$\bar{K}NN$ Absorption within the Framework of the Fixed Center Approximation to Faddeev equations

M. Bayar[1],[2] and E. Oset[1]

[1]*Instituto de Física Corpuscular (centro mixto CSIC-UV)*
Institutos de Investigación de Paterna,
Aptdo. 22085, 46071, Valencia, Spain,
[2]*Department of Physics, Kocaeli University, 41380 Izmit, Turkey*
(Dated: July 9, 2012)

Abstract

We present a method to evaluate the $\bar{K}$ absorption width in the bound $\bar{K}NN$ system. Most calculations of this system ignore this channel and only consider the $\bar{K}N \to \pi\Sigma$ conversion. Other works make a qualitative calculation using perturbative methods. Since the $\Lambda(1405)$ resonance is playing a role in the process, the same resonance is changed by the presence of the absorption channels and we find that a full nonperturbative calculation is demanded, which we present here. We employ the Fixed Center Approximation to Faddeev equations to account for $\bar{K}$ rescattering on the (NN) cluster and we find that the width of the states found previously for $S = 0$ and $S = 1$ increases by about 30 MeV due to the $\bar{K}NN$ absorption, to a total width of about 80 MeV.

arXiv:1207.1661v1 [hep-ph] 6 Jul 2012

I. INTRODUCTION

The interaction antikaon ($\bar{K}$) with nucleon (N) has drawn attention for decades [1–11] since this interaction provides essential elements to understand the strangeness in nuclear hadron physics. The $\bar{K}N$ interaction is strongly attractive in the isospin $I = 0$ and also the $\pi\Sigma$ interaction is attractive at low energies. The $\bar{K}N$ and $\pi\Sigma$ are the most important channels for the $\Lambda(1405)$: these two channels dynamically generate the quasi bound $\Lambda(1405)$.

To understand the many body kaonic nuclear system, the $\bar{K}NN$ is the simplest system. The investigation of this system started in the 60's [12]. More recently, using the variational approach, a $\bar{K}NN$ state was found bound by 48 MeV with width of 60 MeV [13]. In [14], the K^-pp three body system was investigated using the variational approach taking the results of $\bar{K}N$ interactions from a chiral $SU(3)$ model, and they get a much less bound K^-pp state with, B=19 MeV and Γ=40-70 MeV. In a further paper, the same group quantifies the uncertainties and evaluates the extra with coming from $\bar{K}$ absorption from the pair of nucleons, with the results of B=20-40 MeV and Γ as large as 100 MeV with admitted large uncertainties [15] . Coupled channel three body calculations of the quasi bound $\bar{K}NN$ system were investigated in [16, 17] using Faddeev equations, and their results for the binding is 50-70 MeV and the width around 100 MeV. The work of [18] studied three body resonances in the $\bar{K}NN$ system within a framework of the $\bar{K}NN \rightarrow \pi YN$ coupled channel Faddeev equation. In this work they found the binding energy 79 MeV with width 74 MeV. However, in more recent papers [19, 20], where the energy dependence of the potential is taken from chiral dynamics, the same authors find much smaller binding energies, below 20 MeV. We also calculated the $\bar{K}NN$ system [21, 22] for the $S = 0$ and $S = 1$ case, using the Fixed Center Approximation (FCA) to the Faddeev equations taking into account $\pi\Sigma N$ channel explicitly and also including the charge exchange diagrams. We get the binding 26-35 MeV for $S = 0$ while around 9 MeV for $S = 1$ case. The width are in both cases around 50 MeV. A very recent Faddeev calculation [23] finds that the $S = 0$ state is bound by about 16 MeV and the $S = 1$ by less than 11 MeV. The widths, without including $\bar{K}$ absorption by a pair of nucleons, are around 40 MeV.

The FCA to the Faddeev equations is an efficient and useful method to investigate many body systems. There is substantial work to rely on this model. For instance, the three body $\bar{K}NK$ scattering amplitude [24] was calculated using the FCA to the Faddeev equations and the results of this work are in good agreement with the other theoretical works [25, 26] evaluated using variational and Faddeev approaches, respectively. Besides, in [27], using the same model, the authors give a plausible explanation for the $\Delta_{\frac{5}{2}+}(2000)$ puzzle. As important as to know the success of the FCA, it is to know the limitation of this procedure, and it was found in [28] that the approximation collapses for the study of resonant states, where there is plenty of phase space for the excitation of intermediate states.

Meson nucleus bound systems are good laboratories to investigate the finite density QCD. The existence of the $\bar{K}$ bound states in nuclear systems, which are called kaonic nuclei, is theoretically expected because of the strong attractive interactions between the $\bar{K}$ and the nucleon. Models range from empirical ones with a depth of around 200 MeV at normal nuclear matter density [29], to those imposing chiral symmetry constraints that lead to an attraction of around 40-50 MeV [30–34]. However there is no experimental evidence on this system. In spite of experimental claims, which have been proved to be unfounded (see Ref. [35] for a review on the issue). The large width of the predicted states compared to the binding could justify why such states are not being found [36]. The K^-pp is the prototype of the $\bar{K}$ nuclei and to observe the kaonic nuclei in experiments, the precise knowledge of the width of this state is important. In this sense, it is important to recall that in all previous works of the $\bar{K}NN$ system the source of the width is only the $\bar{K}N \rightarrow \pi\Sigma$ conversion channel. The absorption channels $\bar{K}NN \rightarrow \Lambda N, \Sigma N$ are not

considered explicitly, although in some case it is estimated perturbatively [15]. Early experiments on K^--nucleus absorption at rest were done in the 1970s. A precise measurement of the total two nucleon absorption of stopped K^- mesons in deuterium was reported in [37] using deuterium bubble chambers. This is the first detailed measurements of two nucleon absorption branching ratio and these results were compared with the predictions of isospin invariance for strong interaction. The ratio of the rates $R(K^-d \to \Sigma^-p)$ and $R(K^-d \to \Sigma^0n)$ was in good agreement with the results of isospin invariance for strong interaction. A Similar experiment was done using a helium bubble chamber [38].

$\bar{K}$ absorption by two nucleons in nuclei, $\bar{K}NN \to \Lambda N, \Sigma N, \Sigma^*N$, has been considered before in the selfconsistent evaluation of the $\bar{K}$ nucleon optical potential of [32, 39]. More recently it has also been evaluated in a different way in [40]. The $\bar{K}$ nucleus optical potential evaluated in [32] has been checked versus kaonic atoms in [36] with good agreement with experiment. The discrepancies seen for the ^4He atom have been revolved recently with very precise measurement in [41, 42] which agree with the early predictions of [36]. Also a best fit to the data done in [43] showed that a best fit to the data could be obtained with a potential that differed from the predicted one of [32] at the level of 20 %. This agreement with data gives us confidence that the input used in [32] for the $\bar{K}NN$ absorption is realistic and we use this input here to evaluate the absorption width of the $\bar{K}NN$ state.

The article is organized as follows. In Section II, the calculation of the three body $\bar{K}NN$ amplitude including the charge exchange mechanisms is summarized using the Faddeev equations under the FCA. In Section III, The explicit derivation of the K^- absorption by two nucleons is given. The results of the K^- absorption both for spin-0 and spin-1 are shown in Section IV.

II.   CALCULATION OF THE $\bar{K}NN$ AMPLITUDE

Here we will summarize the derivation of the $\bar{K}NN$ three body amplitude within the framework of the FCA to the Faddeev equations taking into account the charge exchange contributions. In order to investigate the three-body amplitude of the $\bar{K}NN$, we have two possible spin states, $S = 0$ and $S = 1$. Let us first concentrate on the calculation of the $S = 0$ case which is done in detail in [21]. At the end we want total isospin-$\frac{1}{2}$ for three-body amplitude. The wave function for this state is

$$|K^-pp> = -(\frac{1}{\sqrt{3}}|3/2, 1/2> + \sqrt{\frac{2}{3}}|1/2, 1/2>) \tag{1}$$

with the basis of $|I_{tot}, I_{3,tot}>$. For the total amplitude with total isospin-$\frac{1}{2}$ we have

$$< 1/2|T|1/2> = \frac{3}{2}(< K^-pp|T|K^-pp> -\frac{1}{3}< 3/2|T|3/2>). \tag{2}$$

First, we start from the $K^-pp \to K^-pp$ amplitude where a K^- interacts with either of the two protons and the diagrammatic representation of this amplitude is shown in Fig. 1, where the shaded ellipse includes all multiple scattering of the K^- where the K^- interacts first with the first nucleon. We call the result of this diagram T_p and hence the total amplitude including the diagrams where the K^- interacts first with the second nucleon is $2T_p$. Looking in detail at the ellipse in Fig. 1 we have three partition functions (see Fig. 2) that fulfill the following coupled channel equations

$$\begin{aligned} T_p &= t_p + t_pG_0T_p + t_{ex}G_0T_{ex}^{(p)} \\ T_{ex}^{(p)} &= t_0^{(p)}G_0T_{ex}^{(n)} \\ T_{ex}^{(n)} &= t_{ex} + t_{ex}G_0T_p + t_0^{(n)}G_0T_{ex}^{(p)} \end{aligned} \tag{3}$$
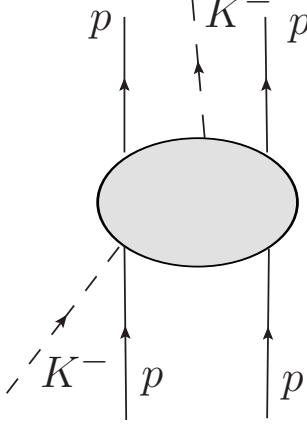
FIG. 1: Diagrammatic representation of the $K^-(pp)$ interaction. The equivalent diagram where K^- first interacts with the second proton should be added.

where $t_p = t_{K^-p,K^-p}$, $t_{ex} = t_{K^-p,\bar{K}^0n}$, $t_0^{(p)} = t_{\bar{K}^0p,\bar{K}^0p}$, $t_0^{(n)} = t_{\bar{K}^0n,\bar{K}^0n}$ and G_0 [44, 45] is

$$G_0 = \int \frac{d^3q}{(2\pi)^3} F_{NN}(q) \frac{1}{q^{02} - \vec{q}^2 - m_{\bar{K}}^2 + i\epsilon}, \tag{4}$$

with the factor $F_{NN}(q)$ standing for the form factor of the bound NN cluster. The diagrammatic representation of the partition functions is shown in Fig. 2. Calculating the three equations in Eq. (3) we get T_p as below

$$T_p^{(1/2)} = \frac{t_p(1 - t_0^{(n)}G_0 t_0^{(p)}G_0) + t_{ex}^2 G_0 t_0^{(p)}G_0}{(1 - t_p G_0)(1 - t_0^{(n)}G_0 t_0^{(p)}G_0) - t_{ex}^2 t_0^{(p)}G_0^3}. \tag{5}$$

Now we need to evaluate the second term of Eq. (2). This term does not have charge exchange and its diagrammatic representation is shown in Fig. 3. Taking into account the equivalent diagram where $\bar{K}^0$ interacts first with the second nucleon we obtain

$$T_p^{(3/2)} = 2\frac{t_0^{(p)}}{1 - G_0 t_0^{(p)}}. \tag{6}$$

Using Eq. (2), the final result of the amplitude for $S = 0$ case is

$$T^{(1/2)} = 3T_p - \frac{t_0^{(p)}}{1 - G_0 t_0^{(p)}}. \tag{7}$$

In the case of $S = 1$, the amplitude of the three-body system is calculated in Ref. [22] following the steps of [46] for the evaluation of the $\bar{K}^-d$ scattering length. The resulting amplitude is

$$T_{K^-d} = \frac{t_p + t_n + (2t_p t_n - t_x^2)G_0 - 2t_x^2 t_n G_0^2}{1 - t_p t_n G_0^2 + t_x^2 t_n G_0^3} \tag{8}$$
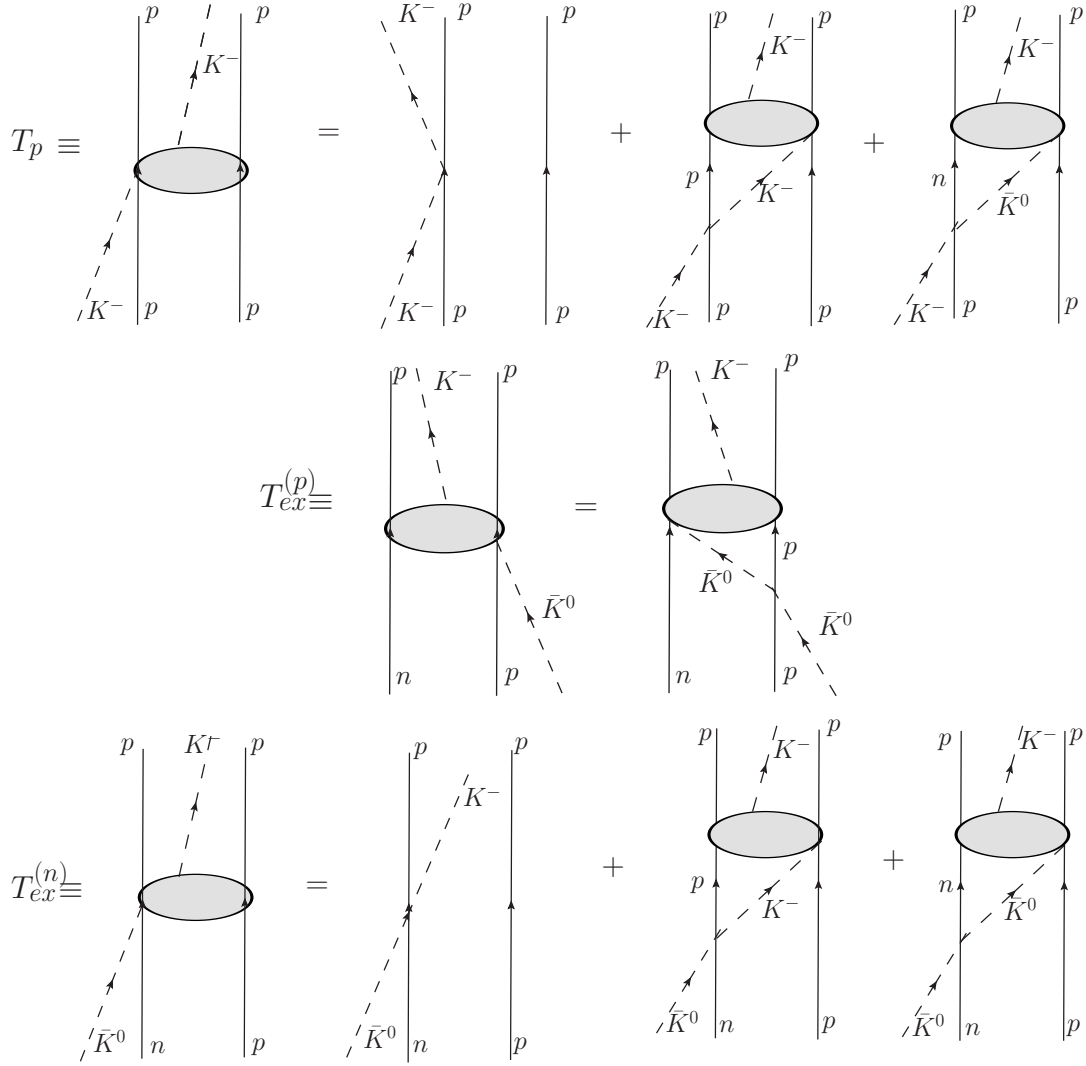
where $t_x = t_{ex}/\sqrt{1 + t_0^{(n)}G_0}$.

FIG. 2: Diagrammatic representation of the partition functions for $K^-pp \to K^-pp$.
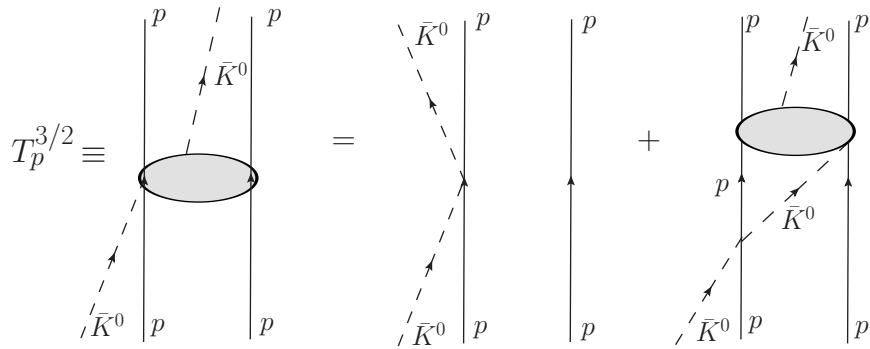

FIG. 3: Diagrammatic representation of the partition function for I=3/2.

III. $\bar{K}NN$ ABSORPTION

We are going to formulate the $\bar{K}$ absorption by two nucleons. The Feynman diagrams for this process are shown in Fig. 4. The S-matrix elements for these diagrams are given as follows:

$$
\begin{aligned}
S &= \frac{1}{V^2} \int \frac{d^3q}{(2\pi)^3} \frac{1}{\sqrt{2\omega_{\bar{K}}}} t_{\bar{K}N \to \bar{K}N} \frac{1}{q^2 - m_{\bar{K}}^2 + i\epsilon} \\
&\quad \times V_{y_{Z(H)}} A(B)\, \tilde{\varphi}(\vec{q} - \vec{p}_{Z(H)} + \frac{\vec{P}}{2})(2\pi)^4 \delta^4(p_i - p_f) \\
&\equiv -iT \frac{1}{\sqrt{2\omega_{\bar{K}}}} \frac{1}{V^2}(2\pi)^4 \delta^4(p_i - p_f),
\end{aligned}
\tag{9}
$$

where $V_{y_{Z(H)}} A(B)$ is the meson baryon Yukawa vertex, Z stands for the $\Lambda, \Sigma^0, \Sigma^+$ spin-$\frac{1}{2}$ octet baryons, H represents the Σ^{*0}, Σ^{*+} spin-$\frac{3}{2}$ decuplet baryons, A is equal to $\vec{\sigma}\vec{q}$ for the $\bar{K}NZ$ vertex and B is equal to $\vec{S}^\dagger \vec{q}$ for the $\bar{K}NH$ vertex where $\vec{S}^\dagger$ is the spin transition operator from spin-$\frac{1}{2}$ to spin-$\frac{3}{2}$. Here $\tilde{\varphi}(\vec{q})$ is the wave function of the pair of nucleons and p_i and p_f are the initial and final momenta, respectively.

Taking the NN system at rest, we can make some approximation for the $\bar{K}$ propagator as below

$$
\frac{1}{q^2 - m_{\bar{K}}^2} \to \frac{1}{(q^0)^2 - \vec{p}_{Z(H)}^2 - m_{\bar{K}}^2},
\tag{10}
$$

where $q_{Z(H)}^0 = E_{Z(H)} - E_N$ with $E_{Z(H)} = (M_{\bar{K}NN}^2 + M_{Z(H)}^2 - M_N^2)/2M_{\bar{K}NN}$.

After redefinition of $\vec{q}\,' \equiv \vec{q} - \vec{p}_{Z(H)}$, the square of the total matrix element summed over the spins of the Z, or H, and averaged over the spin of the nucleons is obtained as

$$
\begin{aligned}
|T|^2 &= V_{y_{Z(H)}}^2 C_{Z(H)} \vec{p}_{Z(H)}^2 \left(\frac{1}{(q^0)^2 - p_{Z(H)}^2 - m_{\bar{K}}^2} \right)^2 \\
&\quad \times \left| \frac{1}{2\pi^2} \int q'^2 dq' \tilde{\varphi}(\vec{q}\,') t_{\bar{K}N,\bar{K}N}(\sqrt{s'}) \right|^2.
\end{aligned}
\tag{11}
$$

with $C_Z = 1$ and $C_H = 2/3$. Using the coefficients in Table I and Eq. (14) in Ref. [39], we have for the coefficients of the Yukawa vertex

$$
\begin{aligned}
V_{yz} &= -\frac{1}{\sqrt{3}} \frac{3F + D}{2f} \quad \text{for} \ \ K^- p \to \Lambda \\
&= \frac{D - F}{2f} \quad \text{for} \ \ K^- p \to \Sigma^0 \\
&= \sqrt{2} \frac{D - F}{2f} \quad \text{for} \ \ \bar{K}^0 p \to \Sigma^+
\end{aligned}
\tag{12}
$$

with $D = 0.795$ and $F = 0.465$. Similarly, for the $\bar{K}NH$ vertex we have

$$
V_{\bar{K}NH} = a \frac{g_H}{2M_N}
\tag{13}
$$

The coefficient a is given in Table 2 in Ref. [39] and the coupling $\frac{g_H}{2M_N}$ is given by,

$$\frac{g_H}{2M_N} = \frac{2\sqrt{6}}{5}\frac{D+F}{2f}.$$ (14)

Thus, the coefficients for the $\bar{K}NH$ Yukawa vertices are

$$V_{y_H} = \frac{2}{5}\frac{F+D}{2f} \quad \text{for} \quad K^-p \to \Sigma^{*0}$$

$$= \frac{2\sqrt{2}}{5}\frac{D+F}{2f} \quad \text{for} \quad \bar{K}^0p \to \Sigma^{*+}$$ (15)

It is useful to relate the T matrix with the cross section. Using the T matrix we calculate the $\bar{K}(NN)$ cross section as

$$\sigma_{\text{abs}} = \frac{1}{2\pi}\frac{M_{NN}M_{Z(H)}M_N}{M_{\bar{K}NN}^2}\frac{p_{Z(H)}}{p_{\bar{K}}}|T|^2.$$

Borrowing the optical theorem we rewrite the cross section in terms of the imaginary part of the $\bar{K}(NN) \to \bar{K}(NN)$ amplitude, $T_{\bar{K}(NN)}$, as follow:

$$\text{Im}\, T_{\bar{K}(NN)} = -\frac{p_{\bar{K}}\sqrt{s}}{M_{NN}}\sigma_{\text{abs}} = -\frac{1}{2\pi}\frac{M_{Z(H)}M_N}{M_{\bar{K}NN}}p_{Z(H)}|T|^2.$$ (16)

From Eq. (16), the imaginary part of the $T_{\bar{K}(NN)}$ is written as

$$\text{Im}\, T_{\bar{K}(NN)} = -\frac{1}{2\pi}\frac{M_N}{M_{\bar{K}NN}}M_{Z(H)}p_{Z(H)}|T_{Z(H)}|^2.$$ (17)

In analogy to [47], we can convert the absorption diagram of Fig. 4 into the "many-body" diagram of Fig. 5 where the nucleon on which the virtual $\bar{K}$ is absorbed is converted into a "hole" line. The purpose of following this path is that we can now include the absorption process by making a modification of the meson baryon loop function, δG, to include the "ph" excitation in the $\bar{K}$ propagator. Then, we can reevaluate the $\bar{K}N$ amplitude in the coupled channels nonperturbative Bethe Salpeter equations and use the new amplitude in the $\bar{K}(NN)$ amplitudes of Eqs. (7) and (8) which sum the multiple scattering series of the $\bar{K}(NN)$ system. In this way we perform fully nonperturbatively the implementation of the $\bar{K}(NN)$ absorption in the $\bar{K}(NN)$ system.

Since we are only concerned about the absorption process it suffices to evaluate $\text{Im}\,\delta G$. This is easy since $\text{Im}\,\delta G$ is the same as $\text{Im}\,T_{\bar{K}(NN)}$ removing the two $t_{\bar{K}N,\bar{K}N}$ amplitudes in Eq. (11). Thus we can immediately write

$$i\text{Im}\,\delta G_{K^-p} = -i\frac{1}{2\pi}\frac{M_N}{M_{\bar{K}NN}}\sum_{i=1}^{3}M_{Z(H)}p_{Z(H)}|\tilde{T}_{Z(H)}|^2.$$ (18)

$$i\text{Im}\,\delta G_{\bar{K}^0n} = -i\frac{1}{2\pi}\frac{M_N}{M_{\bar{K}NN}}\sum_{i=1}^{2}M_{Z(H)}p_{Z(H)}|\tilde{T}_{Z(H)}|^2.$$ (19)

where the sum for 1-3 of δG_{K^-p} runs over Λ, Σ^0 for type Z and Σ^{*0} for type H, and the sum 1-2 of $G_{\bar{K}^0n}$ runs over Σ^+ of type Z and Σ^{*+} of type H (see Fig. 5).
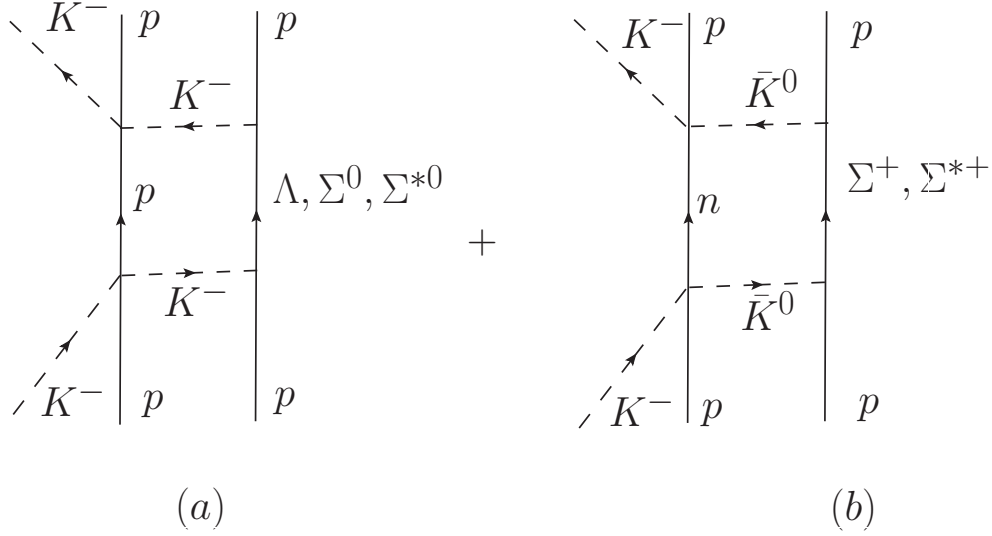
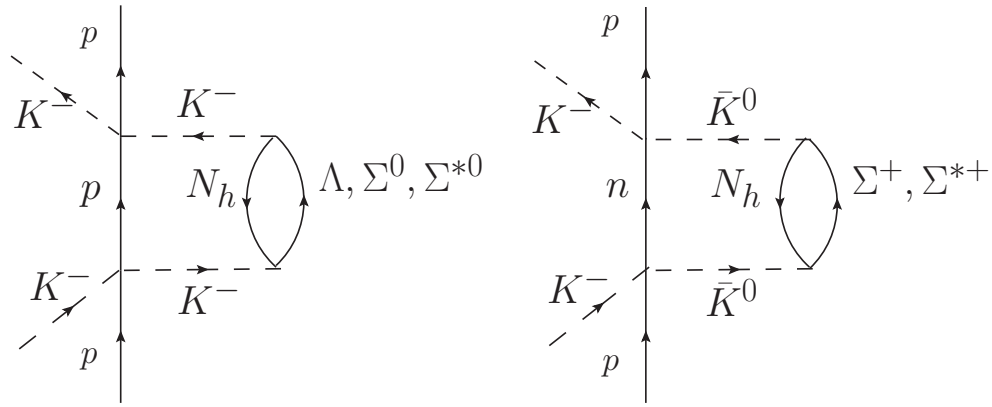FIG. 4: Diagrammatic representation of the the $K^-(pp)$ absorption.

FIG. 5: Representation of the $K^-(pp)$ absorption in many-body diagrams.

In Eqs. (18) and (19), $|\tilde{T}_{Z(H)}|^2$ is given by

$$|\tilde{T}_{Z(H)}|^2 = |\varphi(0)|^2 V_{y_{Z(H)}}^2 C_{Z(H)} \vec{p}_{Z(H)}^{\,2} \left(\frac{1}{(q_i^0)^2 - p_{Z(H)}^2 - m_{\bar{K}}^2} \right)^2. \tag{20}$$

where V_{y_Z}, V_{y_H}, are given in Eqs. (12) and (15), respectively. Note that when removing the $t_{\bar{K}N,\bar{K}N}$ amplitude in the integral of Eq. (11), the remaining integral simply gives the wave function in coordinate space in the origin, $\varphi(0)$.

For $\varphi(r)$ we take the same wave function as in [47] for the NN cluster

$$\varphi(r) = ae^{-\alpha r}, \quad a = \frac{1}{2}\left(\frac{\alpha^3}{2\pi}\right)^{\frac{1}{2}}, \tilde{\varphi}(q) = \frac{4\pi a\alpha}{(\frac{1}{4}\alpha^2 - q^2)^2 + q^2\alpha^2}. \tag{21}$$

The momentum of the baryons are given by

$$p_{Z(H)} = \frac{\lambda^{\frac{1}{2}}(M_{\bar{K}NN}^2, M_N^2, M_{Z(H)}^2)}{2M_{\bar{K}NN}} \tag{22}$$

and in order to account for relativistic corrections, $\vec{p}_{Z(H)}^{\,2}$ accompanied by $V_{y_{Z(H)}}^2$ is given by

$$\vec{p}_{Z(H)}^{\,2} V_{y_{Z(H)}}^2 \longrightarrow V_{y_{Z(H)}}^2 \frac{1}{4M_{Z(H)}^2}(M_N + M_{Z(H)})^2 \vec{p}_{Z(H)}^{\,2} \tag{23}$$

Since we get different contributions for δG in the K^-p or $\bar{K}^0n$ intermediate states, we are now forced to recalculate the $\bar{K}N$ amplitudes using the charge basis of the coupled channels, rather than the isospin base normally used.
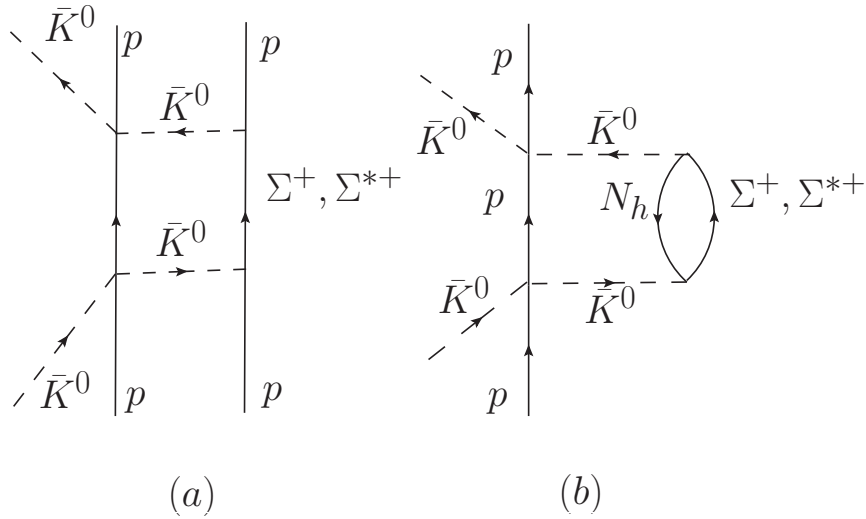

FIG. 6: Representation of the $K^-(pp)$ absorption for $t_{\bar{K}^0p \to \bar{K}^0p}$ amplitude.

Furthermore, as seen in Eqs. (3), we also need to calculate the $t_{\bar{K}^0p \to \bar{K}^0p}$ amplitude. This is a channel that has only one Feynman diagram as shown in Fig. 6. In order to calculate the $t_{\bar{K}^0p \to \bar{K}^0p}$ amplitude we use the same potential with V_{K^-n, K^-n} [48]. The $\bar{K}^0p$ channel is renormalized in the same way as the $\bar{K}^0n$, as one can see, comparing Fig. 6 (a) with Fig. 4 (b). Hence, one can use $G_{\bar{K}^0p} \to G_{\bar{K}^0p} + i\delta G_{\bar{K}^0n}$ for the channel $\bar{K}^0p$.

IV. RESULTS

As we mentioned in the Introduction, the $\Lambda(1405)$ plays a key role in the $\bar{K}$ absorption. As mentioned before, we recalculate the $\bar{K}N$ amplitudes in the charge base. In the case of the $Q = 0$ and $I_3 = 0$, there are ten coupled channels which are K^-p, $\bar{K}^0n$, $\pi^0\Lambda$, $\pi^0\Sigma^0$, $\eta\Lambda$, $\eta\Sigma^0$, $\pi^+\Sigma^-$, $\pi^-\Sigma^+$, $K^+\Xi^-$, $K^0\Xi^0$ [48]. As discussed above, we also need the $\bar{K}^0p \to \bar{K}^0p$ amplitude in pure $I = 1$. For this we take the $I_3 = -1$ component and use the six coupled channels, K^-n, $\pi^0\Sigma^-$, $\pi^-\Sigma^0$, $\pi^-\Lambda$, $\eta\Sigma^-$, $K^0\Xi^-$ [48]. In the coupled channels reevaluation of the $\bar{K}N$ amplitudes for $I_3 = 0$ we add $i\delta G_{K^-p}$, $i\delta G_{\bar{K}^0n}$ to G_{K^-p}, $G_{\bar{K}^0n}$, respectively. For $I_3 = 1$ we add $i\delta G_{K^-n}$ to G_{K^-n}, but for isospin symmetry reasons $i\delta G_{K^-n} \equiv i\delta G_{\bar{K}^0p}$ which we have discussed above. In all the channels we use the same numeric value for f, $f = 1.123f_\pi$

(f_π=93 MeV) as in [48]. For the form factor of Eq. (4) we use the deuteron type form factor with the reduced size of the two N system found in [15].

In Figs. 7 and 8 we show the absolute squared of the T matrix for $\bar{K}$ scattering on the NN cluster, including absorption, for the cases of $S = 0$ and $S = 1$. As we can see from Fig. 7 the most striking thing is a substantial increase in the width, that goes from about 50 MeV to about 80 MeV, and which is due to absorption. The centroid of the distribution is also a bit displaced to lower energies, but what concerns us now is that the $\bar{K}$ absorption on two nucleons has increased the width by about 30 MeV. The order of magnitude is similar to what was estimated in [14, 15], but here we have done a full nonperturbative evaluation of this magnitude. In the case of $S = 1$ in Fig. 8 we observe that the shape of the distribution has been distorted considerably due to the consideration of the $\bar{K}$ absorption and the centroid has not changed appreciably. From this distorted shape one can also estimate that the width has increased in about 30 MeV, like in the previous case.

We should mention that we have also included the effect of the $\pi\Sigma N$ intermediate states in the calculation. Here also the $\bar{K}N \to \pi\Sigma$ amplitudes have been modified due to absorption, since they come from the coupled channels calculation. The contribution of this channel has been done as in [45] and its effect is small, like also found there.
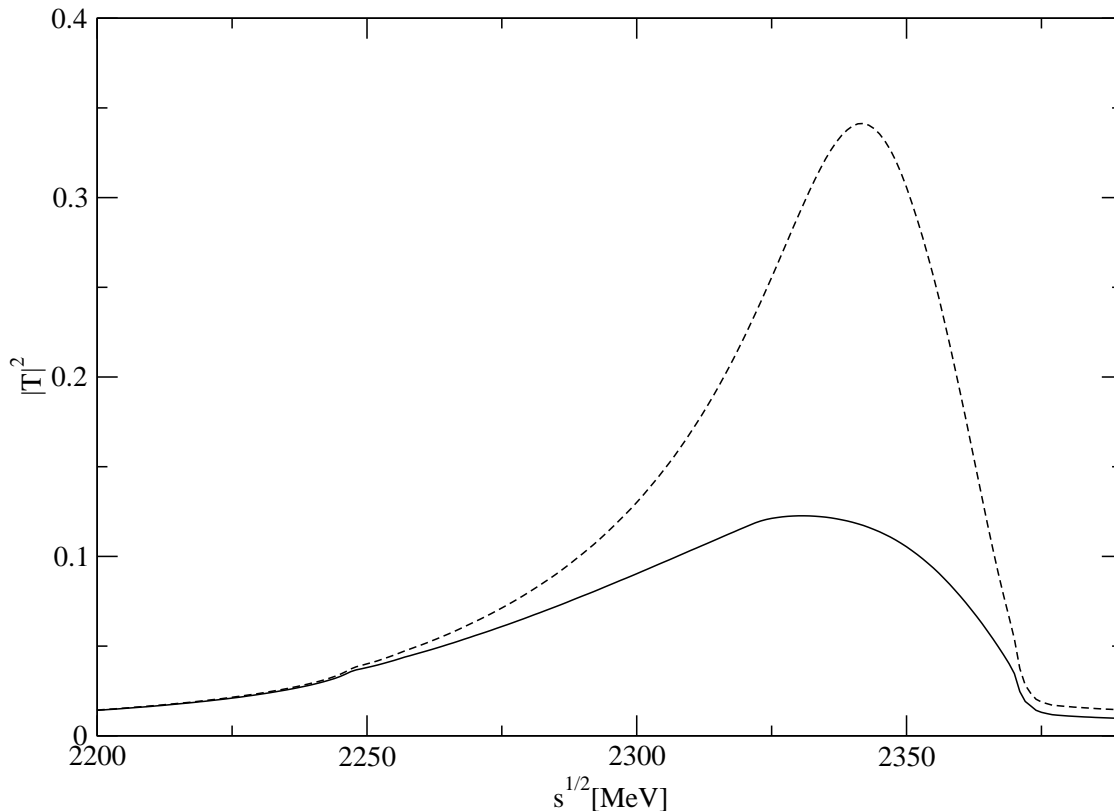

FIG. 7: Modulus squared of the T matrix for $\bar{K}$ scattering on the NN cluster for $S = 0$.
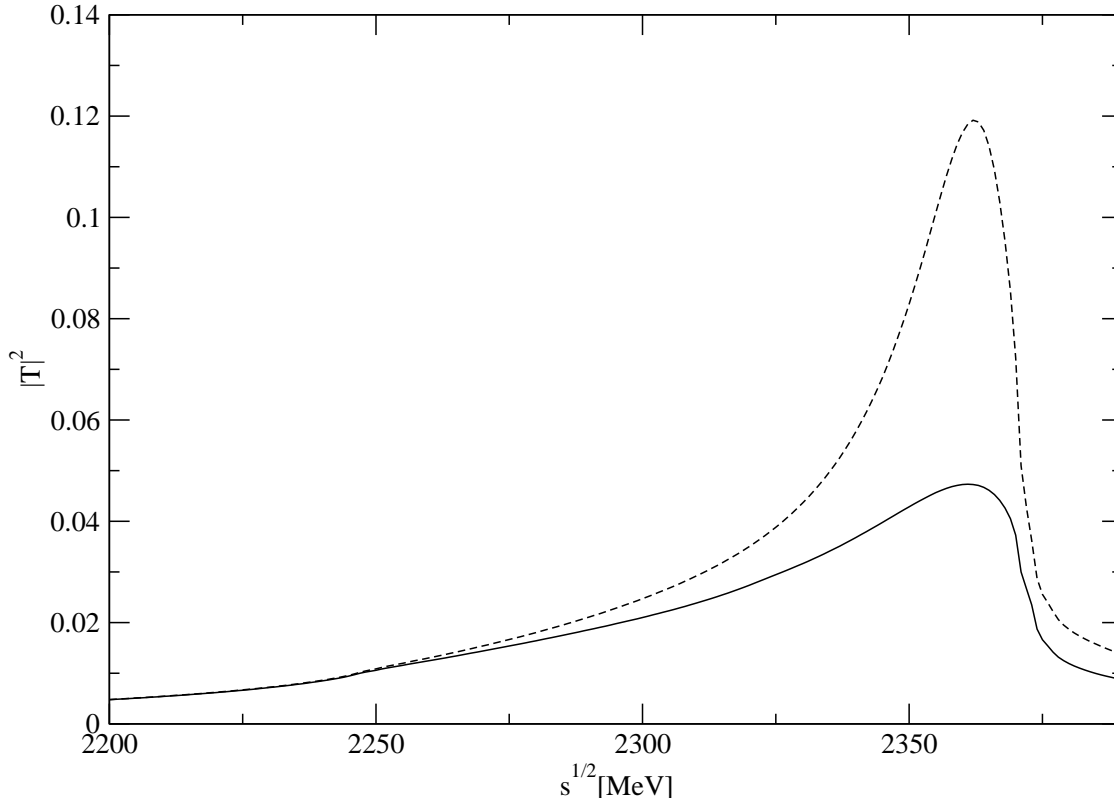
FIG. 8: Modulus squared of the T matrix for $\bar{K}$ scattering on the NN cluster for $S = 1$.

V. CONCLUSIONS

We have made a detailed and accurate calculation of the contribution to the width of the bound $\bar{K}NN$ states from $\bar{K}$ absorption on two nucleons. The evaluation is done nonperturbatively in two aspects: First, the $\bar{K}N$ amplitudes are reevaluated in the unitary coupled channels approach taking into account the absorption of the $\bar{K}$. Second the resulting $\bar{K}N$ amplitudes are used in the nonperturbative formula of the Fixed Center Approximation that takes into account the rescattering of the kaons on the nucleons of the NN cluster.
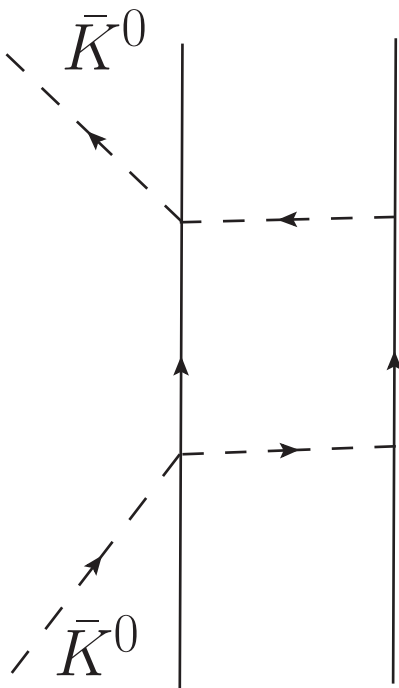
The result of these calculations is that the width of the states with $S = 0$, $S = 1$ is increased by about 30 MeV to values of the total width of 75-80 MeV. With this large width and the small values of the binding, 15-30 MeV, to which the different groups are converging [15, 20, 23], we are facing a situation of states with much larger width than binding, which makes the experimental observation problematic. Further calculations taking advantage of the steps given in the present paper, but using different formalisms, would be most welcome.
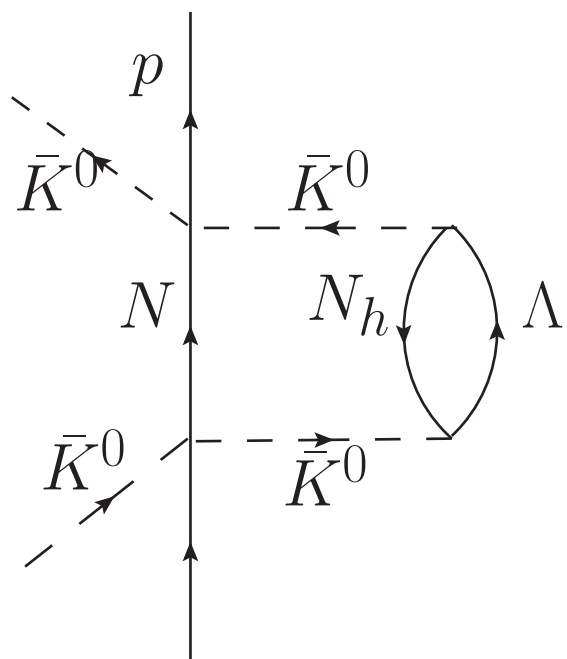
VI. ACKNOWLEDGMENTS

[1] N. Kaiser, P. B. Siegel and W. Weise, Nucl. Phys. A **594**, 325 (1995).

[2] E. Oset and A. Ramos, Nucl. Phys. A **635**, 99 (1998).

[3] J. A. Oller and U. G. Meissner, Phys. Lett. B **500**, 263 (2001).

[4] M. F. M. Lutz and E. E. Kolomeitsev, Nucl. Phys. A **700**, 193 (2002).

[5] E. Oset, A. Ramos and C. Bennhold, Phys. Lett. B **527**, 99 (2002) [Erratum-ibid. B **530**, 260 (2002)].

[6] T. Hyodo, S. I. Nam, D. Jido and A. Hosaka, Phys. Rev. C **68**, 018201 (2003).

[7] D. Jido, J. A. Oller, E. Oset, A. Ramos and U. G. Meissner, Nucl. Phys. A **725**, 181 (2003).

[8] B. Borasoy, R. Nissler and W. Weise, Eur. Phys. J. A **25**, 79 (2005).

[9] J. A. Oller, Eur. Phys. J. A **28**, 63 (2006) [hep-ph/0603134].

[10] B. Borasoy, U. -G. Meissner and R. Nissler, Phys. Rev. C **74**, 055201 (2006) [hep-ph/0606108].

[11] Y. Ikeda, T. Hyodo and W. Weise, Nucl. Phys. A **881**, 98 (2012) [arXiv:1201.6549 [nucl-th]].

[12] Y. Nogami, Phys. Lett. 7, 288 (1963).

[13] T. Yamazaki and Y. Akaishi, Phys. Lett. B **535** (2002) 70.

[14] A. Dote, T. Hyodo and W. Weise, Nucl. Phys. A **804**, 197 (2008).

[15] A. Dote, T. Hyodo and W. Weise, Phys. Rev. C **79**, 014003 (2009) [arXiv:0806.4917 [nucl-th]].

[16] N. V. Shevchenko, A. Gal and J. Mares, Phys. Rev. Lett. **98**, 082301 (2007) [arXiv:nucl-th/0610022].

[17] N. V. Shevchenko, A. Gal, J. Mares and J. Revai, Phys. Rev. C **76**, 044004 (2007)

[18] Y. Ikeda and T. Sato, Phys. Rev. C **76**, 035203 (2007) [arXiv:0704.1978 [nucl-th]].

[19] Y. Ikeda and T. Sato, Phys. Rev. C **79**, 035201 (2009) [arXiv:0809.1285 [nucl-th]].

[20] Y. Ikeda, H. Kamano and T. Sato, Prog. Theor. Phys. **124**, 533 (2010) [arXiv:1004.4877 [nucl-th]].

[21] M. Bayar and E. Oset, Nucl. Phys. A **883**, 57 (2012) [arXiv:1203.5313 [nucl-th]].

[22] E. Oset, D. Jido, T. Sekihara, A. M. Torres, K. P. Khemchandani, M. Bayar and J. Yamagata-Sekihara, Nucl. Phys. A **881**, 127 (2012) [arXiv:1203.4798 [hep-ph]].

[23] N. Barnea, A. Gal and E. Z. Liverts, Phys. Lett. B **712**, 132 (2012) [arXiv:1203.5234 [nucl-th]].

[24] J. -J. Xie, A. Martinez Torres and E. Oset, Phys. Rev. C **83**, 065207 (2011) [arXiv:1010.6164 [nucl-th]].

[25] D. Jido and Y. Kanada-En'yo, Phys. Rev. C **78**, 035203 (2008) [arXiv:0806.3601 [nucl-th]].

[26] A. Martinez Torres, K. P. Khemchandani and E. Oset, Phys. Rev. C **79**, 065207 (2009) [arXiv:0812.2235 [nucl-th]].

[27] J. -J. Xie, A. Martinez Torres, E. Oset and P. Gonzalez, Phys. Rev. C **83**, 055204 (2011) [arXiv:1101.1722 [nucl-th]].

[28] A. Martinez Torres, E. J. Garzon, E. Oset and L. R. Dai, Phys. Rev. D **83**, 116002 (2011) [arXiv:1012.2708 [hep-ph]].

[29] E. Friedman and A. Gal, Phys. Lett. B **459**, 43 (1999) [nucl-th/9902036].

[30] V. Koch, Phys. Lett. B **337**, 7 (1994) [nucl-th/9406030].

[31] M. Lutz, Phys. Lett. B **426**, 12 (1998) [nucl-th/9709073].

[32] A. Ramos and E. Oset, Nucl. Phys. A **671**, 481 (2000) [nucl-th/9906016].

[33] J. Schaffner-Bielich, V. Koch and M. Effenberger, Nucl. Phys. A **669**, 153 (2000) [nucl-th/9907095].

[34] A. Cieply, E. Friedman, A. Gal and J. Mares, Nucl. Phys. A **696**, 173 (2001) [nucl-th/0104087].

[35] A. Ramos, V. K. Magas, E. Oset and H. Toki, Nucl. Phys. A **804**, 219 (2008).

[36] S. Hirenzaki, Y. Okumura, H. Toki, E. Oset and A. Ramos, Phys. Rev. C **61**, 055205 (2000).

[37] V. R. Veirs and R. A. Burnstein, Phys. Rev. D **1**, 1883 (1970).

[38] P. A. Katz, K. Bunnell, M. Derrick, T. Fields, L. G. Hyman and G. Keyes, Phys. Rev. D **1**, 1267 (1970).

[39] E. Oset and A. Ramos, Nucl. Phys. A **679**, 616 (2001) [nucl-th/0005046].

[40] T. Sekihara, J. Yamagata-Sekihara, D. Jido and Y. Kanada-En'yo, arXiv:1204.3978 [nucl-th].

[41] S. Okada, G. Beer, H. Bhang, M. Cargnelli, J. Chiba, S. Choi, C. Curceanu and Y. Fukuda *et al.*, Phys. Lett. B **653**, 387 (2007) [arXiv:0707.0448 [nucl-ex]].

[42] M. Bazzi *et al.* [SIDDHARTA Collaboration], Phys. Lett. B **681**, 310 (2009).

[43] A. Baca, C. Garcia-Recio and J. Nieves, Nucl. Phys. A **673**, 335 (2000) [nucl-th/0001060].

[44] L. Roca and E. Oset, Phys. Rev. D **82**, 054013 (2010).

[45] M. Bayar, J. Yamagata-Sekihara and E. Oset, Phys. Rev. C **84**, 015209 (2011) [arXiv:1102.2854 [hep-ph]].

[46] S. S. Kamalov, E. Oset and A. Ramos, Nucl. Phys. A **690**, 494 (2001).

[47] M. Bayar, C. W. Xiao, T. Hyodo, A. Dote, M. Oka and E. Oset, arXiv:1205.2275 [hep-ph].

12

[48] E. Oset and A. Ramos, Nucl. Phys. A **635**, 99 (1998) [nucl-th/9711022].

(a) (b)