Understanding the branching ratios of $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\phi\omega$ observed at BES-III

Dian-Yong Chen$^{1,3}$, Jun He$^{1,3}$, Xue-Qian Li$^4$, and Xiang Liu$^{1,2,\ast}$

$^1$Research Center for Hadron and CSR Physics, Lanzhou University & Institute of Modern Physics of CAS, Lanzhou 730000, China
$^2$School of Physical Science and Technology, Lanzhou University, Lanzhou 730000, China
$^3$Institute of Modern Physics, Chinese Academy of Sciences, Lanzhou 730000, China
$^4$Physics department, Nankai University, Tianjin 300071, China

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In this work, we discuss the contribution of the mesonic loops to the decay rates of $\chi_{c1} \to \phi\phi$, $\omega\omega$ which are suppressed by the helicity selection rules and $\chi_{c1} \to \phi\omega$ which is a double-OZI forbidden process. We find that the mesonic loop effects naturally explain the clear signals of $\chi_{c1} \to \phi\phi$, $\omega\omega$ decay modes observed by the BES collaboration. Moreover, we investigate the effects of the $\omega - \phi$ mixing which may result in the order of magnitude of the branching ratio $BR(\chi_{c1} \to \omega\phi)$ being $10^{-7}$. Thus, we are waiting for the accurate measurements of the $BR(\chi_{c1} \to \omega\omega)$, $BR(\chi_{c1} \to \phi\phi)$ and $BR(\chi_{c1} \to \omega\phi)$, which may be very helpful for testing the long-distant contribution and the $\omega - \phi$ mixing in $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\omega\phi$ decays.

I. INTRODUCTION

Charm physics is an active field full with chances and challenges [1]. Decays of charmonia may provide an ideal laboratory to study perturbative as well as non-perturbative quantum chromodynamics (QCD). Until now, most of hadronic decays of P-wave charmonium states $\chi_{cJ}$ ($J = 0, 1, 2$) are not well understood compared to the $J/\psi$ decays, so that they cause great interests of the experimentalists and theorists to further explore the decay behavior of $\chi_{cJ}$.

In the Hadron 2009 conference, the BES-III Collaboration announced its observations of $\chi_{cJ}$ ($J = 0, 1, 2$) decaying into light vector mesons, where the data of $\chi_{cJ}$ are taken from 110 million radiative-decay events of $\psi(2S)$ collected at the BES-III. Among those decay modes of $\chi_{cJ}$ decaying into light vector mesons, $\chi_{c1} \to \phi\phi$, $\omega\omega$ processes were measured for the first time and the doubly OZI suppressed process $\chi_{c1} \to \phi\omega$ had not been measured before the BES-III observation [2].

Generally, the decays of $\chi_{c1}$ into two light vector mesons are suppressed compared to the corresponding decays of $\chi_{c0}$ and $\chi_{c2}$ due to the helicity selection rule [3]. Besides, $\chi_{c1} \to \omega\phi$ suffers from the double-OZI suppression [4]. Thus, in the sense, the branching ratios of the channels $\chi_1 \to \phi\phi$, $\omega\omega$, $\phi\omega$ should be small and it would be difficult to observe them in experiments, especially $\chi_{c1} \to \phi\omega$. However, the observation of $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\phi\omega$ seems to be surprising and compel us to reconsider what mechanism plays the dominant role in those decays. It would be definitely different from that responsible for $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\phi\omega$ decays. To understand the governing mechanism which results in sizable ratios for $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\phi\omega$, two questions must be answered: (1) what is the source to alleviate the helicity selection rule for $\chi_{c1}$; (2) why the double-OZI suppression is violated for $\chi_{c1} \to \phi\omega$ decay.

The conventional decay mechanism depicting $\chi_{cJ}$ into two light vector mesons is that $c$ and $\bar{c}$ annihilate into a pair of gluons, which then transit into quark-antiquark pairs to form the light vector mesons in the final state. The helicity selection rule manifests in the $\chi_{cJ}$ decays, and results in the suppression of $\chi_{c1}$ decaying into two light mesons. In $\chi_{c1}$ decays, the non-perturbative QCD effect plays a crucial role. As an important non-perturbative effect, the hadronic loop contributions, which were introduced in Refs. [5, 6] and applied to study charmonium decay [7, 8] and open-charm and hidden charm decays of charmonium-like states $X$, $Y$, $Z$ [9, 10] extensively, would change the whole scenario from the conventional decay mechanism of charmonium.

For $\chi_{c1}$ into two light vector mesons, a quark level description of the hadronic loop contribution is presented in Fig. 1. Here, red fermion line denotes charm quark, blue and green lines represent the light quarks. $\chi_{c1}$ first dissolves into two virtual charmed mesons, then by exchanging an appropriate hadron (i.e. it possesses appropriate charge, flavor, spin and isospin) they turn into two on-shell real light hadrons, which can be caught by detector 9. The matrix element of $\chi_{c1}$ into two light vector
mesons via hadronic loop effect can be described as

$$\mathcal{M}(\chi_{c1} \to VV) = \sum_i \langle VV | \mathcal{H}^{(2)}_i | i \rangle \langle i | \mathcal{H}^{(1)} | \chi_{c1} \rangle.$$  \hspace{1cm} (1)

The depiction at the hadron level corresponding to the quark level diagrams are presented in Fig. 1. Here, the suitable intermediated charmed mesons for the decay of $\chi_{c1}$ into two light vector mesons should be $D_{(s)}^1 \bar{D}_{(s)}^{*0} + h.c.$, which interact with $\chi_{c1}$ via S-wave. The exchanged mesons include pseudoscalar and vector charmed mesons. Thus, by the hadronic loop mechanism, the transition of $\chi_{c1}$ into two light vector mesons would not be suppressed by the helicity selection rule.

Another important motivation is how $\chi_{c1} \to \phi \omega$ evades the double-OZI suppression rule. Here, $\phi$ and $\omega$ mesons are ideal mixtures of the flavor $SU(3)$ octet $\omega_8 = (u\bar{u} + d\bar{d} - 2s\bar{s})/\sqrt{6}$ and the singlet $\phi_0 = (u\bar{u} + d\bar{d} + s\bar{s})/\sqrt{3}$. In terms of the hadronic loop mechanism, such $\chi_{c1} \to \phi \omega$ decay is fully forbidden. In reality $\omega$ and $\phi$ are not ideal mixtures of the flavor $SU(3)$ octet and singlet, which could provide a source which violates the double-OZI suppression rule for $\chi_{c1} \to \phi \omega$.

In this work, we will combine the hadronic loop effect with $\omega - \phi$ mixing to study $\chi_{c1} \to \phi \phi, \omega \omega, \phi \omega$. The paper is organized as follows. After introduction, we present the formula of hadronic loop contribution to $\chi_{c1} \to \phi \phi, \omega \omega, \phi \omega$ with the mixing schemes for $\omega - \phi$. Then, the numerical results about $\chi_{c1} \to \omega \omega, \phi \phi, \omega \phi$ are given in Sec. III. Finally the paper ends with a discussion and a short summary.

II. HADRONIC LOOP EFFECT ON $\chi_{c1} \to VV$ DECAYS UNDER TWO DIFFERENT MIXING SCHEMES OF $\omega - \phi$

Firstly, we present the mixing scheme for $\omega - \phi$ used in this work

$$\begin{pmatrix}
|\phi_0^p \rangle \\
|\omega_0^p \rangle
\end{pmatrix} = \begin{pmatrix}
\cos \theta & \sin \theta \\
-\sin \theta & \cos \theta
\end{pmatrix}
\begin{pmatrix}
|\phi^f_0 \rangle \\
|\omega^f_0 \rangle
\end{pmatrix},$$  \hspace{1cm} (2)

where $|\phi_0^p \rangle$ ($|\omega_0^p \rangle$) and $|\phi^f_0 \rangle$ ($|\omega^f_0 \rangle$) are the physical and ideally mixing states, respectively. The flavor wave functions for the ideally mixing states $|\omega^f \rangle$ and $|\phi^f \rangle$ are $\omega^f = (u\bar{u} + d\bar{d} + 2s\bar{s})/\sqrt{6}$ and $\phi^f = -s\bar{s}$.

Taking mixing angle $\theta = 0^\circ$ corresponds to the ideal mixing. According to the analysis in Refs. [16, 17, 22], the mixing angle $\theta$ should be $(3.4 \pm 0.2)^\circ$ in the mixing scheme in Eq. (2).

The effective Lagrangians which are responsible for the decay amplitudes for the diagrams in Fig. 1 is listed below [23, 24]

$$\mathcal{L}_{\chi_{c1} DD^*} = ig_{\chi_{c1} DD} \cdot D_{i}^f D_{i}^{*f} + h.c., \hspace{1cm} (3)$$

$$\mathcal{L}_{\bar{H}H} = i g_{\bar{H}H} D_{i}^f D_{j}^{*f} \cdot (D_{i}^0 D_{j}^{*0} + h.c.), \hspace{1cm} (4)$$

where Eq. (4) is constructed under the chiral and heavy quark limits. The superfield $H$ is given by $H = \frac{1}{2}(D^{*}\gamma_{\mu} \gamma_{\mu} - ig_5 D)$ and $\bar{H} = \gamma^0 H^\dagger \gamma^0$. $(V_\mu)_{ba}$ and $(A_\mu)_{ba}$ denotes the matrix elements for vector and axial currents, respectively. The expansion in Eq. (4), which is related to the hadronic loop calculation, includes

$$\mathcal{L}_{D^{(*)} D^{(*)} V} =$$

$$-ig_{D^*DV} D_{i}^f D_{j}^{*f} \partial_{\mu} (V_{\mu})_{ij}$$

$$-2f_{D^*DV} \epsilon_{\mu\nu\alpha\beta} (D^\alpha V_{\mu}^{\nu} - D^\nu V_{\mu}^{\alpha}) D_{i}^{*f} \partial_{\beta} D_{j}^{*f}$$

$$-D_{i}^{*f} \partial_{\alpha} D_{j}^{*f} + ig_{D^*DV} D_{i}^{*f} \partial_{\mu} (D_{j}^{*f} V_{\mu}^{*}) + D_{i}^{*f} \partial_{\mu} (D_{j}^{*f} V_{\mu}^{*}),$$  \hspace{1cm} (5)

where $D_{i}^{(s)} = (D_{i}^{(s)0}, D_{i}^{(s)+}, D_{i}^{(s)}-)$. The coupling constants relevant to the calculation include: $g_{D^*DV} =$
If we pre-assume that the mixing parameters in the matrix are not independent, but related to each other by a single variable $\theta$, the coefficients $\kappa, \zeta, \delta, \sigma$ are written as

$$
\kappa = \frac{1}{\sqrt{2}} \cos \theta, \quad \zeta = \frac{1}{\sqrt{2}} \sin \theta,
\delta = -\sin \theta, \quad \sigma = \cos \theta.
$$

(6)

Thus, the decay amplitudes of $\chi_{c1} \rightarrow \omega \omega, \phi \phi, \phi \omega$ due to the hadronic loop effect are written as

$$
\mathcal{M}[\chi_{c1} \rightarrow \omega \omega] = 2 \left( \begin{array}{cccc} \kappa^2 & \kappa^2 & \kappa^2 & \kappa^2 \\ \kappa^2 & \delta & \delta & \delta \\ \kappa^2 & \delta & \delta & \delta \\ \kappa^2 & \delta & \delta & \delta \end{array} \right) \mathcal{M}^{(P)}[D^- \rightarrow D^{++}, D^{--}, \omega, \omega],
$$

(7)

$$
\mathcal{M}[\chi_{c1} \rightarrow \phi \phi] = 2 \left( \begin{array}{cccc} \zeta^2 & \zeta^2 & \zeta^2 & \zeta^2 \\ \zeta^2 & \sigma & \sigma & \sigma \\ \zeta^2 & \sigma & \sigma & \sigma \\ \zeta^2 & \sigma & \sigma & \sigma \end{array} \right) \mathcal{M}^{(P)}[D^- \rightarrow D^{++}, D^{--}, \phi, \phi],
$$

(8)

$$
\mathcal{M}[\chi_{c1} \rightarrow \phi \omega] = 2 \left( \begin{array}{cccc} \kappa \zeta & \kappa \zeta & \kappa \zeta & \kappa \zeta \\ \kappa \zeta & \kappa \zeta & \kappa \zeta & \kappa \zeta \\ \kappa \zeta & \kappa \zeta & \kappa \zeta & \kappa \zeta \\ \kappa \zeta & \kappa \zeta & \kappa \zeta & \kappa \zeta \end{array} \right) \mathcal{M}^{(V)}[D^- \rightarrow D^{++}, D^{--}, \phi, \omega].
$$

(9)

where the factor 2 is from the charge conjugation transformation. In the above expressions, $\mathcal{M}^{(P)}[* \cdots *]$ and $\mathcal{M}^{(V)}[* \cdots *]$ denote the amplitudes corresponding to pseudoscalar and vector charmed meson exchanges. The general expressions of the amplitudes are

$$
\mathcal{M}^{(P)}[A(p_1), B(p_2), C(q), D(p_3), E(p_4)]
$$

$$
= \int \frac{d^4q}{(2\pi)^4} \frac{ig_{\chi_{c1}D^*}\epsilon^*}{[ig_{D^*D}D^*}\epsilon^*][i(2\pi)^4\delta[q-p_1-q_2]]\left[-g^\beta + \frac{p^\beta q^\alpha}{m_D^2}\right]^2 \delta^2(q^2, m_D^2),
$$

(10)

$$
\mathcal{M}^{(V)}[A(p_1), B(p_2), C(q), D(p_3), E(p_4)]
$$

$$
= \int \frac{d^4q}{(2\pi)^4} \frac{ig_{\chi_{c1}D^*}\epsilon^*}{[ig_{D^*D}D^*}\epsilon^*][i(2\pi)^4\delta[q-p_1-q_2]]\left[-g^\beta + \frac{p^\beta q^\alpha}{m_D^2}\right]^2 \delta^2(q^2, m_D^2),
$$

(11)

which correspond to pseudoscalar and vector charmed meson exchanges, respectively. Here, $A$ and $B$ denote
the intermediate charmed mesons. \( C \) is the exchanged charmed meson. \( D \) and \( E \) mean the light vector mesons in the final states. We adopt the form factor with the pole form

\[
\bar{F}_N(q^2, m^2) = \left( \frac{\Lambda^2 - m^2}{\Lambda^2 - q^2} \right)^N, \tag{12}
\]

which depicts the inner structure of the effective vertex of the exchanged charmed meson and intermediated states. Meanwhile, the form factor with pole form also plays the role to make the ultraviolet divergence disappear, in analog to the cut-offs in the Pauli-Villas renormalization scheme. Here, the cutoff \( \Lambda \) can be parameterized as \( \Lambda = m + \alpha \Lambda_{QCD} \) with \( \Lambda_{QCD} = 220 \text{ MeV} \) and \( m \) is the mass of the exchanged meson \([20]\).

### III. NUMERICAL RESULTS

As a free parameter, \( \alpha \) is introduced by the cutoff \( \Lambda \). The value is usually dependent on the particular process and taken to be of the order of unity. The BES-II collaboration reported the branching ratio of \( \chi_{c1} \to K^{(*)0}(892)\bar{K}^{(*)0}(892) \) as \( BR(\chi_{c1} \to K^{(*)0}\bar{K}^{(*)0}) = (1.67 \pm 0.32 \pm 0.31) \times 10^{-3} \) \([30]\), which can be applied to determine \( \alpha \) assuming that the hadronic loop effect is dominant in the process \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \). The formula of hadronic loop contribution to \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \) is similar to that for \( \chi_{c1} \to \omega\omega, \phi\phi, \omega\phi \) decays, where the exchanged charmed meson matching to intermediated states \( D^+D^- + \text{h.c.} \) and \( D^+_1D^-_1 + \text{h.c.} \) are \( D^{(*)\pm} \) and \( D^{(*)} \), respectively. Our study indicates that the dipole form factor, i.e. taking \( N = 2 \) in Eq. (12), can well reproduce the branching ratio of \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \) with \( \alpha = 1.14 \sim 1.28 \), which seems to be reasonable\(^1\).

In Fig. 2 we show the dependence of the ratios of \( BR(\chi_{c1} \to \omega\phi)/BR(\chi_{c1} \to \phi\phi) \) and \( BR(\chi_{c1} \to \omega\phi)/BR(\chi_{c1} \to \omega\omega) \) on \( \theta \) considering the \( \omega - \phi \) mixing. In our calculations, we find these two ratios weakly depend on the parameter \( \alpha \). In the figure a typical value of \( \alpha = 1.20 \) is employed. For the ideal mixing, i.e. \( \theta = 0^\circ \), the coefficients \( \zeta \) and \( \delta \) are zero, thus, the mesonic loop contribution to the decay \( \chi_{c1} \to \omega\phi \) vanishes. Even though considering the deviation from the ideal mixing, the ratios of \( BR(\chi_{c1} \to \omega\phi)/BR(\chi_{c1} \to \phi\phi) \) and \( BR(\chi_{c1} \to \omega\phi)/BR(\chi_{c1} \to \omega\omega) \) are still rather small. In the region \( \theta = (3.4 \pm 0.2)^\circ \), the branching ratio of \( \chi_{c1} \to \omega\phi \) is three orders smaller than those of \( \chi_{c1} \to \omega\omega \) and \( \chi_{c1} \to \phi\phi \). In Fig. 3 we present the \( \theta \) and \( \alpha \) dependence of the branching ratios of \( \chi_{c1} \to \omega\omega, \chi_{c1} \to \phi\phi \) and \( \chi_{c1} \to \omega\phi \). The corresponding ranges of \( BR(\chi_{c1} \to \omega\omega), BR(\chi_{c1} \to \phi\phi), BR(\chi_{c1} \to \omega\phi) \) with \( \alpha = 1.14 \sim 1.28 \) and \( \theta = (3.4 \pm 0.2)^\circ \) are listed in the second column of Table I. In the table, one can notice, the branching ratio of \( \chi_{c1} \to \omega\phi \) is at most of the order \( 10^{-7} \).

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|}
\hline
Channel & Branching ratio & Experimental value \\
\hline
\( K^{(*)0}\bar{K}^{(*)0} \) & (10.83 \sim 23.25) \times 10^{-4} & (16.7 \pm 3.2 \pm 3.1) \times 10^{-4} \\
\omega\omega & (4.822 \sim 10.366) \times 10^{-4} & — \\
\phi\phi & (1.465 \sim 3.238) \times 10^{-4} & — \\
\omega\phi & (2.542 \sim 6.893) \times 10^{-7} & — \\
\hline
\end{tabular}
\caption{The ranges of the branching ratios of \( \chi_{c1} \to \omega\omega, \phi\phi, \omega\phi \). Here, the experimental data of \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \) provide the central value with an error tolerance. The range in the \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \) is determined by the error existing in the experimental data. The experimental error also results in the range of \( \alpha \), i.e. \( \alpha = 1.14 \sim 1.28 \). By this determined range of \( \alpha \) and considering the range of mixing angle \( (\theta = (3.4 \pm 0.2)^\circ) \), we give the possible range of the branching ratios of \( \chi_{c1} \to \omega\omega, \phi\phi, \omega\phi \).}
\end{table}

\(^1\) In our calculation, we also tried to fit the measured \( B(\chi_{c1} \to K^{(*)0}\bar{K}^{(*)0}) \) by adopting the monopole form factor (setting \( N = 1 \) in Eq. (12)). Although we can also describe the \( \chi_{c1} \to K^{(*)0}\bar{K}^{(*)0} \) data, the obtained value of \( \alpha \) is far away from order of unity. Thus it is more reasonable not to take the monopole form factor in our calculation. Then, we choose the dipole form factor instead. In Ref. [24], Cheng et al. preferred the \( N = 1 \) monopole form factor, whereas, in our work we adopt \( N = 2 \). The difference between Ref. [24] and this work is due to that the intermediate states of the hadronic loop for the processes discussed in Ref. [24] can be on-shell while the intermediate states in this work are off-shell.
values may provide an ideal opportunity to test the $\omega-\phi$ mixing.

IV. DISCUSSION AND SUMMARY

To summarize, in this work, we study the mesonic loop contributions and the $\omega-\phi$ mixing effect to the branching ratios of $\chi_{c1} \to \omega\omega$, $\phi\phi$ and $\omega\phi$. From the results, one can note that the $\omega-\phi$ mixing plays an important role in the understanding of the clear signal for $\chi_{c1} \to \omega\omega$ observed in experiments. Our results also indicate that accurate measurements on the ratios $BR(\chi_{c1} \to \omega\omega)/BR(\chi_{c1} \to \omega\omega)$ and $BR(\chi_{c1} \to \omega\phi)/BR(\chi_{c1} \to \phi\phi)$ are very helpful for checking mesonic loop contributions and the $\omega-\phi$ mixing effect.

It is noted from the figures we presented in the text that the uncertainties in the theoretical computations originate from the errors of the data, therefore more accurate measurements are necessary for further studies. Fortunately, a large database on such rare decay modes will be available at BES-III, which will help to draw more solid conclusions.

Notes added: Very recently, a similar work [31] appeared in the arXiv submitted by X.H. Liu and Q. Zhao when this manuscript was close to completion. In Ref. [31], the authors calculated $\chi_{c1} \to \rho\rho$, $\omega\omega$, $\phi\phi$ are obtained. In our work, we mainly focused on $\chi_{c1} \to \omega\omega$, $\phi\phi$, $\omega\phi$ channels, where we also consider the hadronic loop effect. Our discussion is based on the recent preliminary results of $\chi_{c1} \to VV$ presented by the BES collaboration at the Hadron 2009 conference, especially the first observation of $\chi_{c1} \to \omega\phi$ [2]. Furthermore, we proposed that the $\omega-\phi$ mixing can be tested via $\chi_{c1} \to \phi\phi$, $\omega\omega$, $\omega\phi$ decays, which is different from the idea in Ref. [31] to some extent.

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![Fig. 3: (Color online). The contour plot for the dependence of $BR(\chi_{c1} \to \omega\omega)$, $BR(\chi_{c1} \to \phi\phi)$ and $BR(\chi_{c1} \to \omega\phi)$ on $\theta$ and $\alpha$. Here, the range sandwiched between two vertical solid lines is allowed by the 1-sigma standard deviation of the mixing angle $\theta$.](image-url)
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[1] X. Q. Li, X. Liu and Z. T. Wei, Front. Phys. China 4, 49 (2009) [arXiv:0808.2587 [hep-ph]].
[2] Y.P. Huang, for the BESIII Collaboration, talk given in Hadron 2009 conference, Florida State University, November 29-December 4 2009.
[3] V. L. Chernyak and A. R. Zhitnitsky, Nucl. Phys. B 201, 492 (1982) [Erratum-ibid. B 214, 547 (1983)].
[4] S. Okubo, Phys. Lett. 5 (1963) 165.
[5] G. Zweig, CERN Rep. 8419/TH-412.
[6] G. Zweig, CERN Rep. TH-401.
[7] J. Iizuka, Prog. Theor. Phys. Suppl. 37, 21 (1966).
[8] H. J. Lipkin, Phys. Lett. B 179, 278 (1986).
[9] X. Liu, X. Q. Zeng and X. Q. Li, Phys. Rev. D 74, 074003 (2006) [arXiv:hep-ph/0606191].
[10] X. Liu, B. Zhang and X. Q. Li, Phys. Lett. B 675, 441 (2009) [arXiv:0902.0480 [hep-ph]].
[11] Y. J. Zhang, G. Li and Q. Zhao, Phys. Rev. Lett. 102, 172001 (2009) [arXiv:0902.1300 [hep-ph]].
[12] F. K. Guo, C. Hanhart and Ulf-G. Meissner, Phys. Rev. Lett. 103, 082003 (2009) [arXiv:0907.0521 [hep-ph]].
[13] X. Liu, B. Zhang and S. L. Zhu, Phys. Lett. B 645, 185 (2007) [arXiv:0610278].
[14] X. Liu, B. Zhang and S. L. Zhu, Phys. Rev. D 77, 114021 (2008) [arXiv:0803.4270 [hep-ph]].
[15] X. Liu, Phys. Lett. B 680, 137 (2009) [arXiv:0904.0136 [hep-ph]].
[16] M. Benayoun, L. DelBuono, S. Eidelman, V. N. Ivanchenko and H. B. O’Connell, Phys. Rev. D 59, 114027 (1999) [arXiv:hep-ph/9902326].
[17] A. Kucharskian and U. G. Meissner, Mod. Phys. Lett. A 21, 1423 (2006) [arXiv:hep-ph/0603061].
[18] M. Benayoun, L. DelBuono and H. B. O’Connell, Eur. Phys. J. C 17, 593 (2000) [arXiv:hep-ph/9905350].
[19] M. Benayoun, L. DelBuono, P. Leruste and H. B. O’Connell, Eur. Phys. J. C 17, 303 (2000) [arXiv:nucl-th/0004005].
[20] M. Benayoun and H. B. O’Connell, Eur. Phys. J. C 22, 503 (2001) [arXiv:nucl-th/0107047].
[21] M. Benayoun, P. David, L. DelBuono, O. Leitner and H. B. O’Connell, Eur. Phys. J. C 55, 199 (2008) [arXiv:0711.4482 [hep-ph]].
[22] S. I. Dolinsky et al., Phys. Rept. 202, 99 (1991).
[23] H. Y. Cheng, C. Y. Cheung, G. L. Lin, Y. C. Lin, T. M. Yan and H. L. Yu, Phys. Rev. D 47, 1030 (1993) [arXiv:hep-ph/9209262].
[24] T. M. Yan, H. Y. Cheng, C. Y. Cheung, G. L. Lin, Y. C. Lin and H. L. Yu, Phys. Rev. D 46, 1148 (1992) [Erratum-ibid. D 55, 5851 (1997)].
[25] M. B. Wise, Phys. Rev. D 45, R2188 (1992).
[26] G. Burdman and J. F. Donoghue, Phys. Lett. B 280, 287 (1992).
[27] R. Casalbuoni, A. Deandrea, N. Di Bartolomeo, R. Gatto, F. Feruglio and G. Nardulli, Phys. Rept. 281, 145 (1997) [arXiv:hep-ph/9605342].
[28] P. Colangelo, F. De Fazio and T. N. Pham, Phys. Rev. D 69, 054023 (2004) [arXiv:hep-ph/0310084].
[29] H. Y. Cheng, C. K. Chua and A. Soni, Phys. Rev. D 71, 014030 (2005) [arXiv:hep-ph/0409317].
[30] M. Ablikim et al. [BES Collaboration], Phys. Rev. D 70, 092003 (2004) [arXiv:hep-ex/0408012].
[31] X. H. Liu and Q. Zhao, arXiv:0912.1508 [hep-ph].