Are there significant scalar resonances in $B^- \to D^{(*)0}K^-K^0$ decays?

Ron-Chou Hsieh and Chuan-Hung Chen

Department of Physics, National Cheng-Kung University, Tainan 701, Taiwan

(Dated: March 26, 2022)

Abstract

We study the indication of large different branching ratios between $B^- \to D^{(*)0}K^-K^0$ and $\bar{B}_d \to D^{(*)+}K^-K^0$ observed by Belle. Interestingly, the same situation is not found in the decays $B \to D^{(*)}K^-K^{*0}$. If there exist no intermediate resonances in the decays $B^- \to D^{(*)0}K^-K^0$, a puzzle will be arisen. We find that the color-suppressed processes $B^- \to D^{(*)0}a_0(1450)$ with $a_0^{-}(1450) \to K^-K^0$ could be one of the candidates to enhance the BRs of $B^- \to D^{(*)0}K^-K^0$. Our conjecture can be examined by the Dalitz plot technique and the analysis of angular dependence on $K^-K^0$ state at $B$ factories.
Since CP violation was discovered in $K$-meson system in 1964 [1], the same thing has been also realized by Belle [2] and Babar [3] with high accuracy in the $B$ system. Although the new observation doesn’t improve our recognition on CP violation, it stimulates the development. Besides the origin of CP violation, $B$ factories also provide the chance to understand or search the uncertain states, such as the scalar bosons $f_0(400 - 1200)$, $f_0(980)$, $a_0(980)$ and glueball etc. Unlike pseudoscalar mesons, the scalar bosons with wide widths are difficult to measure directly via two-body $B$ meson decays. Inevitably, the study of three body decays will become important for extracting the quasi-two-body resonant states from the Dalitz plot and the invariant mass distribution. By the analyses of Dalitz plot technique, we can recognize whether there exist unusual structures in the phase space.

Recently, Belle has observed the BRs of the decays $B \rightarrow D^{(*)}K^-K^{*0}$ to be $BR(B^- \rightarrow D^{(*)}K^-K^{*0}) = 7.5 \pm 1.3 \pm 1.1 \ (15.3 \pm 3.1 \pm 2.) \times 10^{-4}$ and $BR(\bar{B}^0 \rightarrow D^{(*)+}K^-K^{*0}) = 8.8 \pm 1.1 \pm 1.5 \ (12.9 \pm 2.2 \pm 2.5) \times 10^{-4}$, and the decays $B \rightarrow D^{(*)}K^-K^0$ to be $BR(\bar{B}^- \rightarrow D^{(*)}K^-K^0) = 5.5 \pm 1.4 \pm 0.8 \ (5.2 \pm 2.7 \pm 1.2) \times 10^{-4}$ and $BR(\bar{B}^0 \rightarrow D^{(*)+}K^-K^0) = 1.6 \pm 0.8 \pm 0.3 \ (2.0 \pm 1.5 \pm 0.4) \times 10^{-4}$ [4]. According to Belle’s analyses, we know that $K^-K^{*0}$ system has the state $J^P = 1^+$; and also, it is pointed out that the $B \rightarrow D^{(*)}K^-K^{*0}$ decays in the low $K^-K^{*0}$ invariant mass region can be fitted well by quasi-two-body decays $B \rightarrow D^{(*)}a^-_1(1260)$ with $a^-_1(1260) \rightarrow K^-K^{*0}$. In $K^-K^0$ system of $B^- \rightarrow D^0K^-K^0$ decays, by adopting the fitting distribution $dN/d\cos\theta_{KK} \propto (R_L \cos^2\theta_{KK} + (1 - R_L) \sin^2\theta_{KK})$ Belle observed the value $R_L = 0.97 \pm 0.08$, i.e. the $K^-K^0$ state prefers being $J^P = 1^-$. Interestingly, by looking at the data shown above, we find that the BRs of charged $B$ decaying to $K^-K^0$ final state are much larger than those of neutral $B$ decays. In terms of central values, the ratios, defined by $R^{(*)} = BR(B^- \rightarrow D^{(*)}K^-K^0)/BR(\bar{B}^0 \rightarrow D^{(*)+}K^-K^0)$, could be estimated roughly to be $3.5 \ (2.6)$; however, there is no such kind of implication on the $K^-K^{*0}$ final state. Therefore, if the data display the correct behavior, there must be something happened in $B \rightarrow D^{(*)}K^-K^0$ decays. Before discussing the origin of the differences, we need to understand the mechanism to produce $K^-K^{*0}$ and $K^-K^0$ systems in charged and neutral $B$ decays.

Since $B \rightarrow D^{(*)}K^-K^{*0}$ decays are dictated by the $b \rightarrow c\bar{u}d$ transition, we describe the effective Hamiltonian as

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{ud} [C_1(\mu)\mathcal{O}_1 + C_2(\mu)\mathcal{O}_2] \quad (1)$$

with $\mathcal{O}_1 = \bar{d}_\alpha u_\beta \bar{c}_\beta b_\alpha$ and $\mathcal{O}_2^{(q)} = \bar{d}_\alpha u_\alpha \bar{c}_\beta b_\beta$, where $\bar{q}_\alpha q_\beta = \bar{q}_\alpha \gamma_\mu(1 - \gamma_5)q_\beta$, $\alpha(\beta)$ are the color
FIG. 1: Topologies for $B \to D^{(*)}K^-K^{(*)0}$ decays. (a) color-allowed with $q = u$ and $d$, and (b) color-suppressed.

indices, $V_u = V_{ud}V_{cb}$ are the products of the CKM matrix elements, $C_{1,2}(\mu)$ are the Wilson coefficients (WCs) \[\text{[5]}\]. As usual, the effective WCs of $a_2 = C_1 + C_2/N_c$ and $a_1 = C_2 + C_1/N_c$, with $N_c = 3$ being color number, are more useful. According to the interactions of Eq. (1), we classify the topologies for $B \to D^{(*)}K^-K^{(*)0}$ decays to be color-allowed (CA) and color-suppressed (CS) processes, illustrated by Fig. \[\text{[II]}\] and also use the decay amplitudes $A^Q(K^-K^{(*)0})$ to describe their decays, where $Q = C(N)$ denote the charged (neutral) $B$ decays. From the figure, we see that both CA and CS topologies contribute to charged $B$ decays but neutral $B$ decays are only governed by CA. We note that because we don’t consider the problem of direct CP asymmetry, the annihilation effects, which can contribute to $\bar{B}_d \to D^{(*)-}K^-K^{(*)0}$ decays and are usually smaller than the CA contributions, are neglected. Although Fig. \[\text{[II]}\] corresponds to the picture of three-body-decay, after removing the $s\bar{s}$ pair, the figures are related to quasi-two-body decays if $d\bar{u}$ could form a possible bound state.

Now, we have to examine that if there exists only $J^P = 1^-$ state in $K^-K^0$ system, the differences between $\text{BR}(B^- \to D^{(*)0}K^-K^0)$ and $\text{BR}(\bar{B}_d \to D^{(*)+}K^-K^0)$ are insignificant. To be more understanding the problem, it is useful to start from the discussion on $B \to D^{(*)}K^-K^0$ processes. As mentioned before, charged $B$ decays are governed by CA and CS while neutral $B$ decays are only from CA; therefore, the amplitudes for $B^-$ and $\bar{B}^0$ decays can be written as $A^C(K^-K^0) \propto a_1 M^C_a(K^-K^0) + a_2 M^C_b(K^-K^0)$ and $A^N(K^-K^0) \propto a_1 M^N_a(K^-K^0)$, where $M^Q_{a(b)}$ express the hadronic transition matrix elements for CA (CS). For simplicity, we only consider factorizable parts. As a conjecture, if there exist significant differences in BRs of $B^-$ and $\bar{B}^0$ decays, the source should be from $a_2 M^C_b(K^-K^0)$. In the following, we use two examples to show that it is impossible to
only increase the BRs of $B \to D^{(*)}K^-K^0$ without enhancing the BRs of $B \to D^{(*)}K^-K^{*0}$. Firstly, if the decays $B \to D^{(*)}K^-K^0$ are only governed by quasi-two-body decays, say $B \to D^{(*)}\rho_X$ with $\rho_X$ being arbitrary vector meson, the decay amplitudes can be described by $A^C(K^-K^0) \sim a_1 f_{\rho_X} F^{B\to D^{(*)}}(0) + a_2 f_{D^{(*)}} F^{B\to \rho_X}(0)$ and $A^N(K^-K^0) \sim a_1 f_{\rho_X} F^{B\to D^{(*)}}(0)$, where $f_{D^{(*)}}$ and $f_{\rho_X}$ are the decay constants of $D^{(*)}$ and $\rho_X$ mesons, respectively and $F^{B\to D^{(*)}(\rho_X)}$ denote the $B \to D^{(*)}(\rho_X)$ form factors. However, it is known that $|a_2/a_1| \sim 0.26$ and $f_{\rho_X} F^{B\to D^{(*)}}(0) \sim f_{D^{(*)}} F^{B\to \rho_X}(0)$ [8]. If the resonance is $J^P = 1^-$ state, the differences in BR between charged and neutral $B$ decays will be slight. The same argument can be applied to $B \to D^{(*)}K^-K^{*0}$ decays for $J^P = 1^+$ resonant state. Hence, the consequences are consistent with the observations of $B \to D^{(*)}K^-K^{*0}$ decays. Secondly, we consider nonresonant mechanism on $B \to D^{(*)}K^-K^0$. Since the situation corresponds to a three-body phase space, we use the $K^-K^0$ invariant mass, expressed by $\omega$, as the variable to describe the behavior of the decay. Hence, the decay amplitudes could be written as $A^C(K^-K^0) \sim a_1 F^{0\to KK}(\omega^2) F^{B\to D^{(*)}}(\omega^2) + a_2 f_{D^{(*)}} F^{B\to \rho_X}(0)$ and $A^N(K^-K^0) \sim a_1 F^{0\to KK}(\omega^2) F^{B\to D^{(*)}}(\omega^2)$, and the $0\to KK$ form factors are the associated form factors. It is known that at asymptotic region, in terms of perturbative QCD (PQCD) the behavior of $F^{0\to KK}$ has the power law $1/\omega^2(\ln \omega^2/\Lambda_{QCD}^2)^{-1}$ [6]. If we take the behavior of $F^{B\to D^{(*)}}$ to be $1/(1 - \omega^2/M_X^2)^n$, $M_X$ is the pole mass and $n = 1$ or 2, we clearly see that the decay amplitudes are suppressed at large $\omega$. Furthermore, according to the analysis of Ref. [8], the dominant region actually is around $\omega^2 \sim \bar{\Lambda}_M$. Therefore, if $a_2 F^{B\to \rho_X}(\bar{\Lambda}_M)$ is the source of the differences in BR between charged and neutral $B$ decays, the similar effects $a_2 F^{B\to KK^*}(\bar{\Lambda}_M)$ will also affect the decays $B \to D^{(*)}K^-K^{*0}$. Hence, unless $F^{B\to KK^*}(\bar{\Lambda}_M)$ is much less than $F^{B\to KK}(\bar{\Lambda}_M)$, otherwise, the CS effects should have the similar contributions to $B \to D^{(*)}K^-K^{*0}$. It is known that with the concept of two-meson wave functions [9], the calculations of $F^{B\to KK}(\bar{\Lambda}_M)$ and $F^{B\to KK^*}(\bar{\Lambda}_M)$ can be associated with the wave functions of $KK$ and $KK^*$ systems. At $\omega^2 \sim \bar{\Lambda}_M$ region, two-meson wave functions could be described by two individual meson wave functions [10]. Refer to the derived $K$ and $K^*$ wave functions [11], the value of $F^{B\to KK}(\bar{\Lambda}_M)/F^{B\to KK^*}(\bar{\Lambda}_M)$ should be close to $F^{B\to K}(0)/F^{B\to K^*}(0) \sim 0.8$. As a result, it is hard to imagine that the CS effects on pure three-body decays can make large differences in $B \to D^{(*)}K^-K^0$ processes but not in $B \to D^{(*)}K^-K^{*0}$ processes.

Inspired by Belle’s observation, we think the significant differences in the BRs of $B \to
$D^{(*)}K^-K^0$ should not come from $K^-K^0$ with $J^P = 1^-$ state but with $J^P = 0^+$ state. This could be easily understood as follows: as discussed before, only CS effects, $a_2 M^C_K (K^-K^0)$, could make a discrepancy in BRs of charged and neutral $B$ decays. Therefore, the mechanism which dominantly contributes to CS topology will be our candidate. If $K^-K^0$ has the $J^P = 0^+$ state, due to $F^{0 \to KK} \propto \langle KK | \bar{u} \gamma_\mu d | 0 \rangle$, in terms of equation of motion, we get that $F^{0 \to KK}$ is proportional to $(m_d - m_u)$. That is, the contribution of scalar state to CA topology is negligible. On the contrary, there is no such suppression to the CS topology. Now, the problem becomes how to produce $K^-K^0$ to be a $0^+$ state. By searching particle data group \cite{12}, we find that the preference could be the scalar bosons $a_0(980)$ and $a_0(1450)$ because both are isovector states and have sizable decay rates for $a_0 \to KK$. Hence, to realize our thought, we propose that the decays $B^- \to D^{(*)0}a_0^- (980, 1450)$ with $a_0^- (980, 1450) \to K^-K^0$ can satisfy the required criterion to enhance the BRs of $B^- \to D^{(*)0}K^-K^0$.

Since the quark contents of scalar mesons below or near 1 GeV are still obscure in the literature \cite{13}, for avoiding the difficulty in estimation, we only make the explicit calculations on $a_0(1450)$ which have definite composed structure of $qq$. If regarding $a_0(980)$ consists of $q\bar{q}$ state, the same estimation could be also applied \cite{14}. In theory, it is known that the serious problem on two-body nonleptonic $B$ decays comes from the calculations of hadronic matrix elements. Since the involving processes are governed by CS topologies, like the well known decays $B \to J/\Psi K^{(*)}$ and $\bar{B}_d \to D^{(*)0}\pi^0$ in which nonfactorizable effects play an important role, we adopt the PQCD approach which is described by the convolution of hard amplitude and wave functions \cite{15}, can deal with the factorizable and nonfactorizable contributions and can avoid the suffering end-point singularities self-consistently \cite{16}.

In the calculations of hadronic effects, the problem is how to determine the wave functions of $D^{(*)}$ and $a_0$ mesons. For $D^{(*)}$ meson wave functions, we could model them to fit with the measured BRs of $\bar{B}_d \to D^{(*)0}\pi^0$ decays \cite{17}, in which $B$ meson wave function is chosen to be coincide with the observed BRs of $B \to \pi\pi$ decays while $\pi$ meson ones are adopted from the derivation of QCD sum rules \cite{11}. As to $a_0$ scalar meson, the spin structures of $a_0$ are required to satisfy $\langle 0 | \bar{u} \gamma_\mu d | a_0^-, p_3 \rangle = [(m_d - m_u)/m_{a_0}] \tilde{f} p_3^\mu$ and $\langle 0 | \bar{q}q | a_0 \rangle = m_{a_0} \tilde{f}$, where $m_{a_0}$ and $\tilde{f}$ are the mass and decay constant of $a_0$. In order to satisfy these local current matrix
elements, the distribution amplitude for $a_0$ is adopted as
\begin{equation}
\int_0^1 dx e^{ixP^\mu z} \langle 0 | q(0) q(z) | a_0 \rangle = \frac{1}{\sqrt{2N_c}} \left\{ [\bar{\psi}]_q \Phi_{a_0}(x) + m_{a_0}[1]_q \Phi_{a_0}(x) \right\}.
\end{equation}

By the charge parity invariance and neglecting the effects of light current quark mass $m_{u(d)}$, we obtain $\Phi_{a_0}(x) \approx -\Phi_{a_0}(1-x)$ and $\Phi^p_{a_0}(x) = \Phi^p_{a_0}(1-x)$, and their normalizations are $\int_0^1 dx \Phi_{a_0}(x) = 0$ and $\int_0^1 dx \Phi^p_{a_0}(x) = \int/2\sqrt{2N_c}$. Although vector $D^{*0}$ meson carries the spin degrees of freedom, in the $B^- \to D^{*0}a_0^-$ decay only longitudinal polarization is involved. The results should be similar to $D^0a_0^-$ case. Therefore, we only present the representative formulas for $B^- \to D^0a_0^-$ decay. Hence, the decay amplitude for $B^- \to D^0a_0^-$ is read as
\begin{equation}
A = V_u \left[ f_D F_{D^0a_0}^e + M_{D^0a_0}^e \right]
\end{equation}

where $F^e(M^e)$ is the factorized (non-factorized) emission hard amplitudes. According to Eq. (2), the typical hard functions are expressed as
\begin{equation}
F_{D^0a_0}^e = \eta \int_0^1 dx_1 dx_3 \int_0^\infty d_1 d_3 d_2 \Phi_B(x_1, b_1)
\left\{ \left[ (\zeta_1 + \zeta_2 x_3) \Phi_{a_0}(x_3) + r_a(1 - 2x_3) \Phi^p_{a_0}(x_3) \right] \right.
\left. + \mathcal{E}(t_3^e) + r_a(2\Phi^p_{a_0}(x_3) - r_a \Phi_{a_0}(x_3)) \right\}
\end{equation}

\begin{equation}
M_{D^0a_0}^e = 2\eta \int_0^1 dx d_1 d_3 d_2 \Phi_B(x_1, b_1) \Phi_D(x_2)
\left\{ \left[ \left( 1 - r_a^2 \right) x_2 + \left( 1 - r_D^2 \right) x_3 \right] \Phi_{a_0}(x_3)
+ r_a x_3 \Phi^p_{a_0}(x_3) \right\} \mathcal{E}(t_3^e) + \left[ \left( \zeta_1 - r_a^2 \right) (1 - x_2)
+ r_a x_3 \Phi^p_{a_0}(x_3) \right] \mathcal{E}(t_3^e)
\end{equation}

with $\eta = 8\pi C_F M_B^2$, $r_{a(D)} = m_{a(D)}/M_B$, $\zeta_1 = 1 - r_a^2 - r_D^2$, $\zeta_2 = \zeta_1 - r_D^2$, $\mathcal{E}(t_3^e) = \alpha_s(t_3^e) a_2(t_3^e)$ $S \mu \nu_{B^0a_0}(t_3^e) h_\nu \{x, \{b\}\}$ and $\mathcal{E}(t_3^e) = \alpha_s(t_3^e)(C_2(t_3^e)/N_c)$ $S \mu \nu_{B^0a_0}(t_3^e) h_\nu \{x, \{b\}\}$. $t_{e,d}$, $S\mu \nu$ and $h_{e,d}$ denote the hard scales of $B$ decays, Sudakov factors and hard functions which are arisen from the propagators of gluon and internal valence quark, respectively. Their explicit expressions can be found in Ref. [19].
For numerical estimations, the $B$, $D^{(*)0}$ and $a_0$ meson wave functions are simply chosen as

$$
\Phi_B(x, b) = N_B x^2 (1 - x)^2 \exp \left[ -\frac{1}{2} \left( \frac{x m_B}{\omega_B} \right)^2 - \frac{\omega_B^2 b^2}{2} \right]
$$

$$
\Phi_{D^{(*)}}(x) = \frac{3}{\sqrt{2 N_c}} f_{D^{(*)}} x (1 - x) [1 + 0.7(1 - 2x)],
$$

$$
\Phi_{a_0}(x) = \frac{\tilde{f}}{2\sqrt{2 N_c}} \left[ 6x(1 - x)C_1^{3/2}(1 - 2x) \right],
$$

$$
\Phi'_{a_0}(x) = \frac{\tilde{f}}{2\sqrt{2 N_c}} \left[ 3(1 - 2x)^2 \right],
$$

(5)

where $N_B$ is determined by the normalization of $\int_0^1 dx \Phi_B(x, 0) = f_B/(2\sqrt{2 N_c})$ and $C_1^{3/2}(y)$ is the Gegenbauer polynomial. The values of theoretical inputs are set as: $\omega_B = 0.4$, $f_B = 0.19$, $f_{D^{(*)}} = 0.20$ (0.22), $\tilde{f} = 0.20$, $m_B = 5.28$ and $m_{D^{(*)}} = 1.87$ (2.01) GeV. By the taken values and using Eq. (5) with excluding WC of $a_2$, we immediately obtain the $B \to a_0(1450)$ form factor at large recoil to be 0.44. The result is quite close to the value 0.46 which is estimated by the light-cone sum rules [20]. Hence, the magnitudes of the considered hard functions are given in Table I. Consequently, the BRs of $B^− \to D^{(*)}a_0^−(1450)$ are

| $f_{D^{(*)}}^{F_{D^{(*)0}a_0}}$ | $M_{D^{(*)0}a_0}^r$ | $f_{D^{(*)}}^{F_{D^{(*)0}a_0}}$ | $M_{D^{(*)0}a_0}^r$ |
|-----------------|-----------------|-----------------|-----------------|
| -1.42          | -2.22 + i0.97  | -1.56          | -2.39 + i1.06  |

obtained to be $8.21 \times 10^{-4}$. Since the predictions of PQCD on the BRs of $\bar{B} \to D^0\pi^0$ [17] and $B \to J/\Psi K^{(*)}$ and the helicity amplitudes of $B \to J/\Psi K^{(*)}$ [21] are consistent with the current experimental data, with the same approach, our results should be reliable. Furthermore, the BR products of $Br(B^− \to D^{(*)0}a_0^−(1450)) \times Br(a_0^−(1450) \to K^−K^0) \approx 1.81 \times 10^{-4}$ with $Br(a_0^−(1450) \to K^−K^0) \approx 0.22$ [12]. Clearly, the contributions of quasi-two-body decays to $B^− \to D^{(*)}K^−K^0$ modes are close to the pure three-body decays $\bar{B} \to D^{(*)}K^−K^0$ [22].

In summary, we have investigated that when a proper scalar meson is considered, the BRs of $B^− \to D^{(*)0}K^−K^0$ will deviate from those of $\bar{B}_d \to D^{(*)}K^−K^0$. Although we only concentrate on $a_0(1450)$, the same discussion is also applicable to $a_0(980)$. Since our purpose
is just to display the importance of scalar boson on the decays $B^- \rightarrow D^{(*)0}K^-K^0$, we don’t consider the theoretical uncertainties at this stage. And also, we neglect discussing the interfering effects of resonance and non-resonance. It is worthwhile to mention that by using powerful Dalitz plot technique, many scalar mesons in charm decays have been observed by the experiments at CLEO [23], E791 [24], FOCUS [25], and Babar [26]. Expectably, the significant evidences of scalar productions should be also observed at $B$ factories. In addition, since the scalar meson doesn’t carry spin degrees of freedom, there is no specific direction for $K^-K^0$ production so that we should see a different angular distribution of $K^-K^0$, such as the coefficient associated with the term $\sin^2 \theta_{KK}$ in the fitting distribution mentioned early will be enhanced. Hence, with more data accumulated, our conjecture can be examined by the Dalitz plot and the analysis of angular dependence on $K^-K^0$ state.

Acknowledgments:

The author would like to thank C.Q. Geng, H.N. Li and H.Y. Cheng for their useful discussions. This work was supported in part by the National Science Council of the Republic of China under Grant No. NSC-92-2112-M-006-026 and by the National Center for Theoretical Sciences of R.O.C..

[1] J.H. Christenson, J.W. Cronin, V.L. Fitch and R. Turly, Phys. Rev. Lett. 13, 138 (1964).
[2] Belle Collaboration, A. Abashian et al., Phys. Rev. Lett. 86, 2509 (2001).
[3] Babar Collaboration, B. Aubert et al., Phys. Rev. Lett. 86, 2515 (2001).
[4] Belle Collaboration, A. Drutskoy et al., Phys. Lett. B542, 171 (2002).
[5] G. Buchalla, A.J. Buras and M.E. Lautenbacher, Rev. Mod. Phys. 68, 1230 (1996).
[6] H.Y. Cheng and K.C. Yang, Phys. Rev. D59, 092004 (1999).
[7] S.J. Brodsky and G.R. Farrar, Phys. Rev. D11, 1309 (1975).
[8] C.H. Chen and H.N. Li, Phys. Lett. B561, 258 (2003).
[9] D. Muller et al., Fortschr. Physik. 42, 101 (1994); M. Diehl, T. Gousset, B. Pire, and O. Teryaev, Phys. Rev. Lett. 81, 1782 (1998); M.V. Polyakov, Nucl. Phys. B555, 231 (1999).
[10] M. Diehl, Th. Feldmann, P. Kroll, and C. Vogt, Phys. Rev. D61, 074029 (2000); M. Diehl, T.
[11] P. Ball et al., Nucl. Phys. B529, 323 (1998); P. Ball, JHEP 01, 010 (1999).
[12] Particle data group, K. Hagiwara et al., Phys. Rev. D66, 010001 (2002).
[13] H.Y. Cheng, Phys. Rev. D67, 034024 (2003).
[14] C.H. Chen, Phys. Rev. D67, 014012 (2003).
[15] G.P. Lepage and S.J. Brodsky, Phys. Lett. B87, 359 (1979); Phys. Rev. D22, 2157 (1980).
[16] H.N. Li, Phys. Rev. D64, 014019 (2001).
[17] H.N. Li, presented at FPCP, Philadelphia, Pennsylvania, 16-18 May 2002, hep-ph/0210198.

T. Kurimoto et al., Phys. Rev. D67, 054028 (2003); Y. Y. Keum et al., hep-ph/0305335.
[18] V.L. Chernyak and A.R. Zhitnitsky, Phys. Rep. 112, 173 (1984).
[19] C.H. Chen and H.N. Li, Phys. Rev. D63, 014003 (2001).
[20] V. Chernyak, Phys. Lett. B509, 273 (2001).
[21] C.H. Chen, Phys. Rev. D67, 094011 (2003).
[22] C.K. Chua, W.S. Hou, S.Y. Shiau and S.Y. Tsai, Phys. Rev. D67, 034012 (2003); Z.T. Wei, hep-ph/0301174.
[23] CLEO Collaboration, H. Muramatsu, et al., Phys. Rev. Lett. 89, 251802 (2002).
[24] E791 Collaboration, E.M. Aitala et al., Phys. Rev. Lett. 89, 121801 (2002).
[25] FOCUS Collaboration, J.M. Link et al., Phys. Lett. B541, 227 (2002).
[26] Babar Collaboration, B. Aubert et al., hep-ex/0207089.