The study of three-body systems is one of the important issues of contemporary nuclear physics and has attracted continuous attention. Chronicled examples are baryonic three-body bound systems, such as tritium, $^3\text{He}$ ($NNN$) and hypertriton ($pn\Lambda$). Recently, interest in three-body systems has developed, and resonance systems including mesons as the constituents are considered based on current understanding of hadronic interaction.

Evidences for several existing and new states which can be interpreted as three-body resonances are being reported from theoretical and experimental studies. For example, it has been claimed that the $Y(4660)$ resonance found in $e^+e^-\to\gamma_{ISR}\pi^+\pi^-\psi$ can be interpreted as a $f_0(980)\psi$ bound state \cite{1}. For the $\pi KN$ system, the Faddeev equations were solved using unitary chiral dynamics and coupled channels and dynamical generation of all the $\Sigma$ and $\Lambda$ resonance states with $J^P = 1/2^+$ listed by the Particle Data Group (PDG) in the energy region 1500-1800 MeV was found in Ref. \cite{2}. The same formalism applied to the $\pi\pi N$ system and its coupled channels revealed the dynamical generation of the $N^*(1710)$, $N^*(2100)$ and $\Delta(1910)$ \cite{3}. The $X(2175)$ state, reported by different experimental groups \cite{4} in the $\phi f_0(980)$ invariant mass, has been explained as a $\phi K\bar{K}$ resonance with $K\bar{K}$ forming the $f_0(980)$ resonance \cite{4}.

Further, kaonic nuclear few-body systems are of special interest in relation with strangeness nuclear physics. Possible existence of $KKNN$ bound states was pointed out in 60’s \cite{9} by considering the $\Lambda(1405)$ hyperon resonance as a quasibound state of $KN$ as suggested in Ref. \cite{10} \cite{11}. Recently thorough theoretical investigations of the $KKNN$ system in various approaches \cite{12} \cite{13} \cite{14} \cite{15} \cite{16} indicate a quasibound state with a large width. Baryonic systems with two kaons were also investigated in Refs. \cite{17} \cite{18} with a single channel variational method. For the $K\bar{K}N$ system \cite{18}, a quasibound state of these hadrons was found around 1910 MeV for an $N^*$ with $I = 1/2$ and $J^P = 1/2^+$ using the effective $KN$ potential derived in Ref. \cite{19} and the $K\bar{K}$ interactions reproducing $f_0(980)$ and $a_0(980)$ as $K\bar{K}$ quasibound states with 980 MeV mass and 60 MeV width. In case of the $N^*$ state found in Ref. \cite{18} the $KN$ pair forms the $\Lambda(1405)$ and, simultaneously, the $K\bar{K}$ pair is resonating as the $a_0(980)$. The same state was also found independently in a study of the $NKK$, $N\pi\pi$ and $N\pi\eta$ coupled channels based on solution of the Faddeev equations \cite{20}. There it was concluded that a state indeed appears around 1920 MeV when the hadrons rearrange themselves to form a $a_0(980)N$ system and that the contribution of the $N\pi\pi$ and $N\pi\eta$ channels was negligible in the dynamical generation of this resonance. A $N^*$ state with these properties is not listed by the PDG, however, there have been some proposals that this state can be seen in the data for the $\gamma p \to K^+\Lambda$ reaction \cite{21} \cite{22}, although the situation is still controversial \cite{23} \cite{24}.

The single channel variational approach \cite{18} found that this new $N^*$ state substantially contains $\Lambda(1405)$ in the $KN$ subsystem. In the Faddeev analysis \cite{20}, although two-body coupled channels were fully considered, explicit three-body channels of $K\pi\Sigma$ and $K\pi\Lambda$ were not taken into account. The $K\pi\Sigma$ channel could produce some changes on the characteristics of the $N^*$ state found in Ref. \cite{18} \cite{20}, since the $\Lambda(1405)$ is dynamically generated in coupled channels such as $KN$ and $\pi\Sigma$ \cite{24} \cite{30}, especially the lower pole \cite{2} of the $\Lambda(1405)$ couples strongly to the $\pi\Sigma$ channel \cite{29} \cite{32}. Thus, this article is devoted to further clarification of the properties of this new $N^*$ state. To do that we follow the coupled-channel Faddeev approach developed in Ref. \cite{2} but taking into account coupled channels of $K\pi\Sigma$ and $K\pi\Lambda$ together with $KN$.

Let us briefly explain our formulation to study three-body coupled channels of two mesons and one baryon with $J^P = 1/2^+$. Paying special attention to the dynamical generation of the $\Lambda(1405)$ in the $K\bar{K}N$ and $\pi\Sigma$ subsystems, we consider the $KKNN$, $K\bar{K}\Sigma$ and $K\bar{K}\pi\Sigma$ channels with total charge zero, namely, in the charge base, $K^0\bar{K}^0$, $K^0\bar{K}^0\Sigma$, $K^0\bar{K}^0\Sigma$, $K^0\bar{K}^0\pi^\pm\Sigma$, $K^0\bar{K}^0\pi^\mp\Sigma$, $K^0\bar{K}^0\pi^\pm\Sigma$, $K^0\bar{K}^0\pi^\mp\Sigma$, $K^0\bar{K}^0\pi^\pm\Sigma$, $K^0\bar{K}^0\pi^\mp\Sigma$, $K^+\bar{K}^-\pi^+\Sigma$, $K^+\bar{K}^-\pi^+\Sigma$, $K^+\bar{K}^-\pi^+\Sigma$, $K^+\bar{K}^-\pi^+\Sigma$, $K^+\bar{K}^-\pi^+\Sigma$, and calculate the three-body $T$ matrix for the different transitions. To determine the three-body $T$ matrix we follow the formalism developed in Refs. \cite{2} \cite{3} \cite{20} \cite{34} which is based on the Faddeev equations \cite{35}. In terms of the Faddeev partitions, $T^1, T^2$ and $T^3$, the three-body $T$-matrix is written as

$$T = T^1 + T^2 + T^3. \quad (1)$$

\footnote{A recent investigation using a coupled channel approach based on chiral dynamics also confirmed that the $\Lambda(1405)$ resonance can be described substantially by a meson-baryon molecular state \cite{11}.}

\footnote{It was pointed out in Refs \cite{24} \cite{25} \cite{31} \cite{32} that there exist two poles around the $\Lambda(1405)$ energy region, having them different coupling nature to meson-baryon states. Detailed discussion on the double pole structure of the $\Lambda(1405)$ in chiral unitary approaches can be found in Refs. \cite{29} \cite{35}.}
In our formalism these partitions are expressed as

\[ T^i = t^i \delta^3(\vec{k}'_i - \vec{k}_i) + \sum_{j\neq i=1}^{3} T^{ij}_R, \]  

for \( i = 1, 2, 3 \) with \( \vec{k}_i (\vec{k}'_i) \) being the initial (final) momentum of the particle \( i \) and \( t^i, i = 1, 2, 3 \), the two-body \(-\)matrix which describes the interaction for the \((jk)\) pair of the system, \( j \neq k \neq i = 1, 2, 3 \). In our approach, this two-body \(-\)matrix is calculated by solving the Bethe-Salpeter equation with the potential obtained from chiral Lagrangians \([27, 29, 36, 37]\). Namely we consider all possible two-body channels of meson \((\pi, \eta, K, \bar{K})\) and baryon \((N, \Lambda, \Sigma, \Xi)\) which couple to \( KK \), \( \pi K, \bar{K} N, K \Sigma \) and \( \bar{K} N \), but except for the \( nn \) channel, which is not important in \( \pi \) and \( K \) dynamics \([36]\).

In Eq. (2), the \( T^{ij}_R \) partitions include all the different contributions to the three-body \(-\)matrix in which the last two interactions are given in terms of the two-body \(-\)matrices \( t^i \) and \( t^j \), respectively, and satisfy the following set of coupled equations

\[ T^{ij}_R = t^i g^{ij} t^j + t^i \left[ G^{ijk} T^{ij}_R + G^{ijk} T^{jk}_R \right], \]  

for \( i = 1, 2, 3 \). In Eq. (3), \( g^{ij} \)'s correspond to the three-body Green's function of the system and its elements are defined as

\[ g^{ij}(\vec{k}_i, \vec{k}_j) = \left( \frac{-N_k}{2E_k(\vec{k}_i + \vec{k}_j)} \right) \times \frac{1}{\sqrt{s} - E_i(\vec{k}_i) - E_j(\vec{k}_j) - E_k(\vec{k}_i + \vec{k}_j) + i\epsilon}, \]  

with \( N_k = 1 \) for mesons and \( N_k = 2M_k \) for baryons with baryon mass \( M_k \), and \( E_l, l = 1, 2, 3 \), is the energy of the particle \( l \). The \( G^{ijk} \) matrix in Eq. (3) represents a loop function of three-particles and it is written as

\[ G^{ijk} = \int \frac{d^3k''}{(2\pi)^3} \bar{g}^{ij}(\vec{k}'', \vec{s}_m) F^{ijk}, \]  

with the elements of \( \bar{g}^{ij} \) being

\[ \bar{g}^{ij}(\vec{k}'', \vec{s}_m) = \frac{N_k}{2E_k(\vec{k}'')} \frac{N_m}{2E_m(\vec{k}'')} \times \frac{1}{\sqrt{s}_{lm} - E_i(\vec{k}'') - E_m(\vec{k}'') + i\epsilon}, \]  

for \( i \neq l \neq m \), and the matrix \( F^{ijk} \), with explicit variable dependence, is given by

\[ F^{ijk}(\vec{k}'', \vec{k}_i, \vec{k}_j, \vec{s}_m) = t^{ij}(\vec{s}_m) \bar{g}^{ik}(\vec{k}'', \vec{k}_i) g^{jk}(\vec{k}_j, \vec{k}_k) \left[ t^{ij}(\vec{s}_m) \right]^{-1}, \]  

for \( j \neq r \neq u = 1, 2, 3 \). In Eq. (6), \( \sqrt{s}_{lm} \) is the invariant mass of the \((lm)\) pair and can be calculated in terms of the external variables. The upper index \( k'' \) in the invariant mass \( s_{k''} \) of Eq. (7) indicates its dependence on the loop variable, as it was shown in Ref. \([2]\). The main advantage of Eq. (3) is that they are algebraic coupled equations and not integral equations as it was shown in Refs. \([2, 20]\). In these works, for the first time, cancelling between the contribution of the off-shell parts of the chiral two-body \(-\)matrices to the three-body diagrams and the contact term with three particles in the initial and final state, whose origin is in the chiral Lagrangian used to describe the interaction, was found analytically (see Refs. \([2, 3, 20]\) for more details).

In terms of the \( T^{ij}_R \) partitions (Eq. 3), the expression for the full three-body \(-\)matrix can be obtained combining Eq. (1) and Eq. (2). Nontrivial three-body dynamics appears in

\[ T_R = \sum_{i=1}^{3} \sum_{j \neq i=1}^{3} T^{ij}_R. \]  

This amplitude is a function of the total three-body energy, \( \sqrt{s} \), and the invariant mass of the particles 2 and 3, \( \sqrt{s_{23}} \). The other invariant masses, \( \sqrt{s_{21}} \) and \( \sqrt{s_{32}} \), can be obtained in terms of \( \sqrt{s} \) and \( \sqrt{s_{23}} \), as it was shown in Ref. \([2, 3]\).

To present our results, we have chosen \( \sqrt{s} \) and the invariant mass of one of the two-body subsystems, \( \sqrt{s_{23}} \). All the matrices in Eq. (3) are projected in S-wave, thus the quantum numbers of the three-body system and, hence, the resulting resonances are \( J^P = 1/2^+ \).

Let us discuss the results obtained for the three-body amplitude \( T_R \). Our discussion concentrates on showing \( |T_R|^2 \) for real values of \( \sqrt{s} \) and of the invariant mass of one of the two-body subsystems, in concrete, \( \sqrt{s_{KN}} \) and \( \sqrt{s_{\bar{K}N}} \). We have solved Eq. (3) in the charge base, thus, to study the existence of a three-body \( N^* \) resonance around 1920 MeV we have to project \( T_R \) on the isospin base with total isospin \( I = 1/2 \). Our interest is to examine the possibility of existence of a \( N^* \) resonance which appears as a \( K \bar{K} N \) bound state when the \( K \bar{K} \) subsystem and its coupled channels generate the \( \Lambda(1405) \), as pointed out in Ref. \([18]\). For this purpose, we specify also the isospin of the two-body subsystem. First, we calculate the three-body \( T_R \) matrix for the \( K \bar{K} N \) channel projected on total isospin \( I = 1/2 \) and the \( K \bar{K} N \) subsystem with \( I_{KN} = 0 \), which we represent as \( \langle T^{KKKN}(I=1/2,J_{KN}=0) \rangle = \langle I = 1/2, I_{KN} = 0 | T^{KKKN}_R | I = 1/2, I_{KN} = 0 \rangle \). In addition, to understand the characteristics of this \( N^* \) further, we also show the \( T_R \) matrix projected on \( I = 1/2 \) with the \( K \bar{K} \) subsystem in isospin one, \( I_{K \bar{K}} = 1 \), which is denoted by \( \langle T^{KKKN}(I=1/2,J_{K \bar{K}}=1) \rangle = \langle I = 1/2, I_{K \bar{K}} = 1 | T^{KKKN}_R | I = 1/2, I_{K \bar{K}} = 1 \rangle \). In the \( K \bar{K} \) two-body subsystem with \( I_{K \bar{K}} = 1 \), \( a_0(980) \) is dynamically generated (Ref. \([36, 37]\)).

In Fig. 1 we show the contour plots corresponding to the three-dimensional plots of the squared three-body \( T_R \) matrix, \( \langle T^{KKKN}_R(I=1/2,J_{KN}=0) \rangle \) (upper panel) and \( \langle T^{KKKN}_R(I=1/2,J_{K \bar{K}}=1) \rangle \) (lower panel) plotted as functions of the total energy of the three-body system, \( \sqrt{s} \), and the \( K \bar{K} \) invariant mass, \( \sqrt{s_{\bar{K}N}} \), and the \( K \bar{K} \) invariant mass, \( \sqrt{s_{KK}} \), respectively. As it can be seen in the upper panel, a peak in the
The present case, the $KN$ subsystem has purely tiny contributions of the $KK$ subsystem. We also find the $NN$ component is favored in the $KK$ subsystem. Group theory tells us that, in case the $KN$ subsystem is strong enough to compensate the repulsion in the $KN$ subsystem, the $KN$ subsystem has an invariant mass close to 1428 MeV. In the lower panel, the peak shows up when the invariant mass of the $KK$ subsystem is around 987 MeV.

We also find the $N^*$ resonance at the same value of $\sqrt{s}$ in the $T_R$ matrices for different isospin combinations of the $KN$ and $KK$ subsystems. For the case in which the $KN$ subsystem is in isospin 1, $I_{KN} = 1$, the $(T_R^{KKN})^{(I=1/2,I_{KN}=1)}$ matrix shows a less pronounced peak structure for the $N^*$, due to the fact that the projected $T_R$ matrix on $I_{KN} = 1$ has tiny contributions of the $\Lambda(1405)$ in the intermediate states. The ratio of the $|T_R|^2$ matrices with $I_{KK} = 0$ and $I_{KK} = 1$ at the resonance point, $(\sqrt{s},\sqrt{s_{KK}}) = (1923, 1428)$ MeV, is found to be a tiny value, $\sim 0.008$. Also, for the $KK$ subsystem, the ratio of $|T_R|^2$ with $I_{KK} = 0$ and $I_{KK} = 1$ at $(\sqrt{s},\sqrt{s_{KK}}) = (1923, 987)$ MeV is $\sim 1$. Although the magnitude for these two matrix elements with $I_{KK} = 0$ and $I_{KK} = 1$ is very similar, it does not mean that the fraction of the $I_{KK} = 0$ and $I_{KK} = 1$ components in the $N^*$ state is similar, since this fraction depend on the isospin configuration of the $KN$ subsystem. Group theory tells us that, in case the $KN$ subsystem has purely $I_{KN} = 0$, the ratio of the $I_{KK} = 1$ and $I_{KK} = 0$ components is 3 to 1 for total $I = 1/2$. Since, in the present case, the $KN$ pair dominantly has $I_{KN} = 0$, the $I_{KK} = 1$ component is favored in the $N^*$ state. This implies that the $N^*$ resonance contains mostly $a_0(980)$ in the $KK$ subsystem with $I_{KK} = 1$ and less contribution from $f_0(980)$ with $I_{KK} = 0$. The attraction in the $KK$ and $KN$ subsystems is strong enough to compensate the repulsion in the $KN$ subsystem and form a bound state.

It is well known that the $KN$ interaction and coupled channels generate the $\Lambda(1405)$ state, with a double pole structure associated with the energy in which the $KN$ subsystem has an invariant mass of around 987 MeV. This fact could be related to the double pole structure of the $\Lambda(1405)$.

![FIG. 1. Contour plots of the three-body squared amplitude $|T_R|^2$ for the $N^*$ resonance in the $KKN$ system as functions of the total three-body energy, $\sqrt{s}$, and the invariant mass of the $KN$ subsystem with $I_{KN} = 0$ (upper panel) or the invariant mass of the $KK$ subsystem with $I_{KK} = 1$ (lower panel).](image1)

![FIG. 2. Dashed Lines: (Left) $K\pi \rightarrow K\pi t$ matrix with isospin $1/2$. The $K$ resonance is generated around 850 MeV. (Right) $t$ matrix for the $K\Sigma \rightarrow K\Sigma$ transition with isospin $1/2$, which shows the presence of the $N^*(1535)$. Solid line: Energy range used in the three-body calculation for the $K\pi$ interaction (Left) and for the $K\Sigma$ interaction (Right). The units are arbitrary.](image2)
Recently, it has been pointed out in Ref. [19] that two resonance poles are found in the \( KNN \) system with total isospin \( I = 1/2 \) and \( J^P = 1/2^+ \) between the \( KNN \) and \( \pi\Sigma N \) threshold, when a Weinberg-Tomozawa energy dependent two-body interaction is used for a three-body Faddeev calculation: one pole is located moderately below the \( \Lambda(1405) \) threshold with a narrow width, while the other one appears above the \( \pi\Sigma N \) threshold with a substantially large width. These poles are associated with the two-pole nature of the \( \Lambda(1405) \) generated in the \( K\bar{N} \) and \( \pi\Sigma \) subsystems. We have also looked for another resonance state associated with the lower \( \Lambda(1405) \) pole, but we could not find any signal for such a deep state in the \( T_R \) matrix evaluated with real values of \( \sqrt{\lambda} \). This means either that there is no such a resonance state or that there is a resonance state having such a large width that the resonance contribution cannot be seen in the real axis of \( \sqrt{\lambda} \).

To summarize, we have studied the \( K\bar{K}N \), \( K\pi\Sigma \) and \( K\pi\Lambda \) systems by solving the Faddeev equations in a coupled channel approach. The input two-body \( t \)-matrices have been obtained by using potentials from chiral Lagrangians and solving the Bethe-Salpeter equations in a unitary coupled channel approach. We have found the contribution of the \( K\pi\Sigma \) and \( K\pi\Lambda \) channels to the three-body \( T \) matrix to be negligible around a total energy for the three-body system close to 1920 MeV and, thus, one could solve the Faddeev equations considering only the \( K\bar{K}N \) channel. The resolution of these equations has lead to the dynamical generation of a \( N^* \) resonance around 1920 MeV with \( J^P = 1/2^+ \), as was predicted in [18] and found in [20]. The resonance is generated when the \( K\bar{N} \) subsystem is resonating as the \( \Lambda(1405) \) and, at the same time, the \( K\bar{K} \) interaction generates the \( a_0 (980) \) resonance.

The authors thank Dr. Hyodo for his useful comments, a part of which motivated this work, and Professors E. Oset, K. Kanada-Eny’o and Dr. K. P. Khemchandani for collaboration in the previous works on which the present one is based. The work of A. M. T. is supported by the Grant-in-Aid for the Global COE Program “The Next Generation of Physics, Spun from Universality and Emergence” from the Ministry of Education, Culture, Sports, Science and Technology (MEXT) of Japan. This work is supported in part by the Grant for Scientific Research (No. 22105507) from MEXT of Japan. A part of this work was done in the Yukawa International Project for Quark-Hadron Sciences (YIPQS).

[1] F. K. Guo, C. Hanhart and U. G. Meissner, Phys. Lett. B 665, 26 (2008).
[2] A. Martinez Torres, K. P. Khemchandani and E. Oset, Phys. Rev. C 77, 042203 (2008).
[3] K. P. Khemchandani, A. Martinez Torres and E. Oset, Eur. Phys. J. A 37, 233 (2008).
[4] B. Aubert et al. [BABAR Collaboration], Phys. Rev. D 74, 091103 (2006); Phys. Rev. D 76, 012008 (2007).
[5] M. Ablikim et al. [BES Collaboration], Phys. Rev. Lett. 100, 102003 (2008).
[6] A. Martinez Torres, K. P. Khemchandani, L. S. Geng, M. Nap-sciale and E. Oset, Phys. Rev. D 78, 074031 (2008).
[7] L. Alvarez-Ruso, J. A. Oller, J. M. Alarcon, Phys. Rev. D 80, 054011 (2009).
[8] Susana Coito, George Rupp, Eef van Beveren, Phys. Rev. D 80, 094011 (2009).
[9] Y. Nagami, Phys. Lett. B, 288 (1963).
[10] R. H. Dalitz and S. F. Tuan, Phys. Rev. Lett. 2, 425 (1959); Annals Phys. 10, 307 (1960).
[11] T. Hyodo, D. Jido and A. Hosaka, Phys. Rev. C 78, 025203 (2008).
[12] Y. Akaishi and T. Yamazaki, Phys. Rev. C 65, 044005 (2002); T. Yamazaki and Y. Akaishi, Phys. Rev. C 76, 045201 (2007).
[13] N. V. Shevchenko, A. Gal and J. Mares, Phys. Rev. Lett. 98, 082301 (2007); N. V. Shevchenko, A. Gal, J. Mares and J. Revai, Phys. Rev. C 76, 044004 (2007).
[14] Y. Ikeda and T. Sato, Phys. Rev. C 76, 035203 (2007); Y. Ikeda and T. Sato, Phys. Rev. C 79, 035201 (2009).
[15] A. Dote, T. Hyodo and W. Weise, Nucl. Phys. A 804, 197 (2008).
[16] S. Wycech and A. M. Green, Phys. Rev. C 79, 014001 (2009).
[17] Y. Kanada-En’yo and D. Jido, Phys. Rev. C 78, 025212 (2008).
[18] D. Jido and Y. Kanada-En’yo, Phys. Rev. C 78, 035203 (2008).
[19] T. Hyodo and W. Weise, Phys. Rev. C 77, 035204, (2008).
[20] A. Martinez Torres, K. P. Khemchandani and E. Oset, Phys. Rev. C 79, 065207 (2009).
[21] T. Mart and C. Bennhold, Phys. Rev. C 61, 012201 (2000).
[22] A. Martinez Torres, K. P. Khemchandani, U. G. Meissner and E. Oset, Eur. Phys. J. A 41, 361 (2009).
[23] K. H. Glander et al., Eur. Phys. J. A 19, 251 (2004).
[24] R. Bradford et al. [CLAS Collaboration], Phys. Rev. C 73, 035202 (2006).
[25] M. Sumihama et al. [LEPS Collaboration], Phys. Rev. C 73, 035214 (2006).
[26] N. Kaiser, P. B. Siegel, and W. Weise, Nucl. Phys. A 594, 325 (1995).
[27] E. Oset and A. Ramos, Nucl. Phys. A 635, 99 (1998).
[28] J. A. Oller and U. G. Meissner, Phys. Lett. B 500, 263 (2001).
[29] D. Jido, J. A. Oller, E. Oset, A. Ramos and U. G. Meissner, Nucl. Phys. A 725, 181 (2003).
[30] T. Hyodo, S. I. Nam, D. Jido, and A. Hosaka, Phys. Rev. C 68, 018201 (2003); Prog. Theor. Phys. 112, 73 (2004).
[31] In a different context, the double pole structure of the \( \Lambda(1405) \) was discussed in P. J. Fink, Jr., G. He, R. H. Landau and J. W. Schnick, Phys. Rev. C 41, 2720 (1990).
[32] D. Jido, A. Hosaka, J. C. Nacher, E. Oset and A. Ramos, Phys. Rev. C 66, 025203 (2002).
[33] D. Jido, T. Sekihara, Y. Ikeda, T. Hyodo, Y. Kanada-En’yo and E. Oset, Nucl. Phys. A 835, 59 (2010).
[34] K. P. Khemchandani, A. Martinez Torres and E. Oset, Phys. Lett. B 675, 407 (2009).
[35] L. D. Faddeev, Sov. Phys. JETP 12, 1014 (1961) [Zh. Eksp. Teor. Fiz. 39, 1459 (1960)].
[36] J. A. Oller and E. Oset, Nucl. Phys. A 620, 438 (1997) [Erratum-ibid. A 652, 407 (1999)].
[37] J. A. Oller, E. Oset and J. R. Pelaez, Phys. Rev. D 59, 074001 (1999) [Erratum-ibid. D 60, 099906 (1999) Erratum 75, 099903 (2007)].
[38] T. Inoue, E. Oset and M. J. Vicente Vacas, Phys. Rev. C 65, 035204 (2002).
[39] Y. Ikeda, H. Kamano and T. Sato, arXiv:1004.4877 [nucl-th].