Vol. 39 (2008) *ACTA PHYSICA POLONICA B* No 2

ON NUCLEAR STATES OF STRANGE MESONS*

Sławomir Wycech†,

The A. Sołtan Institute for Nuclear Studies
Hoża 69, 00-681 Warsaw, Poland

Anthony M. Green‡

Helsinki Institute of Physics, P.O. Box 64, 00014, Finland

(Received January 3, 2008)

The search for nuclear states of K-mesons and η-mesons is presented. Methods of theoretical descriptions and the related difficulties: off shell extrapolation of meson–nucleon scattering amplitudes, behavior of hadronic resonances in nuclei and extrapolation to high density nuclear regions are discussed. Variational calculations for the binding energies in light nuclei are described.

PACS numbers: 13.75.Jz, 21.45.+v, 25.10.+s, 25.80.–e

1. Introduction

States involving mesons bound in nuclei have a long history. The first examples were μ and π atomic states. The muons were a valuable instrument to study the nuclear shape and nuclear dynamics. The pions gave information on πN interactions, the behavior of π's in a nuclear medium and, in particular, on the question of chiral symmetry restoration (see ref. [1]). Beginning with loosely bound atomic states, studied by X-ray measurements, one has now reached deeply bound states formed by the combination of Coulomb and strong interactions. This history has been repeated with K-mesons. In addition to atomic states, recent experiments have indicated that there could be systems where K-mesons are deeply bound by nuclear forces. Such systems are likely to be strongly compressed opening a way to study new physics involving multi-quark configurations. Another mechanism to form

* Presented at the XXX Mazurian Lakes Conference on Physics, Piaski, Poland, September 2–9, 2007.

† wycech@fuw.edu.pl

‡ anthony.green@helsinki.fi

such states is connected to excitations of nuclear resonances. This along with the η-nucleus bound states are the topics discussed in this report. These systems of open-strangeness, in the K case, and hidden-strangeness in the η case have a common mechanism for the nuclear attraction. It is related to baryon resonant states: the strange $\Lambda(1405)$ and non-strange $N(1535)$, respectively.

1.1. Nuclear states of the η-meson

The first experimental attempts to find an η bound to a nucleus [2] were unsuccessful apparently due to large experimental backgrounds or large widths of such states. The first signal came from light nuclei. In the formation reaction $pd \rightarrow ^3\mathrm{He}\,\eta$ a strong enhancement of the amplitude is observed close to the η production threshold [3]. A similar, although weaker effect is seen in the $pn \rightarrow ^2\mathrm{D}\,\eta$ reaction, [4]. The information obtained from these reactions is rather indirect and based on final state interactions of the η. Roughly, such interactions are described by a factor

$$\mid T_\eta \mid^2 = \left| \frac{F_\eta}{1 - i q_\eta A_\eta} \right|^2 , \tag{1}$$

where F_η is a formation amplitude, q_η is the momentum and A_η is the scattering length of η on the residual nucleus. Strong threshold enhancements happen when scattering lengths are large, and this signals states of small binding. Measurements of cross sections permit one to determine the modulus of T_η and also to extract the modulus of A_η. The $^3\mathrm{He}\,\eta$ experiments yield $\mid A_\eta \mid \simeq 5\,\mathrm{fm}$ [5]. Such a length is large compared to the size of the nucleus and indicates a quasi-bound (if Re $A_\eta < 0$) or a quasi-virtual state (if $\mathrm{Re}\,A_\eta > 0$). The first case presents an analog of the deuteron, the second case represents an anti-bound state known in the np spin singlet system in "ordinary" nuclear physics. In terms of Eq. (1) these cases correspond to different locations of the singularity in T_η. More detailed information on the nature of the bound state may be obtained from inelastic π^+ and π^0 production reactions which require a multichannel K-matrix description instead of Eq. (1). Along this line, the value $A_\eta = 4.24(29) + i0.72(81)\,\mathrm{fm}$ was obtained in Ref. [6]. It indicates the existence of a virtual state in the $^3\mathrm{He}\,\eta$ system and also tells that the absorption due to $\eta \rightarrow \pi$ conversion is rather weak. Lengths of this order are within the range of the contemporary understanding of ηN interactions based on the dominance of $N(1535)$, and suggest the existence of bound states in heavier nuclei. Recently, a new and dramatic evidence has been given by a (preliminary) Juelich result which determines cross sections much closer to the threshold [7]. The length now becomes $A_\eta = \pm 10.7(0.8) + i1.5(2.6)\,\mathrm{fm}$. Such a result, if confirmed, indicates a very narrow state, only a fraction of an MeV away from the threshold. In order to generate it in terms of an optical potential one would need very strong blocking of the main $\eta \to \pi$ decay process. Such a mechanism is indicated in the next section.

1.2. Nuclear states of the K-meson

It has been known since the first measurements of X-rays from K^- atoms that the optical potential used to describe atomic levels is attractive. The dominant term

$$U_K(r) = \frac{2\pi}{\mu_{KN}} f_{KN} \rho(r) \tag{2}$$

is parameterized by an effective scattering length f_{KN}. Here ρ denotes the nuclear density, μ_{KN} the reduced mass. It was soon realized that the attraction is related to the $\Lambda(1405)$, isospin 0, baryon located some 28 MeV below the KN threshold. This baryon may be a KN quasi-bound state, which decays into the $\pi\Sigma$ channel, or to some extent a genuine quark state. The K atoms are tested only in states of large atomic angular momenta. The meson barely grazes the nucleus and the energy of an average KN pair stays negative. The f_{KN} which is needed involves the subthreshold energy region

$$f_{KN} = f_{KN}(-E_B - E_{\mathrm{recoil}}) \,, \tag{3}$$

where E_B is the KN separation energy and E_{recoil} the recoil energy of the KN pair relative to the rest of the system. In atomic states the energies in Eq. (3) span the region from a few to about 50 MeV. There the $\Lambda(1405)$ dominated amplitude is roughly

$$f_{Kp}(E) \simeq \frac{\gamma^2}{E - E^*} \,, \tag{4}$$

where γ is a coupling constant and $E^* = E_{\mathrm{r}} - i\Gamma_{\mathrm{r}}/2$ is the complex resonance energy. Below the resonance energy E_{r} this amplitude is negative and the optical potential becomes attractive. This kind of attractive mechanism is operating also in the η case. The difference is that the resonance $N(1535)$ is of a quark nature and is located above the ηN threshold. Such a mechanism of attraction is the usual "level repulsion rule", found in atomic physics. The baryon resonance states are external to the meson nucleon scattering states. For energies below $E < E_{\mathrm{r}}$ the resonant states generate attraction while for $E > E_{\mathrm{r}}$ the effect is repulsive.

The optical potential for K-mesons is well tested only on the nuclear surface, optical potentials for η so far have not been tested by real experiments. However, there exist many theoretical or semi-phenomenological calculations. Two problems are met in such calculations: (1) Which way

the resonances behave in nuclear matter; (2) How to extrapolate to central nuclear densities.
A phenomenological approach is to expand the effective amplitude f_{KN} in terms of nuclear density, fit the parameters to atomic level data and extrapolate to the nuclear center. Such a method simulates the energy dependence inherent in $f_{KN}(E)$ and turns it into an effective density dependence [8]. It yields deep potentials $\mathrm{Re}\, U_K(0) \sim -180\,\mathrm{MeV}$ which could support bound states. This calculation is confirmed in a RMF study by the same group [9]. Another method uses the effective amplitude calculated in terms of the chiral meson–nucleon interaction embedded into a nuclear medium [10]. Some parameters have to be fitted to the atomic data. This leads to a shallow potential of $\mathrm{Re}\, U_K(0) \sim -50\,\mathrm{MeV}$. With both methods based on atomic data the absorptive parts $\mathrm{Im}\, U_K(0)$ are large about ~ 50–$100\,\mathrm{MeV}$. Thus deep potentials could generate bound but very broad states, the shallow potentials are unlikely to form bound states. A somewhat different approach was used in the calculations of Ref. [11]. There, a multichannel, phenomenological potential U, which yields reasonable atomic levels is used to calculate the energy levels of K-mesons in nuclear matter. The method solves multichannel scattering equations in the nuclear medium

$$\hat{f}(E) = \hat{V} + \hat{V} \frac{Q^{\mathrm{ex}}}{E - E^{\mathrm{int}}} \hat{f}(E)\,, \tag{5}$$

where Q^{ex} is the Pauli exclusion operator. Intermediate state energies $E^{\mathrm{int}} = E_N + U_K + M_K + E_K^{\mathrm{kinetic}}$ include complex energies determined by the self-consistent value of U_K. The on-shell energies are $E = E'_N + M_K - B_K$. This yields a value $\mathrm{Re}\, U_K(0) \sim -90\,\mathrm{MeV}$ consistent with the recent extraction of $\mathrm{Re}\, U_K(0) = -80\,\mathrm{MeV}$ obtained from heavy ion collisions [12]. Another result obtained from Eq. (5) is a $20\,\mathrm{MeV}$ upward shift of the $\Lambda(1405)$ in the nuclear surface region. The latter is supported by the nuclear emulsion data [13]. However, it has to be stressed that the extrapolation to central densities is model unstable and depends strongly on the form of E^{int} and the on-shell E introduced into Eq. (5). This problem is also related to the widths of nuclear states. It was pointed out in Ref. [11] that Kaon states may be narrow, since the main decay channel $\pi\Sigma$ is closed in the deeply bound states in large nuclei. The main decay mode becomes $KNN \to \Sigma N$ which was estimated to yield $\sim 20\,\mathrm{MeV}$ widths. This question is still open and requires further study.

In recent years, the existence of deeply bound states has been very actively discussed. It followed the discovery at KEK of $30\,\mathrm{MeV}$ wide peaks in the nucleon spectra following K^- absorption in $^4\mathrm{He}$, [14, 15]. Additional evidence was given by the FINUDA measurement of the invariant Λp mass distribution obtained in K^- absorption in light nuclei [16]. The existence of such bound states could have been expected in view of the kaonic atom experience. However, the KEK and FINUDA experiments indicate unexpectedly strong bindings of the order of 100–200 MeV in the lightest KNN and $KNNN$ systems. These experiments require confirmation, and the interpretation of the observed peaks has been disputed [17, 19]. On the other hand, calculations indicate that such states are expected, although they might be very broad and difficult to detect.

The first calculations performed by Akaishi and Yamazaki in Ref. [18] exploited the S-wave resonant attraction due to the $\Lambda(1405)$ state. With an optical model type of approach it was shown that $\mathrm{Re}\, U_K(0)$ at the center of small nuclei can be as strong as 500 MeV generating very strong bindings of the K-meson and strong contraction of the few-nucleon systems. To reproduce the KEK data, these calculations had to relax the NN repulsion at short distances which, only then, would allow the existence of strongly bound and very dense systems. The central part of such a system would now be better represented as a multi-quark state. These calculations raise an important question of how to implement the proper short range NN repulsion in the kaonic systems.

Another open question is related to the strength and range of KN interactions. The mathematical description of few body systems requires knowledge of NN and KN off-shell scattering amplitudes. Those related to NN interactions are controlled fairly well in terms of modern NN potentials. But for bound K-mesons the required amplitudes involve distant subthreshold energy regions. If the separation energy is as strong as 100 MeV, meson momenta become $\approx 250\,\mathrm{MeV}/c$ and E_{recoil} may be as large as 40 MeV. The energies of interest $(-E_B - E_{\mathrm{recoil}})$ are then located well below the $\Lambda(1405)$ state. The amplitudes there are not well known.

This paper uses a variational model to elucidate some of the problems in this field. The basic ingredients are:

1. To account properly for the KN force range and NN repulsion, a two step calculation is performed. First a wave function involving strongly correlated K–N subsystems is found in a fixed nucleon approximation. This step enables one to find potentials due to the K-meson which tend to contract the nucleons. Next, these correlated wave functions and contracting potentials are used as input in variational calculations for K-few-nucleon binding.

2. While the dominant mechanism of attraction is related to the $\Lambda(1405)$ state, it is found that another resonant state $\Sigma(1385)$ also contributes to the binding in $KNNN$, $KNNNN$ systems [23]. The latter state generates an additional branch of K-few-nucleon spectroscopy.

2. Variational approach

This section presents an introduction to the method. In the first step the meson binding energy and wave function is found with the nucleons fixed at coordinates $x_i (i = 1, 2, \ldots, n)$. In the next step this solution is used as an input to a variational calculation of the binding energy of the meson and n-nucleons. The method is presented in detail for a single KNN channel. It may be extended to a multichannel description in few-nucleon systems and with some changes it is applicable to the η-meson case.

Consider the scattering of a light meson on two identical, heavy nucleons fixed at coordinates $x_i (i = 1, 2)$. The wave function is assumed to be of the form

$$\Psi(x, x_1, x_2) = \psi_K(x, x_1, x_2)\chi_{NN}(x_1, x_2)\,, \tag{6}$$

where x is the meson coordinate. The meson wave function ψ_K is given by the solution of the multiple scattering equation

$$\psi_K(x, x_1, x_2) = \psi_K(x)^0 - \Sigma_i \int \boldsymbol{dy} \frac{\exp[ip \mid \boldsymbol{x} - \boldsymbol{y} \mid]}{4\pi \mid \boldsymbol{x} - \boldsymbol{y} \mid} U_{KN}(y, x_i)\psi_K(y, x_1, x_2) \tag{7}$$

obtained with fixed positions of the nucleons. A related procedure with a zero range meson–nucleon pseudo-potential U was used by Brueckner [20] to calculate the scattering length of a meson on two nucleons. At high energies this equation was extensively discussed by Foldy and Walecka [21] with separable interactions U. Here, the method is extended to the bound state problem. In the K-meson case one looks for solutions of Eq. (7) with no incident wave $\psi_K(x)^0$. The momentum $p \to p(x_i)$ becomes a complex eigenvalue which determines the energy and width of the system for given x_i.

Eq. (7) is written in terms of the Klein–Gordon propagator. This reflects on the normalization of the interaction $U_{KN} = 2\mu_{KN}V$. The potential V for an S-wave interaction is chosen in a separable form

$$V(\boldsymbol{x} - \boldsymbol{x_i}, \boldsymbol{x'} - \boldsymbol{x_i}) = \lambda_{\rm S} v_S(\boldsymbol{x} - \boldsymbol{x_i}) v_S(\boldsymbol{x'} - \boldsymbol{x_i})\,, \tag{8}$$

where v_S is the Yamaguchi form-factor and $\lambda_{\rm S}$ is a strength parameter. The eigenvalue equation now becomes

$$\psi_K(x) = -\Sigma_j \lambda_{\rm S} \int \boldsymbol{dy} \frac{\exp[ip \mid \boldsymbol{x} - \boldsymbol{y} \mid]}{4\pi \mid \boldsymbol{x} - \boldsymbol{y} \mid} v_S(\boldsymbol{y} - \boldsymbol{x_j}) \int \boldsymbol{dy'} v_S(\boldsymbol{y'} - \boldsymbol{x_j})\psi_K(\boldsymbol{y'})\,. \tag{9}$$

Eq. (9) may be reduced to a matrix equation for wave amplitudes ψ_i defined at each scatterer i

$$\psi_i = \int \boldsymbol{dx}\, v_S(\boldsymbol{x} - \boldsymbol{x_i})\psi_K(\boldsymbol{x}, \boldsymbol{x_1}, \boldsymbol{x_2})\,. \tag{10}$$

A few steps, described in some detail in Ref. [21], are necessary to bring Eq. (9) into standard multiple scattering equations for ψ_i. These are: (1) Integration over the i-th form-factor as in Eq. (10); (2) Selecting the i-th term from the right-hand side which has the form $-\lambda G$; (3) Inverting the diagonal in the i term of $1+\lambda G$. In this way the kernel of the multiple scattering equation can be expressed in terms of the S-wave scattering amplitudes $f^{\rm s}$ at each nucleon i and propagators describing the passage from nucleon i to the other nucleon j. The latter are given by

$$G_{i,j} = \int \boldsymbol{dydx}\, \boldsymbol{v}(\boldsymbol{x}-\boldsymbol{x}_i)\, \frac{\exp(ik \mid \boldsymbol{x}-\boldsymbol{y} \mid)}{4\pi \mid \boldsymbol{x}-\boldsymbol{y} \mid}\, \boldsymbol{v}(\boldsymbol{y}-\boldsymbol{x}_j)\,. \tag{11}$$

And so one arrives at a set of linear equations

$$\psi_i + \Sigma_{j\neq i} G_{i,j} f_j \psi_j = 0\,, \tag{12}$$

which may be solved numerically. For the sake of illustration we present the KNN case in some detail. The condition for a bound state with two amplitudes ψ_i is a 2×2 equation

$$\psi_1 + f^{\rm s}G^{ss}\psi_2 = 0\,, \qquad \psi_2 + f^{\rm s}G^{ss}\psi_1 = 0\,. \tag{13}$$

When the determinant $D = 1-(f^{\rm s}G^{ss})^2$ is set to zero, the binding "momenta" $p(x_i)$ may be obtained numerically. Two different solutions corresponding to $1+f^{\rm s}G^{ss}=0$ or $1-f^{\rm s}G^{ss}=0$ may exist. The first solution is symmetric $\psi_2=\psi_1$ and describes the meson in the S-wave state with respect to the NN center of mass. The second solution is antisymmetric $\psi_2=-\psi_1$ and, if it exists, describes a P-wave solution.

Eigenvalues corresponding to unstable quasi-bound states are obtained in the second quadrant of the complex $p(x_i) = p_{\rm R} + ip_{\rm I}$ plane. In this quadrant the kernel

$$f^{\rm s}G^{ss} = f^{s}_{KN}(p)[\exp(-p_{\rm I}r)\exp(ip_{\rm R}r) - \exp(-\kappa r) - r\frac{\kappa^2-p^2}{2\kappa}\exp(-\kappa r)]/r \tag{14}$$

is exponentially damped at large distances as required by the asymptotic form of the bound meson wave function. At short distances G^{ss} is regularized by the KN form-factor. It is essential for the stability of the method that the KN interaction range is very short. Therefore, uncertainties related to the actual value of κ are largely eliminated by the short range repulsion in NN systems.

In the specific case, where the scattering amplitude in the $\Lambda(1405)$ is dominated by the f_{Kp} of Eq. (4), the solution $1+f^sG^{ss}=0$ for the S-state becomes

$$E = E^* - \gamma^2 G^{ss}(r, p(E))\,. \tag{15}$$

Since $G^{ss}(r,p(E))$ close to resonance is positive, this solution results in binding stronger than the $\approx 28\,\mathrm{MeV}$ binding in $\Lambda(1405)$. Asymptotically, for $r \mapsto \infty$ one obtains $E = E^*$, *i.e.* a kaon bound to a proton to give the $\Lambda(1405)$. At short NN distances the binding energy reaches $\sim 300\,\mathrm{MeV}$.

With symmetrized nucleon states, this type of asymptotic behavior is characteristic for all solutions in few nucleon cases of interest. Eq. (15) also allows one to determine some broadening of $\Lambda(1405)$ in the NN system or more precisely the width of the KNN state. This effect becomes more transparent in the two channel $KN, \pi\Sigma$ formulation where ψ_i has two components and Eq. (12) becomes a 4×4 multichannel equation. Close to resonance the scattering amplitude is given by

$$f^{a,b} \approx \frac{\gamma_a \gamma_b}{E - E_0 + i\Gamma/2}\,. \tag{16}$$

Consistency requires the width to be $\Gamma/2 = p_\pi(\gamma_b)^2$ where p_π is the momentum in the decay channel. The singular term in Eq. (16) permits one to find an eigenvalue solution in a fairly simple form. It is presented below in the limit of zero range KN (and $\Sigma\pi$) force. The binding energy

$$\mathrm{Re}E = E_0 - (\gamma_K)^2 \frac{\cos(p_{\mathrm{R}}r)}{r}\exp(-p_{\mathrm{I}}r) - (\gamma_\pi)^2\,\frac{\cos(p_\pi r)}{r} \tag{17}$$

becomes larger than the binding of the resonance but the effect of the decay channel indicates oscillations. This oscillatory behavior is also seen in the width of the system

$$\mathrm{Im}E = -\Gamma/2\,\left[1 + \frac{\sin(p_\pi r)}{p_\pi r}\right] + (\gamma_K)^2\,\frac{\sin(p_{\mathrm{R}}r)}{r}\exp(-p_{\mathrm{I}}r)\,. \tag{18}$$

The effect of KN scattering represented by the second term enlarges the width as p_{R} is negative. In practice, this term describing the collision broadening is small. The contribution from multiple scattering in the decay channel oscillates and may under some conditions reduce the total width. That is an effect of destructive interference in the decay channel, known to exist in multichannel systems. In practice this effect is negligible in the KNN case but may be of significance in heavier systems. In particular it may reduce the widths of the mesonic state in K–$NNNN$ and η–$NNNN$ systems.

The parametrization of separable interactions is adjusted to reproduce the resonances in the KN center of mass system. For consistency, the meson propagator in Eq. (7) is chosen to make this equation equivalent to a differential equation

$$(E^{\mathrm{c}} 2\mu_{KN} + \Delta_x - U_{KN})\psi_K(x, x_i) = 0\,. \tag{19}$$

With nucleons fixed at some distance, the energy $E^{\rm c} = p(x_i)^2/2\mu_{KN}$ generates the potential contracting the NN system to a smaller radius. This effect is discussed now in terms of the Schrödinger equation.

The solution of the full KNN bound state problem is given by

$$\left(E - \frac{\Delta_x}{2m} - \frac{\Delta_1}{2M} - \frac{\Delta_2}{2M} - V_{KN} - V_{NN}\right)\Psi = 0\,, \tag{20}$$

where m, M are the masses of the meson and nucleon. To find the solution we assume the wave function to be given by Eq. (6) *i.e.* $\psi_K(x, x_1, x_2)\chi_{NN}(r)$.

Multiplying Eq. (20) on the left by ψ_K and integrating over the meson coordinate x one obtains the Schrödinger equation for the NN wave function

$$\chi_{NN}(r) = \int \boldsymbol{dx}\,\psi_K(x, x_i)\Psi(x, x_i) \tag{21}$$

in the form

$$\left(E - E^{\rm c}(r) + \frac{\Delta_1}{2M} + \frac{\Delta_2}{2M} - V_{NN}\right)\chi + \Delta E_{\rm kin}\chi = 0\,, \tag{22}$$

where the last term $\Delta E_{\rm kin}$ is a correction to the nucleon kinetic energies. It is given by

$$\Delta E_{\rm kin}\chi = -\frac{1}{M}\Sigma_i \int \boldsymbol{dx}\,\partial_i\psi_K\partial_i\Psi + \Delta_\delta\chi\,. \tag{23}$$

The first term turns out to be very small due to angular averaging and sign changes in both derivatives. The second comes from the δ-function in $\Delta\psi_K$ from Eq. (9) and is small but not negligible. These corrections may be calculated as perturbations. The generalization to few nucleons is in principle straightforward but the corresponding Schrödinger equation has to be solved by the variational method.

The actual binding energy is determined by the repulsive core in the NN interactions. With the Argonne [24] NN potential this method yields the lowest state of the KNN to be bound by about 40 MeV. Similar numbers are obtained with different calculations based on the same NN repulsion [25,26]. On the other hand, full three body calculations in Ref. [27] based on a limited separable description of the NN repulsion generate 50–70 MeV binding.

The corresponding $KNNN$ state is bound by about 90 MeV. In the KNN system the effect of the $\Sigma(1385)$ is small, but in the $KNNN$ state it adds some 15 MeV contribution to the binding. Another effect of $\Sigma(1385)$ is the existence of a new branch of P-wave states bound by 50–60 MeV. These states are very broad with $\Gamma \sim$ 40–50 MeV. However, the analogous

$KNNNN$ state seems to be of greater interest as now the binding is large enough to block the main two body decay channels and to reduce strongly the three body decay mode.

This work is supported by the KBN grant 1P0 3B 04229 and the EU Contract No. MRTN-CT-2006-035482, "FLAVIAnet". S.W. is grateful for the hospitality of Helsinki Institute of Physics.

REFERENCES

[1] E. Friedmann, Lecture presented at the XXX Mazurian Lakes Conference on Physics, Piaski, Poland, September 2–9, 2007.
[2] B. Chrien *et al., Phys. Rev. Lett.* **60**, 2595 (1988).
[3] B. Mayer *et al., Phys. Rev.* **C53**, 2068 (1996).
[4] H. Calen *et al., Phys. Rev. Lett.* **79**, 39 (1996).
[5] C. Wilkin, *Phys. Rev.* **C47**, R938 (1993); A. Sibirtsev *et al., Eur. Phys. J.* **A22**, 495 (2004).
[6] A.M. Green, S. Wycech, *Phys. Rev.* **C68**, 061601 (2003).
[7] A. Khoukaz *et al.*, II Eta Meson Net Workshop, `nucl-ex/0710.1809`
[8] E. Friedmann, A. Gal, J. Mares, A. Cieply, *Phys. Rev.* **60**, 024314 (1990).
[9] J. Mares, E. Friedmann, A. Gal, *Nucl. Phys.* **A777**, 84 (2006).
[10] A. Ramos, E. Oset, *Nucl. Phys.* **A671**, 481 (2000).
[11] S. Wycech, *Nucl. Phys.* **A450**, 399c (1986).
[12] P. Senger, *Acta. Phys. Pol. B* **39**, 321 (2007) these proceedings.
[13] R. Staronski *et al., J. Phys. G* **13**, 263 (2004).
[14] T. Suzuki *et al., Phys. Lett.* **B597**, 263 (2004).
[15] M. Iwasaki, *Proc. EXA 05, Austr. Acad. Sc. Press*, 2005, p. 93 `nucl-ex/0706.0297`.
[16] M. Agnello *et al.* [FINUDA Collaboration], *Phys. Rev. Lett.* **94**, 212303 (2005).
[17] V.K. Magas, E. Oset, A. Ramos, H. Toki, *Phys. Rev.* **C74**, 025206 (2006).
[18] Y. Akaishi, T. Yamazaki, *Phys. Rev.* **C65**, 044005 (2002).
[19] Y. Akaishi, T. Yamazaki, *Nucl. Phys.* **A792**, 229 (2007).
[20] K.A. Brueckner, *Phys. Rev.* **89**, 834 (1953).
[21] W.L. Foldy, J.D. Walecka, *Ann. Phys.* **54**, 447 (1969).
[22] A. Martin, *Phys. Lett.* **B179**, 33 (1981).
[23] S. Wycech, A.M. Green, *Int. J. Mod. Phys.* **A22**, 629 (2007).
[24] R.B. Wiringa *et al., Phys. Rev.* **C51**, 38 (1995).
[25] T. Yamazaki, Y. Akaishi, *Phys. Rev.* **C76**, 045201 (2007).
[26] A. Dote, W. Weise, *Eur. Phys. J.* `nucl-th/0701050`
[27] N. Shevchenko, A. Gal, J. Mares, *Phys. Rev.* **C76**, 044004 (2007).