}$ decay, which occurs through the operators $\mathcal{O}_1 = (\overline{d}_i u_j)(\overline{c}_j b_i)$ and $\mathcal{O}_2 = (\overline{d}_i u_i)(\overline{c}_j b_j)$ with the product of the CKM matrix element $V_{cb} V_{ud}^*$. 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 [25,26]. The derivation of the factorization formulas at leading power in $1/m_B$ and leading order in α_s follows that for the $B \to D^{(*)}\pi(\rho,\omega)$ decays in [20]. We shall present the explicit expressions elsewhere. Adopting the same D_{sJ}^* meson distribution amplitudes 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 II. The predictions for the $B_d \to D_{sJ}^{*-} K^+$ mode are close to the Belle and BaBar measurements [17,18], $B(B_d \to D_s^- K^+) = (3.2 \pm 0.9 \pm 1.0, 3.2 \pm 1.0 \pm 1.0) \times 10^{-5}$, which are in agreement with the PQCD predictions [27,28]. The estimation for the $B_d \to D_{sJ}^{*-} K^{(*)+}$ 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 $B_d \to D_{sJ}^{*-}(2317)K^+$ branching ratio is expected to be around 4.2×10^{-6}.

Our estimation given above applies to other non-quark-antiquark models for the D_{sJ}^* content, such as a molecule, up to an 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 $B \to D_{sJ}^* M$ branching ratios, $M = D$, π, and K, can provide more information on the nature of the new D_{sJ}^* mesons. If the D_{sJ}^* mesons are multiquark bound states, they will be more difficult to produce than the conventional $D_s^{(*)}$ mesons in exclusive B meson decays: the branching ratios will be one order of magnitude smaller. The suppression is a combined effect of α_s, $1/N_c$, and $(m_{D_{sJ}^*}/m_B)^2$, which arise from the additional hard scattering to which the four valence quarks of the D_{sJ}^* mesons participate. More precise data are necessary for drawing a conclusion, though the recently measured $B \to D_{sJ}^* D$ branching ratios have indicated an unconventional picture.

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[1] BaBar Collaboration, B. Aubert *et al.*, Phys. Rev. Lett. **90**, 242001 (2003).
[2] CLEO Collaboration, D. Besson *et al.*, hep-ex/0305017.
[3] S. Godfrey and N. Isgur, Phys. Rev. D **32**, 189 (1985); S. Godfrey and R. Kokoski, *ibid.* **43**, 1679 (1991); M. Di Pierro and E. Eichten, *ibid.* **64**, 114004 (2001).
[4] R.N. Cahn and J.D. Jackson, Phys. Rev. D **68**, 037502 (2003).
[5] E. van Beveren and G. Rupp, Phys. Rev. Lett. **91**, 012003 (2003).
[6] T. Barnes, F.E. Close, and H.J. Lipkin, Phys. Rev. D **68**, 054006 (2003).
[7] A.P. Szczepaniak, Phys. Lett. B **567**, 23 (2003).
[8] H.Y. Cheng and W.S. Hou, Phys. Lett. B **566**, 193 (2003); K. Terasaki, Phys. Rev. D **68**, 011501(R) (2003).
[9] T. Browder, S. Pakvasa, and A.A. Petrov, Phys. Lett. B **578**, 365 (2004).
[10] H.J. Lipkin, hep-ph/0306204.
[11] G.S. Bali, Phys. Rev. D **68**, 071501(R) (2003).
[12] Y.B. Dai *et al.*, Phys. Rev. D **68**, 114011 (2003).
[13] S. Nussinov, hep-ph/0306187.
[14] S. Godfrey, Phys. Lett. B **568**, 254 (2003); P. Colangelo and F. De Fazio, *ibid.* **570**, 180 (2003).
[15] W.A. Bardeen, E.J. Eichten, and C.T. Hill, Phys. Rev. D **68**, 054024 (2003); M.A. Nowak, M. Rho, and I. Zahed, hep-ph/0307102.
[16] P. Ball, J. High Energy Phys. **09**, 005 (1998); A. Khodjamirian *et al.*, Phys. Rev. D **62**, 114002 (2000).
[17] Belle Collaboration, P. Krokovny *et al.*, Phys. Rev. Lett. **89**, 231804 (2002).
[18] BaBar Collaboration, B. Aubert *et al.*, Phys. Rev. Lett. **90**, 181803 (2003).
[19] A. Zhang, Phys. Rev. D **61**, 114021 (2000).
[20] Y.Y. Keum, T. Kurimoto, H-n. Li, C.D. Lu, and A.I. Sanda, hep-ph/0305335.
[21] T. Kurimoto, H-n. Li, and A.I. Sanda, Phys. Rev. D **67**, 054028 (2003).
[22] Belle Collaboration, P. Krokovny *et al.*, Phys. Rev. Lett. **91**, 262002 (2003).
[23] Particle Data Group, K. Hagawara *et al.*, Phys. Rev. D **66**, 010001 (2002).
[24] A. Datta and P.J. O'Donnell, Phys. Lett. B **572**, 164 (2003).
[25] G.P. Lepage and S.J. Brodsky, Phys. Lett. **87B**, 359 (1979); Phys. Rev. D **22**, 2157 (1980).

[26] Y.Y. Keum, H-n. Li, and A.I. Sanda, Phys. Lett. B **504**, 6 (2001); Phys. Rev. D **63**, 054008 (2001); Y.Y. Keum and H-n. Li, *ibid.* **63**, 074006 (2001).
[27] H. Hayakawa, K. Hosokawa, and T. Kurimoto, Mod. Phys. Lett. A **18**, 1557 (2003).
[28] C.H. Chen, Phys. Lett. B **560**, 178 (2003); C.D. Lu and K. Ukai, Eur. Phys. J. C **28**, 305 (2003); Y. Li and C.D. Lu, J. Phys. G **29**, 2115 (2003).