I. INTRODUCTION

The $f_2(1270)$ is a resonance that plays a prominent role in many reactions. For instance in $\pi^0\pi^0 \rightarrow \gamma\gamma$ it shows with a strength much bigger than the corresponding one of the $f_0(980)$ \cite{1, 2}. From the theoretical point of view, the $f_2(1270)$ has been widely accepted as an ordinary quark state belonging to a $P$-wave nonet of tensor mesons \cite{3, 4}. However, this wisdom was challenged by the work of Ref. \cite{5} where the interaction of vector mesons was studied and the $f_2(1270)$ emerged as one of the dynamically generated states from the $\rho$-$\rho$ interaction. The vector-vector interaction is described by means of the local hidden gauge Lagrangians \cite{6, 7} and the unitarization of this interaction leads to vector-vector scattering amplitudes that show poles for some states of spin $J = 0, 1, 2$ in $L = 0$. One of them is the $f_2(1270)$. The generalization to $SU(3)$ of the $\rho$-$\rho$ interaction was done following the same lines but with coupled channels \cite{9} and the $f_2(1270)$ appeared as a dynamically generated state of the $\rho$-$\rho$ interaction with parallel spins, corroborating the findings of Ref. \cite{5}. The reason for the large binding of the $f_2(1270)$ is a very large interaction of the $\rho$-$\rho$ in $J = 2$, much larger than for other channels, and larger than other interactions known in hadron physics.

The molecular nature of the $f_2(1270)$ claimed in Refs. \cite{5, 9} has been successfully tested in a large number of processes. In Ref. \cite{9} it was shown that it leads to a very good description of the decay rate for $f_2(1270) \rightarrow \gamma\gamma$. It was also shown in Ref. \cite{9} that the decay rates for two photons and one photon-one vector decays of the $f_0(1370)$, $f_2(1270)$, $f_0(1710)$, $f_2'(1525)$ and $K_2^*(1430)$ were in good agreement with available data. A study of the $J/\psi \rightarrow \phi(\omega)f_2(1270)$, $f_2'(1525)$ and $J/\psi \rightarrow K^{*0}(892)K_2^{*0}(1430)$ decays was also done within this picture in Ref. \cite{13} and good results were obtained. Another test conducted along these lines was the radiative decay of $J/\psi$ into $f_2(1270)$, $f_2'(1525)$, $f_0(1370)$ and $f_0(1710)$ which was done in Ref. \cite{14} and good results were obtained compared with the available experimental information. Similarly, predictions for $\psi(2S)$ decay into $\omega(\phi)f_2(1270)$, $\omega(\phi)f_2'(1525)$, $K^{*0}(892)K_2^{*0}(1430)$ and radiative decay of $\Upsilon(1S, 2S, 2S')$ into $f_2(1270)$, $f_2'(1525)$, $\gamma f_0(1370)$, $\gamma f_0(1710)$ were done in Ref. \cite{15} with also good agreement with available experiments. More recently it was also shown in Ref. \cite{16} that the ratio of the decay widths of $B^0_s \rightarrow J/\psi f_2(1270)$ to $B^0 \rightarrow J/\psi f_2'(1525)$ is compatible with the experimental values of \cite{17}. The nature of these states as vector meson-vector meson composite states has undergone a large number of test than any other model. Yet, all the test have been done in the mesonic sector and not in the baryonic sector. The recent measurement of the photoproduction of $f_2(1270)$ in Jefferson Lab \cite{18} offers us the first opportunity to do this new test, which we conduct here.

On the theoretical side there is a study of this reaction in Ref. \cite{19}, where the idea is to create the $f_2(1270)$ as the final state interaction of two pions in $D$-wave. Hence, one constructs a mechanism for two pion production and then lets the pions interact in $D$-wave to produce the $f_2(1270)$ resonance. Our picture has some similarity with this idea in the sense that we also generate the resonance from the final state interaction of two mesons, but these two mesons are two $\rho$ instead of two pions. Hence, the mechanism consists on the production of two $\rho$ mesons, for which vector meson dominance is used, and then the two $\rho$ mesons, upon interaction, produce the $f_2(1270)$. The results obtained are in good agreement with experi-
ment, both for the invariant mass dependence and the $t$ dependence, offering new support for the picture of the $f_2(1270)$ as a $\rho\rho$ molecular state.

II. FORMALISM

In Ref. [5] the $\rho\rho$ amplitudes in $L = 0$ were classified in terms of the spin of the system, which was produced by the $\rho$ polarization. Spin-projector operators were written in terms of the $\rho$ polarization vectors and concretely, for spin $S = 2$, the projector was given by

$$P^{(2)} = \frac{1}{2} \left( \epsilon_i^{(1)} \epsilon_j^{(2)} + \epsilon_i^{(2)} \epsilon_j^{(1)} \right) - \frac{1}{3} \epsilon_i^{(1)} \epsilon_i^{(2)} \delta_{ij},$$

where $\epsilon_i^{(k)}$ are the three spatial components ($i$) of the vector polarization of the $\rho$ meson $k$, in the order of $\rho(1) + \rho(2) \rightarrow \rho(3) + \rho(4)$. The momenta of the external $\rho$ mesons is assumed to be small with respect to the mass of the $\rho$, such that the time component of the $\epsilon^\mu$ ($\epsilon^0$) is neglected. This is the case of the polarization of the photons, where we only have transverse components. It was also discussed in Ref. [20] that extra terms linear in the momentum of the particles were small in this type of processes, justifying the success of the radiative decay, in spite of the photons not having small momenta.

As we can see, the $P^{(2)}$ projector is factorized into one block corresponding to the initial vectors and another one corresponding to the final vectors. The amplitude close to a pole that represents a resonance is then written as

$$t_{\text{pole}} \simeq \frac{g_T^2 P^{(2)}_{\text{initial}} P^{(2)}_{\text{final}}}{s - s_R},$$

with $s_R$ the pole position and $g_T$ the coupling of the resonance to the $\rho\rho$ component in isospin $I = 0$ and spin $S = 2$. Eq. (2) is a representation of a resonance amplitude, for instance the $f_2(1270)$ in the present case, as shown in Fig. 1 (a).

$$P^{(2)}_{\text{initial}} = \frac{1}{2} \left( \epsilon_i^{(1)} \epsilon_j^{(2)} + \epsilon_i^{(2)} \epsilon_j^{(1)} \right) - \frac{1}{3} \epsilon_i^{(1)} \epsilon_i^{(2)} \delta_{ij},$$

where $s_R$ is the pole position and $g_T$ the coupling of the resonance to the $\rho\rho$ component in isospin $I = 0$ and spin $S = 2$. Eq. (2) is a representation of a resonance amplitude, for instance the $f_2(1270)$ in the present case, as shown in Fig. 1 (a).

$$P^{(2)}_{\text{final}} = \frac{1}{2} \left( \epsilon_i^{(3)} \epsilon_j^{(4)} + \epsilon_i^{(4)} \epsilon_j^{(3)} \right) - \frac{1}{3} \epsilon_i^{(3)} \epsilon_i^{(4)} \delta_{ij},$$

where $s_R$ is the pole position and $g_T$ the coupling of the resonance to the $\rho\rho$ component in isospin $I = 0$ and spin $S = 2$. Eq. (2) is a representation of a resonance amplitude, for instance the $f_2(1270)$ in the present case, as shown in Fig. 1 (a).

$$t_{\text{pole}} \simeq \frac{g_T^2 P^{(2)}_{\text{initial}} P^{(2)}_{\text{final}}}{s - s_R},$$

with $s_R$ the pole position and $g_T$ the coupling of the resonance to the $\rho\rho$ component in isospin $I = 0$ and spin $S = 2$. Eq. (2) is a representation of a resonance amplitude, for instance the $f_2(1270)$ in the present case, as shown in Fig. 1 (a).

As we can see, the photon gets converted into one $\rho^0$, a characteristic of the local hidden gauge Lagrangian, and the other $\rho^0$ becomes a virtual $\rho$ connecting to the proton. Apart from the vertex of Eq. (5), we need now two ingredients, the $\gamma\rho\rho$ conversion vertex and the $\rho NN$ vertex. The $\gamma\rho\rho$ conversion vertex is easily obtained from the local hidden gauge Lagrangian [6-10] (see Ref. [21] for a practical set of rules) and we have [11]

$$-it_{\rho\gamma} = -i \frac{eM^2}{\sqrt{2} g} \epsilon_\mu (\rho) e^\mu(\gamma),$$

with

$$g = \frac{M_\rho}{2f}; \quad f = 93 \text{ MeV}; \quad \frac{e^2}{4\pi} = \frac{1}{137}.$$ (7)

The $\rho^0$ polarization of Eq. (4) is contracted with the one of the $\rho$ in the vertex of Eq. (5) and the polarization vector of the $\rho$ in Eq. (4) gets converted into the one of the photon. The other ingredient that we need is the vector-nucleon-nucleon vertex, which is given by the Lagrangian

$$L_{BBV} = \frac{g}{2} \left( \bar{B} \gamma^\mu [\gamma^\nu, B] > + < \bar{B} \gamma^\nu B > \gamma^\mu V_\mu > \right),$$

which provides a $\rho^0 pp$ vertex

$$-it_{\rho^0 pp} = \frac{g}{\sqrt{2}} p^\mu p_\mu (\rho^0).$$
There is one more consideration to make. The amplitude of Eq. (2) is evaluated for a $\rho\rho$ state in the unitary normalization, which for $I = 0$ is given by (recall $|\rho^+> = -|11>$)

$$|\rho\rho, I = 0> = -\frac{1}{\sqrt{6}}(\rho^+\rho^- + \rho^-\rho^+ + \rho^0\rho^0).$$

(10)

This has a factor $\frac{1}{\sqrt{2}}$ extra with respect to the good normalization and is taken to account for the identity of the particles in the intermediate states. The good amplitude will have $g_{\rho\rho}$ of Eqs. (2) and (5) multiplied by $\sqrt{2}$. In addition, in order to project over $\rho^0\rho^0$ we must multiply the coupling to the $I = 0$ state by $\frac{1}{\sqrt{3}}$. Altogether, the coupling $\tilde{g}_{T}$ to be used for $\rho^0\rho^0$ is

$$\tilde{g}_{T} = -\frac{\sqrt{2}}{\sqrt{3}}g_{T}.$$  

(11)

Considering the vertices described above, the $T$ matrix for the diagram of Fig. 2 is given by

$$-iT_{\gamma p\rightarrow f_{2}(1270)p} = -\frac{\tilde{g}_{T}}{2} \left\{ \frac{1}{2} \epsilon_{i}(\gamma)\epsilon_{j}(\rho) + \epsilon_{j}(\gamma)\epsilon_{i}(\rho) \right\}$$

$$-\frac{1}{3}\epsilon_{m}(\gamma)\epsilon_{m}(\rho)\delta_{ij} \frac{1}{q^2 - M^2} < M|\gamma\rho|\mu_{i}\mu_{j}|M > ,$$

(12)

with $M$ and $M'$ the spin third component of the initial and final proton. We must perform the sum over the polarizations of the $\rho$ meson exchanged in Fig. 2 and then we get

$$T_{\gamma p\rightarrow f_{2}(1270)p} = \frac{\tilde{g}_{T}}{2} \left\{ \frac{1}{2} \epsilon_{i}(\gamma)(-g_{j\mu} + \frac{q_{j}q_{\mu}}{M_{\rho}^2})$$

$$+\frac{1}{2}\epsilon_{j}(\gamma)(-g_{i\mu} + \frac{q_{i}q_{\mu}}{M_{\rho}^2}) - \frac{1}{3}\epsilon_{m}(\gamma)\epsilon_{m}(\rho)\delta_{ij} (-g_{m\mu} + \frac{q_{m}q_{\mu}}{M_{\rho}^2}) \right\}$$

$$< M|\gamma\rho|\mu_{i}\mu_{j}|M > .$$

(13)

In Eqs. (12) and (13) all the components in $g_{\mu\nu}$, etc., $q_{i}$, $q_{j}$, $q_{\mu}$ are covariant. The latin indices run over 1, 2, 3 and the $\mu$ index from 0, 1, 2, 3.

In order to calculate $\sum \sum |T|^2$, we calculate

$$\frac{1}{2} \sum_{M',M} < M'|\gamma\rho|M > < M|\gamma\rho|M >$$

$$= \frac{1}{2} \sum_{r,r'} < \bar{u}_{r'}(p')\gamma^{\mu}u_{r}(p)\bar{u}_{r}(p)\gamma^{\mu'}u_{r'}(p')$$

$$= \frac{1}{2} Tr\left[ \frac{\not{p} + m_{\rho}\gamma^{\mu}}{2m_{\rho}} \frac{\not{p} + m_{\rho}\gamma^{\mu'}}{2m_{\rho}} \right]$$

$$= \frac{1}{2m_{\rho}} (2p^\mu p'^\mu + q^\mu p'^\mu + q^\mu' p'^\mu - p\cdot qg^{\mu\mu'}).$$

(14)

where we have used $p' = p + q$. Hence, we have

$$\sum \sum |T|^2 = \frac{e^2g_{T}^2}{8(q^2 - M_{\rho}^2)^2} \sum_{\gamma} \sum_{i,j} \sum_{m,l} \sum_{\mu,\mu'}$$

$$\left\{ \frac{1}{2}\epsilon_{i}(\gamma)(-g_{j\mu} + \frac{q_{j}q_{\mu}}{M_{\rho}^2}) + \frac{1}{2}\epsilon_{j}(\gamma)(-g_{i\mu} + \frac{q_{i}q_{\mu}}{M_{\rho}^2})$$

$$-\frac{1}{3}\epsilon_{m}(\gamma)\delta_{ij} (-g_{m\mu} + \frac{q_{m}q_{\mu}}{M_{\rho}^2}) \right\}$$

$$\sum \sum |T|^2,$$

and we sum explicitly over all the indices and the two photon polarization which we write explicitly as

$$\epsilon^{(1)}(\gamma) = \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix}, \quad \epsilon^{(2)}(\gamma) = \begin{pmatrix} 0 \\ 1 \\ 0 \end{pmatrix},$$

(16)

where we have assumed that the photon travels in the $Z$ direction.

If we consider the $f_{2}(1270)$ as an elementary particle, we find

$$\frac{d\sigma}{d\Omega} = \frac{m_{\rho}^2}{16\pi s |\vec{k}|} \sum \sum |T|^2,$$

(17)

where $\vec{p}_{f_{2}}$ and $\vec{k}$ are the three momenta of final $f_{2}(1270)$ and initial photon in the center of mass frame (c.m.), and taking into account that $t = q^2 = (p - p')^2 = 2m_{\rho}^2 - 2E(p)E(p') + 2p \cdot p'$, we get

$$\frac{d\sigma}{dt} = \frac{m_{\rho}^2}{16\pi s|\vec{k}|^2} \sum \sum |T|^2.$$  

(18)

FIG. 3: Feynman diagram for the $\gamma p \rightarrow p\pi^+\pi^-$ reaction.

Eq. (18) can be generalized for the case when the $f_{2}(1270)$ is explicitly allowed to decay into two pions as shown in Fig. 3 by working out the three body phase space and we find

$$\frac{d\sigma}{dM_{\rho\rho} dt} = \frac{m_{\rho}^2}{8\pi s|\vec{k}|^2} \frac{M_{\rho\rho}^2 M_{f_{2}}^2}{|M_{inv} - M_{f_{2}} + i M_{inv} f_{2} |^2} \times \sum \sum |T|^2,$$

(19)

where $M_{inv}$ is the invariant mass distribution of the two pions, $\Gamma_{f_{2}}$ is the total decay width of the $f_{2}(1270)$ and

$^1$ In Ref. [11] the $\sqrt{2}$ factor is not implemented because it is compensated by not dividing by two the integrated width, as it corresponds to two final identical particles (two photons).
\( \Gamma_{\pi} \) is the partial decay width of the \( f_2(1270) \) into the \( \pi\pi \) system, in our case \( \pi^+\pi^- \). The \( \pi^+\pi^- \) decay accounts for \( \frac{\sqrt{2}}{4} \) of the \( \pi\pi \) decay width of the \( f_2(1270) \), which is 85\% of \( \Gamma_{f_2} \). Since the \( \pi \) decay is in \( D \)-wave, in order to have \( \Gamma_{\pi} \) and \( \Gamma_{f_2} \) in the range of invariant masses that we consider (close to the \( f_2(1270) \) resonance), we take

\[
\Gamma_{\pi}(M_{\text{inv}}) = \Gamma_{\pi}^0 \left( \frac{\hat{q}}{\hat{q}_0} \right)^5 \left( \frac{M_{f_2}}{M_{\text{inv}}} \right)^2,
\]

\[
\Gamma_{f_2}(M_{\text{inv}}) = 0.85 \Gamma_{\pi}^0 \left( \frac{\hat{q}}{\hat{q}_0} \right)^5 \left( \frac{M_{f_2}}{M_{\text{inv}}} \right)^2 + 0.15 \Gamma_{\pi}^0,
\]

with \( \Gamma_{\pi}^0 = 185 \text{ MeV} \), \( \Gamma_{\pi}^0 = 105 \text{ MeV} \), and \( M_{f_2} = 1275 \text{ MeV} \). And

\[
\hat{q} = \frac{\lambda^{1/2}(M_{\text{inv}}^2, m_\pi^2, m_\pi^2)}{2M_{\text{inv}}^2},
\]

\[
\hat{q}_0 = \frac{\lambda^{1/2}(M_{f_2}^2, m_\pi^2, m_\pi^2)}{2M_{f_2}^2},
\]

where \( \lambda \) is the Källén function with \( \lambda(x, y, z) = (x - y - z)^2 - 4yz \).

### III. NUMERICAL RESULTS

In Ref. [18] the \( \gamma p \rightarrow \pi^+\pi^-p \) reaction was studied in the photon energy range \( 3.0 - 3.8 \text{ GeV} \). The magnitude \( \frac{d\sigma}{dM_{\text{inv}}dt} \) was measured, where \( M_{\text{inv}} \) is the invariant mass of the two pions, \( t \) the momentum transfer squared and \( \Omega_{\pi} \), the angle of one pion measured in the \( \pi^+\pi^- \) helicity rest frame. This multiple differential cross section was then projected over the appropriate partial waves with the integral over the corresponding spherical-harmonic functions. After correcting for detector acceptance and detector efficiency, \( \frac{d\sigma}{dM_{\text{inv}}dt} \) for the corresponding partial waves was produced, such that a direct comparison with a theoretical model can be done. The projection over \( \pi^+\pi^- \) \( D \)-wave was done and a peak around 1270 MeV, corresponding to the \( f_2(1270) \) resonance, was found (see Fig. 1) for a certain cut of the photon energy and the variable \( t \).

Simultaneously, a photon energy averaged cross section \( \frac{d\sigma}{dt} \) was determined by considering a wide range of energies and integrating \( \frac{d\sigma}{dM_{\text{inv}}dt} \) over \( M_{\text{inv}} \) around the resonance peak. These are magnitudes that we can easily address with our theoretical framework and we show below the results obtained.

In Fig. 2 we shown the results for \( \frac{d\sigma}{dM_{\text{inv}}dt} \) as a function of \( M_{\text{inv}} \) for \( E_y = 3.3 \text{ GeV} \) and \( t = -0.55 \text{ GeV}^2 \), which we compare with the results of Ref. [18] obtained in the range \( 3.2 < E_y < 3.4 \text{ GeV} \), \( -0.6 < t < -0.5 \text{ GeV}^2 \).

As we can see, the experimental data have a wide band of allowed values, but our results around the peak go through the middle of the band. Discrepancies below the resonance peak can be attributed to background which we have not considered in our approach, since only the resonance contribution is taken into account.

Further information can be obtained from the \( t \) dependence of the cross section. In Ref. [18] one finds \( \frac{d\sigma}{dt} \) as a function of \( t \), obtained by integrating \( M_{\text{inv}} \) in the range \( 1.275 \pm 0.185 \text{ GeV} \) and for photon energy in the range \( E_y = 3.0 - 3.8 \text{ GeV} \). We perform the calculation for this observable and the results are shown in Fig. 3. Once again, we see that we get a good agreement with experiment.

Although the data seem to imply a slightly bigger slope than provided by the theory, the fact is that the results, with no free parameters, agree well with the data within errors. It seems clear that the slope provided by the rho meson propagator in our approach accounts for the bulk of the \( t \) dependence of the cross section, but one cannot exclude extra elements to the theory that gradually change the cross section at larger values of \( t \), too large to be accommodated within an effective theory as we have used. Yet, for the range of energies and momentum transfers measured, it looks clear that our approach provides the basic features of the experiment and globally we can claim a good agreement with experiment.

### IV. CONCLUSIONS

In this paper we have studied the \( \gamma p \rightarrow \pi^+\pi^-p \) reaction performed in Jefferson Lab, where the two pions have been separated in \( D \)-wave, producing the \( f_2(1270) \) resonance. This resonance has been much studied recently from the point of view of a \( \rho \rho \) molecule and has passed all tests in the reactions where it has been studied. Yet, all the reactions were mesonic reactions. This is the first time where this idea has been tested in a baryonic reaction. The elements needed for the test are very simple, which offers a special transparency in the interpretation of the results.

![Graph showing theoretical prediction for D-wave ππ mass distribution at Eγ = 3.3 GeV and t = −0.55 GeV² compared with the CLAS data taken from Ref. [18]](image-url)
FIG. 5: Differential cross section $\frac{d\sigma}{dt}$ as a function of $t$. The experimental data are taken from Ref. [18].

in $I = 0$ and the value of the coupling has been obtained before in the theory that provides the $f_2(1270)$ as a $\rho \rho$ molecule based on the local hidden gauge formalism for the interaction of vector mesons. With this coupling and the vector meson dominance hypothesis, incorporated in the local hidden gauge approach, the photon gets converted into one of the $\rho^0$ of the $\rho \rho$ formation state of the $f_2(1270)$, and the other $\rho^0$ acts as a mediator between the photon and the proton.

With this simple picture we determine both the differential cross section and the $t$ dependence of the integrated cross section over the invariant mass around the resonance without any free parameter. The agreement with the experimental differential cross sections and the $t$ dependence is good, thus, providing new support for the $\rho \rho$ molecular picture of the $f_2(1270)$.

Acknowledgments

One of us, E. O., wishes to acknowledge support from the Chinese Academy of Science (CAS) in the Program of Visiting Professorship for Senior International Scientists. This work is partly supported by the Spanish Ministerio de Economia y Competitividad and European FEDER funds under the contract number FIS2011-28853-C02-01 and FIS2011-28853-C02-02, and the Generalitat Valenciana in the program Prometeo II-2014/068. We acknowledge the support of the European Community-Research Infrastructure Integrating Activity Study of Strongly Interacting Matter (acronym HadronPhysics3, Grant Agreement n. 283286) under the Seventh Framework Programme of EU. This work is also partly supported by the National Natural Science Foundation of China under Grant Nos. 11105126 and 11475227. The Project is Sponsored by the Scientific Research Foundation for the Returned Overseas Chinese Scholars, State Education Ministry.

[1] H. Marsiske et al. [Crystal Ball Collaboration], Phys. Rev. D 41, 3324 (1990).
[2] T. Oest et al. [JADE Collaboration], Z. Phys. C 47, 343 (1990).
[3] E. Klempt and A. Zaitsev, Phys. Rept. 454, 1 (2007).
[4] V. Crede and C. A. Meyer, Prog. Part. Nucl. Phys. 63, 74 (2009).
[5] R. Molina, D. Nicmorus and E. Oset, Phys. Rev. D 78, 114018 (2008).
[6] M. Bando, T. Kugo, S. Uehara, K. Yamawaki and T. Yanagida, Phys. Rev. Lett. 54, 1215 (1985).
[7] M. Bando, T. Kugo and K. Yamawaki, Phys. Rept. 164, 217 (1988).
[8] M. Harada and K. Yamawaki, Phys. Rept. 381, 1 (2003).
[9] U. G. Meiissner, Phys. Rept. 161, 213 (1988).
[10] L. S. Geng and E. Oset, Phys. Rev. D 79, 074009 (2009).
[11] H. Nagahiro, J. Yamagata-Sekihara, E. Oset, S. Hirenzaki and R. Molina, Phys. Rev. D 79, 114023 (2009).
[12] T. Branz, L. S. Geng and E. Oset, Phys. Rev. D 81, 054037 (2010).
[13] A. Martinez Torres, L. S. Geng, L. R. Dai, B. X. Sun, E. Oset and B. S. Zou, Phys. Lett. B 680, 310 (2009).
[14] L. S. Geng, F. K. Guo, C. Hanhart, R. Molina, E. Oset and B. S. Zou, Eur. Phys. J. A 44, 305 (2010).
[15] L. Dai and E. Oset, Eur. Phys. J. A 49, 130 (2013).
[16] J. J. Xie and E. Oset, Phys. Rev. D 90, 094006 (2014).
[17] R. Aaij et al. [LHCb Collaboration], Phys. Rev. D 89, no. 9, 092006 (2014).
[18] M. Battaglieri et al. [CLAS Collaboration], Phys. Rev. D 80, 072005 (2009).
[19] L. Bibbyczynski and R. Kaminski, Phys. Rev. D 87, 114010 (2013).
[20] A. Ramos and E. Oset, Phys. Lett. B 727, 287 (2013).
[21] H. Nagahiro, L. Roca, A. Hosaka and E. Oset, Phys. Rev. D 79, 014015 (2009).
[22] K. A. Olive et al. [Particle Data Group Collaboration], Chin. Phys. C 38, 090001 (2014).