Is there an isovector companion of the $X(2175)$?

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Abstract

In this letter we study the reaction $e^+e^- \rightarrow \phi\pi^0\eta$ with the $\pi^0\eta$ system in an isoscalar s-wave configuration. We use a formalism recently developed for the study of $e^+e^- \rightarrow \phi\pi\pi$ and $e^+e^- \rightarrow \phi K^+K^-$. The obtained cross section is within the reach of present $e^+e^-$ machines. Measuring this channel would test the absence of an isovector companion of the $X(2175)$ as predicted by the three body approach to this resonance.

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I. INTRODUCTION

Recently, a detailed study of the reactions \( e^+e^- \to \phi\pi\pi \) and \( e^+e^- \to \phi K^+K^- \) was performed for \( \sqrt{s} \geq 2 \text{ GeV} \) \cite{1,2}. These processes involve the \( \gamma\phi\pi\pi \) and \( \gamma\phi K^+K^- \) vertex functions respectively which are calculated at low photon virtuality and low di-meson invariant mass using Resonance Chiral Perturbation Theory (\( R\chi\PT \)). The dynamics is shown to be dominated by electromagnetic form factors and meson-meson amplitudes. This calculation, valid for low intermediate photon virtuality is improved replacing the lowest order terms in the form factors by the full form factors at the required \( \sqrt{s} \) as extracted from experiment or by the unitarized form factors. Likewise, the lowest order terms in the meson-meson amplitudes are replaced by the unitarized meson-meson amplitudes containing the scalar poles. The related process \( e^+e^- \to \phi f_0 \) has been also studied in the context of Nambu-Jona-Lasinio models \cite{3}.

In the \( e^+e^- \to \phi\pi\pi \) case, the calculation reproduces the background in the total cross section but not the peak around 2175 MeV, discovered by BaBar \cite{4} and confirmed by Bes in \( J/\Psi \to \eta\phi f_0(980) \) \cite{5} and Belle also in \( e^+e^- \to \phi\pi^+\pi^- \) \cite{6}.

The BaBar values for the mass and width of this new resonance, known as \( X(2175) \) or \( Y(2175) \), are \( M_X = 2175 \pm 10 \text{ MeV} \) and \( \Gamma_X = 58 \pm 16 \pm 20 \text{ MeV} \) \cite{4}, which are consistent with Bes, \( M_X = 2186 \pm 10 \pm 6 \text{ MeV} \) and \( \Gamma_X = 65 \pm 23 \pm 17 \text{ MeV} \) \cite{5}, and Belle, \( M_X = 2079 \pm 13^{+79}_{-28} \text{ MeV} \) and \( \Gamma_X = 192 \pm 23^{+25}_{-61} \text{ MeV} \) \cite{6} results. As pointed out in a recent review \cite{7}, the narrow width is at odds with the large decay width into two mesons predicted in models for \( \bar{s}s \) states \cite{8}, diquark-anti-diquark states \cite{9}, tetraquark \( s\bar{s}s\bar{s} \) states \cite{10} or gluon hybrids \( s\bar{s}g \) \cite{11} (see also \cite{12}).

Another proposal for the nature of \( X(2175) \), which turns out to be consistent with the observed mass and width, is a three-body \( \phi \bar{K}K \) system \cite{13}. The result of this framework is a neat resonance peak around a total mass of 2150 MeV and an invariant mass for the \( K\bar{K} \) system around 970 MeV, quite close to the \( f_0(980) \) mass. The state appears in the isospin \( I = 0 \) sector, and qualifies as a \( \phi f_0(980) \) resonance. Interestingly, the theory also shows that there is no resonance in the isovector channel. This is the main motivation to study the reaction \( e^+e^- \to \phi\pi^0\eta \), whose final state is in a pure isovector state.
II. CALCULATION OF $e^+e^- \rightarrow \phi\pi^0\eta$

In this work we apply the same formalism as in [1, 2] to the production of $\phi\pi^0\eta$. We start form the $R\chi PT$ Lagrangian and follow the conventions in [14]. The relevant interactions in their notation are

\[
\mathcal{L} = \mathcal{L}^{(2)} + \mathcal{L}^{(F)} + \mathcal{L}^{(G)}
\]

\[
\mathcal{L}^{(2)} = \frac{1}{4} f^2 tr \left( (D_\mu U)^\dagger D^\mu U + \chi U^\dagger + \chi^\dagger U \right)
\]

\[
\mathcal{L}^{(F)} = \frac{F_V}{2\sqrt{2}} tr(V_{\mu\nu} f_{\mu\nu}^+) \quad (2)
\]

\[
\mathcal{L}^{(G)} = \frac{i G_V}{\sqrt{2}} tr(V_{\mu\nu} u^\mu u^\nu) \quad (3)
\]

where

\[
u_\mu = i u^\dagger D_\mu U u^\dagger, \quad U = u^2, \quad u = e^{-i\sqrt{2}\phi}, \quad \Phi = \frac{1}{\sqrt{2}} \lambda_i \varphi_i \quad (5)
\]

\[
f_{\mu\nu}^+ = u F_{L}^{\mu\nu} u^\dagger + u^\dagger F_{R}^{\mu\nu} u, \quad D_\mu U = \partial_\mu U - i [v_\mu, U]. \quad (6)
\]

We introduce the photon field through $v_\mu = eQ A_\mu$ and $F_{L}^{\mu\nu} = F_{R}^{\mu\nu} = eQ F^{\mu\nu}$ ($e > 0$) where $F^{\mu\nu}$ denotes the electromagnetic strength tensor.

There are no tree-level contributions to $e^+e^- \rightarrow \phi\pi^0\eta$ in $R\chi PT$. At one loop level, this process is induced by the diagrams shown in Fig. (1) where the fermionic lines are shown only in the first diagram and particles in the loops are neutral and charged kaons. For the sake of simplicity a shaded circle and a dark circle account for the diagrams i) plus j) and k) plus l) respectively, which differentiate the direct photon coupling from the coupling through an intermediate vector meson. We will address the corresponding diagrams as a), b), when we have the direct photon coupling and a'), b'), when the coupling goes through the exchange of a vector meson.

At the energy of the reaction, $\sqrt{s} \geq 1.7$GeV, it is quite probable to excite higher states. Here, we consider the excitation of virtual $K^*K$ states and their contribution to $e^+e^- \rightarrow \phi\pi^0\eta$ through the chain $e^+e^- \rightarrow K^*\overline{K} \rightarrow \phi K \overline{K}$ with the rescattering of the kaon pair to $\pi^0\eta$ as shown in Fig. (2).

The $VV'P$ interaction in this diagram is dictated by the anomalous chiral Lagrangian [15] which we rewrite in terms of the tensor field as

\[
\mathcal{L}_{\text{anom}} = \frac{G}{\sqrt{2}} \epsilon_{\mu\nu\alpha\beta} tr(\partial^\mu V^\nu \partial^\alpha V^\beta \Phi) = \frac{G_T}{4\sqrt{2}} \epsilon_{\mu\nu\alpha\beta} tr(V^{\mu\nu} V^{\alpha\beta} \Phi). \quad (7)
\]
FIG. 1: Feynman diagrams for $e^+e^- \rightarrow \phi\pi^0\eta$ in $R\chi PT$. See [1] for the details and conventions.

![Feynman Diagrams](image1.png)

FIG. 2: Feynman diagram for $e^+e^- \rightarrow K^*\bar{K} \rightarrow \phi K\bar{K} \rightarrow \phi\pi^0\eta$.

The calculation of the invariant amplitude is similar to the previously studied $\phi\pi\pi$ final state and we refer the reader to Ref. [1] for the details. Both $R\chi PT$ contributions in Fig. (11) and the contributions of Fig. (2) turn out to be divergent. However, it is shown in

with $G_T = M_V M_{V'} G$. This Lagrangian was obtained in [15]. In this formalism no direct $VP\gamma$ coupling emerges and this vertex is generated by the Lagrangian in Eq. (7) in combination with the $V\gamma$ interaction in Eq. (3) (see the discussion above Eq. (56) of [15]).
that divergences of $R\chi PT$ contributions come from diagram c) of Fig. 1 and involve the same divergent scalar integral which appears in the unitarization of the meson-meson amplitudes. On the other side, a careful decomposition of the loop integrals of Fig. 2 into scalar integrals was performed in the appendix of Ref. [1]. Using dimensional regularization it is shown there that the only divergent piece corresponds to the same scalar integral appearing in meson-meson scattering. The substraction constant required for this integral was discussed in [16, 17] and we briefly review it below.

The total amplitude for $e^+(p^+) e^-(p^-) \rightarrow \phi(Q, \epsilon) \pi^0(p) \eta(p')$ arising from the diagrams in Figs. (1,2) is

\[ -iM = \frac{ie^2}{2\pi^2 m_K^2} \frac{V_{KK\pi\eta}^{\mu}}{\sqrt{2}} \frac{L_{\mu\nu}^{(1)} - J_{\mu\nu}^{(2)}}{k^2} \epsilon^\nu \]

where $k^2 = (p^+ + p^-)^2$, $V_{KK\pi\eta}$ denotes the leading order on-shell amplitude for $K\bar{K} - \pi^0\eta$ scattering, $L^{\mu} \equiv \bar{\tau}(p^+) \gamma^\mu u(p^-)$ and the Lorentz covariant structures are given by

\[ L_{\mu\nu}^{(1)} \equiv Q \cdot k g_{\mu\nu} - Q_{\mu} k_{\nu}, \quad L_{\mu\nu}^{(2)} = k^2 g_{\mu\nu} - k_{\mu} k_{\nu}. \]  

The $I, J$ functions entering Eq. (8) are given by

\[ I = \frac{\sqrt{2} M_\phi}{2 f^2} G_V F_K^{(1)}(k^2) I_P - \frac{G}{4\sqrt{2}} F_K^{(1)}(k^2) \times \]

\[ \left\{ Q \cdot k J_V - m_K^2 \left\{ I_G - I_2 + 4 - \frac{1}{2} [a(\mu) + 3] + \frac{1}{2} \log \frac{m_K^2}{\mu^2} \right\} \right\} \]  

\[ J = \frac{\sqrt{2} M_\phi}{2 f^2} G_V F_K^{(1)}(k^2) \left[ J_P + \frac{(4\pi m_K)^2}{4k^2} G_K K^2 G_{\pi\eta} \right] \]

\[ -\frac{GM_\phi}{4\sqrt{2}} F_K^{(1)}(k^2) J_V, \]  

with the divergent loop integral

\[ G_K K^2 G_{\pi\eta} = \mu^{2\epsilon} \int \frac{d^d l}{(2\pi)^d} \frac{i}{\square_K (l + Q) \square_K (l - k)} \]

\[ = \frac{1}{(4\pi)^2} \left\{ a(\mu) + \log \frac{m_K^2}{\mu^2} + I_G(m_{\pi\eta}^2) \right\}, \]

where $\mu$ is the dimensional regularization scale, $\square_i(l) = l^2 - m_i^2$ and

\[ I_G(p^2) = \int_0^1 dx \log \left[ 1 - \frac{p^2}{m_K^2} x(1 - x) - i\epsilon \right] = -2 + \sigma(p^2) \log \frac{\sigma(p^2) + 1}{\sigma(p^2) - 1}. \]  


Here $\sigma(p^2) = \sqrt{1 - \frac{4m^2}{p^2}}$ and the integrals in Eqs. (10, 11) are given by

$$I_P = \int_0^1 dx \int_0^x dy \frac{y(1-x)}{f_1(x,y)}, \quad J_P = \frac{1}{2} \int_0^1 dx \int_0^x dy \frac{y(1-2y)}{f_1(x,y)},$$

$$J_V = \int_0^1 dx \int_0^x dy \frac{y(1-x)}{f_2(x,y)}, \quad I_2 = \int_0^1 dx \int_0^x dy \log|f_2(x,y)|,$$

$$f_1(x,y) = 1 - \frac{Q^2}{m_K^2} x(1-x) + \frac{2Q \cdot k}{m_K^2} (1-x)y - \frac{k^2}{m_K^2} y(1-y) - i\varepsilon,$$

$$f_2(x,y) = f_1(x,y) - \frac{(m_K^2 - m_K^2)}{m_K^2}(y - x). \quad (14)$$

The divergences in the loop integrals are contained in $a(\mu)$ in Eqs. (10, 12) and we discuss its physical value in detail below.

From Eqs. (8, 11) we can see that the dynamics is dominated by two main effects in the $\gamma^* \phi \pi^0 \eta$ vertex function which occur at two different energy scales: the leading order terms for the $K^*K - \pi^0 \eta$ on-shell scattering amplitude at the di-meson invariant mass scale, $m_{\pi\eta}$, and the electromagnetic meson form factors at the energy of the reaction, $\sqrt{s}$. The calculation of the $\gamma^* \phi \pi^0 \eta$ vertex function involved in this amplitude is strictly accurate for low di-meson invariant mass and low virtuality of the photon. We improve this results in two respects: i) we replace the leading order result for the $K^*K - \pi^0 \eta$ on-shell scattering amplitude by the unitarized amplitude containing the scalar poles and ii) we replace the leading order terms in the kaon form factor by the unitarized one and, following [2], the leading order terms of the $K^*K$ transition form factor are replaced by the complete form factor at the energy of the reaction as extracted from data.

The isovector $s$-wave $K^*K - \pi^0 \eta$ unitarized scattering amplitude denoted $t_{K^*K\pi\eta}^0$ gives a successful description of the corresponding cross section up to $m_{\pi\eta} \approx 1\text{GeV}$ and is based on the imposition of unitary constraints in coupled channels of $\chi$PT. Following [16, 17], unitarization reduces to the solution of the Bethe-Salpeter equation

$$T = T_{(2)} + T_{(2)} \cdot G \cdot T, \quad (15)$$

where $T$ is the matrix of unitarized amplitudes of the desired isospin channel

$$T = \begin{pmatrix} t_{KKK}^0 & t_{KK\pi\eta}^0 \\ t_{KK\pi\eta}^0 & t_{\pi\eta\pi\eta}^0 \end{pmatrix}. \quad (16)$$

The elements of the matrix $T_{(2)}$ are the corresponding on-shell amplitudes calculated in $\chi$PT.
at $O(p^2)$. A straightforward calculation yields

$$
\mathcal{T}_\mathcal{M}(2) = \begin{pmatrix}
V_{KKK} & V_{K\pi\eta} \\
V_{K\pi\eta} & V_{\pi\eta\eta}
\end{pmatrix}
= \begin{pmatrix}
\frac{m_{\pi\eta}^2}{4f^2} & \frac{9m_{\pi\eta}^2 - 8m_K^2 - 3m_\pi^2 - m_\eta^2}{6\sqrt{6}f^2} \\
\frac{9m_{\pi\eta}^2 - 8m_K^2 - 3m_\pi^2 - m_\eta^2}{6\sqrt{6}f^2} & -\frac{m_\pi^2}{3f^2}
\end{pmatrix}. 
$$

(17)

The diagonal matrix $\mathcal{G}$ contains the loop integrals

$$
\mathcal{G} = \begin{pmatrix}
G_{KK}(m_{\pi\eta}^2) & 0 \\
0 & G_{\pi\eta}(m_{\pi\eta}^2)
\end{pmatrix}
$$

(18)

with $G_{KK}(m_{\pi\eta}^2)$ given in Eq. (12) and

$$
G_{\pi\eta}(m_{\pi\eta}^2) = \mu^2 \epsilon \int \frac{d^d l}{(2\pi)^d} \frac{i}{\Box_\pi (l + Q) \Box_\eta (l - k)}
= \frac{1}{(4\pi)^2} \left[ a(\mu) + \log \frac{m_\eta^2}{\mu^2} + I_\eta(m_{\pi\eta}^2) \right]
$$

(19)

with

$$
I_\eta = -2 + \frac{m_{\pi\eta}^2 - m_\eta^2 + m_\pi^2}{2m_{\pi\eta}^2} \log \left( \frac{m_{\pi}^2}{m_\eta^2} \right)
+ \frac{\nu}{2m_{\pi\eta}^2} \left\{ \log \left[ \left( \frac{m_{\pi\eta}^2 + \nu}{m_\pi^2} \right)^2 - \left( m_\eta^2 - m_\pi^2 \right)^2 \right] - 2\pi i \right\}
$$

(20)

and

$$
\nu = \sqrt{[m_{\pi\eta}^2 - (m_\eta - m_\pi)^2] [m_{\pi\eta}^2 - (m_\eta + m_\pi)^2]}.
$$

(21)

The subtraction constant has been fixed in analogy with Ref. [17] to $a(\mu_0) = 0.87$ for $\mu_0 = 1.2$ GeV matching Eq. (19) to the cutoff regularized integral with a cutoff parameter $\Lambda = 1$ GeV. It is related at different scales as $a(\mu) = a(\mu_0) + \log \frac{\mu^2}{\mu_0^2}$, and therefore the loop function is scale independent.

The unitarized amplitudes are obtained solving the algebraic equation (15). In particular, it was shown in [16] that $t^0_{KK\pi\eta}$ has a pole corresponding to the $a_0(980)$ resonance which manifests as a peak in the squared amplitude as shown in Fig. (3).

Concerning the meson electromagnetic form factors, the high virtuality of the intermediate photon involved in our process requires to work out the complete $\gamma K\bar{K}$ vertex functions. The calculation of the isovector kaon form factor ($F_{K}^{(1)}$) has been done in the context of
FIG. 3: Square modulus of the unitarized amplitude $t_{KK\pi\eta'}^0$.

$U\chi PT$ and the result can be found in Eq.(31) of Ref. [18] (see the discussion below this equation). On the other hand, the extraction of the isovector $K^*K$ transition form factor $(F_{K^*K}^{(1)})$ follows from [2], where an appropriate characterization of this form factor at the energy of the reaction is done. As this form factor is physically dominated by resonances, in [2] it is described with the lowest order terms obtained in $R\chi PT$ complemented with the exchange of a $\rho'$

$$F_{K^*K}^{(1)}(k^2) = \frac{F_V G}{3} \left( \frac{3m_\rho}{k^2 - m_\rho^2} \right) + b_1 \left( \frac{3m_{\rho'}^2}{k^2 - m_{\rho'}^2 + i\sqrt{s}\Gamma_{\rho'}(s)} \right),$$

with the energy dependent width used in [19]

$$\Gamma_{\rho'}(s) = \Gamma_{\rho'} \left[ \frac{P_{4\pi}(s)}{P_{4\pi}(m_{\rho'}^2)} B_{4\pi}^{\rho'} + (1 - B_{4\pi}^{\rho'}) \right],$$

where

$$P_{4\pi}(s) = \frac{(s - 16m_\pi^2)^{3/2}}{s}. \tag{24}$$

The values for the constants appearing here are extracted from the central values of Table XV in [19] as $B_{4\pi}^{\rho'} = 0.65$, $m_{\rho'} = 1504$ MeV and $\Gamma_{\rho'} = 438$ MeV (see also [20]). The parameter $b_1$ is then fitted to the experimental data for the isovector cross sections reported for $e^+e^- \rightarrow K^*K$ at $\sqrt{s} = 1400 - 3000$ MeV in table VII of [19]. The fit to the isovector cross section within 1$\sigma$ yields $b_1 = -(0.255^{+0.030}_{-0.040}) \times 10^{-3}$ MeV$^{-1}$.

We use these results to calculate the $\pi^0\eta$ spectrum, obtaining
FIG. 4: Differential cross section as a function of the $\pi^0\eta$ invariant mass and of the center of mass energy.

\[
\frac{d\sigma}{dm_{\pi\eta}} = \frac{\alpha^2}{24\pi^5m_K^4} \frac{|Q|\tilde{p}|}{s^{\frac{3}{2}}} \frac{|\mathfrak{t}_{K\pi\pi}^0|^2}{2} h(s, m_{\pi\eta}) \tag{25}
\]

where

\[
h(s, m_{\pi\eta}) = |I|^2 \left( M_\phi^2 + 2\omega^2 \right) - 6\text{Re}(IJ^*) \sqrt{s} \omega + |J|^2 \frac{s}{M_\phi^2} \left( 2M_\phi^2 + \omega^2 \right). \tag{26}
\]

Here $(\omega, Q)$ stands for the momentum of the $\phi$ in the center of momentum system of the reaction and $\tilde{p}$ denotes the momentum of the final pion in the $\pi^0\eta$ center of momentum system

\[
\omega = \frac{s + M_\phi^2 - m_{\pi\eta}^2}{2\sqrt{s}}, \quad |Q| = \frac{\lambda\frac{1}{2}(s, M_\phi^2, m_{\pi\eta}^2)}{2\sqrt{s}}, \quad |\tilde{p}| = \frac{\lambda\frac{1}{2}(m_{\pi\eta}^2, m_{\pi}^2, m_\eta^2)}{2m_{\pi\eta}}. \tag{27}
\]

with

\[
\lambda(m_1^2, m_2^2, m_3^2) = (m_1^2 - (m_2 - m_3)^2)(m_1^2 - (m_2 + m_3)^2). \tag{28}
\]

III. NUMERICAL RESULTS

We evaluate numerically the integrals and the differential cross section. Using the physical masses and coupling constants $m_K = 495$ MeV, $m_\phi = 1019.4$ MeV, $\alpha = 1/137$, $G_V =$
FIG. 5: Cross section for $e^+e^- \rightarrow \phi \left[ \pi^0\eta \right]_{I=1,J=0}$ as a function of $\sqrt{s}$. The dashed lines correspond to the 1σ region in the extraction of the isovector $K^*K$ transition form factor in Ref. [2], i.e. to the fit of $b_1$ in Eq. (22).

$53 \text{ MeV}$, $F_V = 154 \text{ MeV}$, $f_\pi = 93 \text{ MeV}$, and $G = 0.016 \text{ MeV}^{-1}$ in Eq. (25) we obtain the spectrum shown in Fig. (4) where the presence of the $a_0(980)$ is well visible. This is a consequence of the fact that the $a_0(980)$ poles are well reproduced in the unitarization of meson-meson $s$-wave isovector amplitudes present in our calculation. Next we integrate $m_{\pi\eta}$ in the $a_0(980)$ region, $m_{\pi\eta} = 850 - 1100 \text{ MeV}$. The obtained cross section is shown in Fig. (5).

The mechanisms studied here were shown to be responsible for the production of most of the events in the case of the $\phi\pi\pi$ final state except for the resonant events due to the $X(2175)$ [1]. On the other hand, in the case of the $\phi K^+K^-$ final state studied in [2], within the limitations due to the extraction of the $K^*K$ isovector transition form factor from data, a good description of experimental points is obtained but there seems to be room for additional contributions around 2200 MeV. In this case, the $\phi K^+K^-$ system can be in both isoscalar and isovector states hence the natural candidate for additional contributions is an isovector companion of the $X(2175)$. Interestingly, in the three-body description of the $X(2175)$ proposed in [13], no peak is generated in the isovector channel. Our calculation of the pure isovector channel $e^+e^- \rightarrow \phi a_0(980)$ under the mechanisms studied in [1] yields a cross section of the same size as that of $e^+e^- \rightarrow \phi f_0$ and thus within the reach of present $e^+e^-$ machines. If there is no such a thing as an isovector companion...
of the $X(2175)$ as expected from [13], then our result is a concrete theoretical prediction for the $e^+e^- \to \phi a_0(980)$ cross section; otherwise, the isovector companion must contribute to this process yielding valuable information on the nature of mesons at this energy. Hence we encourage experimentalists to measure the $e^+e^- \to \phi \pi^0\eta$ channel.

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[1] M. Napsuciale, E. Oset, K. Sasaki and C. A. Vaquera-Araujo, Phys. Rev. D 76, 074012 (2007) [arXiv:0706.2972 [hep-ph]].

[2] S. Gomez-Avila, M. Napsuciale and E. Oset, Phys. Rev. D 79, 034018 (2009) [arXiv:0711.4147 [hep-ph]].

[3] Yu. M. Bystritskiy, M. K. Volkov, E. A. Kuraev, E. Bartos and M. Secansky, Phys. Rev. D 77, 054008 (2008) [arXiv:0712.0304 [hep-ph]].

[4] B. Aubert et al., Phys. Rev. D 74 (2006) 091103, B. Aubert et al., Phys. Rev. D 76 (2007) 012008.

[5] M. Ablikim et al. [BES Collaboration], Phys. Rev. Lett. 100, 102003 (2008) [arXiv:0712.1143 [hep-ex]].

[6] C. P. Shen et al., Phys. Rev. D 80, 031101 (2009).

[7] S. L. Zhu, Int. J. Mod. Phys. E 17, 283 (2008) [arXiv:hep-ph/0703225].

[8] T. Barnes, N. Black and P. R. Page, Phys. Rev. D 68, 054014 (2003) [arXiv:nucl-th/0208072].

[9] G. J. Ding and M. L. Yan, Phys. Lett. B 643 (2006) 33.

[10] Z. G. Wang, Nucl. Phys. A 791 (2007) 106.

[11] G. J. Ding and M. L. Yan, Phys. Lett. B 657 (2007) 49.

[12] F. E. Close, In the Proceedings of 5th Flavor Physics and CP Violation Conference (FPCP 2007), Bled, Slovenia, 12-16 May 2007, pp 020[arXiv:0706.2709 [hep-ph]].

[13] A. Martinez Torres, K. P. Khemchandani, L. S. Geng, M. Napsuciale and E. Oset, Phys. Rev. D 78, 074031 (2008) [arXiv:0801.3635 [nucl-th]].

[14] G. Ecker et al. Nucl. Phys. B 321, 311 (1989).

[15] J. Bijnens, A. Bramon and F. Cornet, Z. Phys. C 46, 599 (1990).
[16] J. A. Oller and E. Oset, Nucl. Phys. A 620, 438 (1997) [Erratum-ibid. A 652, 407 (1999)]
   \texttt{arXiv:hep-ph/9702314}.

[17] J. A. Oller, E. Oset and J. R. Pelaez, Phys. Rev. Lett. 80, 3452 (1998) \texttt{arXiv:hep-ph/9803242};
    J. A. Oller, E. Oset and J. R. Pelaez, Phys. Rev. D 59, 074001 (1999) [Erratum-ibid. D 60,
    099906 (1999)] \texttt{arXiv:hep-ph/9804209}.

[18] J. A. Oller, E. Oset and J. E. Palomar, Phys. Rev. D 63, 114009 (2001)
    \texttt{arXiv:hep-ph/0011096}.

[19] B. Aubert et al. [BaBar Collaboration], Phys. Rev. D 77, 092002 (2008) \texttt{arXiv:0710.4454}
    [hep-ex].

[20] D. Bisello et al., Z. Phys. C 52, 227 (1991); F. Mane, D. Bisello, J. C. Bizot, J. Buon,
    A. Cordier and B. Delcourt, Phys. Lett. B 112, 178 (1982).