Chiral doublet model for positive and negative parity nucleons

D. Jido

Research Center for Nuclear Physics (RCNP), Osaka University
Ibaraki Osaka 567-0047, Japan

We discuss the properties of the nucleon ($N$) and its excited state with odd parity ($N^*$) under the aspect of chiral symmetry, identifying them with suitable representations of the chiral group. It is shown that there are two distinctive schemes to assign the chiral multiplet to $N$ and $N^*$. We construct linear $\sigma$ models for $N$ and $N^*$ with the two assignments to show their physical implications. We also discuss the in-medium properties of $N(1535)$, comparing the present model with the other chiral model based on the different picture of $N^*$.

§1. Introduction

Chiral symmetry (ChS) is one of the key concepts for understanding the structure and dynamics of hadrons at low energies from the viewpoint of QCD. The importance of the dynamical breakdown of ChS is summarized in the existence of the light pion as a Nambu-Goldstone boson. On the other hand, we believe restoration of the broken symmetry at hot and/or dense matter, and one of the interesting implications there is the appearance of the degenerate spectra in parity partners. It has been recently pointed out that partial restoration of ChS may take place in a nucleus and it will be observed as effective modifications of hadron properties in the nuclear medium.

The $N(1535)$ ($N^*$) is an especially interesting nucleon resonance. It is the first excited state with odd parity and is considered to be a possible candidate of a chiral partner of the nucleon. In addition, it is well known that creation of $N^*$ in intermediate states is identified by emission of the $\eta$ meson in the final state as a result of the strong $\eta NN^*$ coupling. The recent theoretical studies of $N^*$ respecting ChS have been done in two distinct pictures of $N^*$: (1) a chiral partner of the nucleon, (2) a dynamically generated object in meson-baryon scattering.

First of all, it is worth emphasizing that ChS is unambiguously defined in the QCD Lagrangian for the quark field as separated rotations in the flavor space:

$$ q_l \rightarrow L q_l \quad q_r \rightarrow R q_r \quad (1.1) $$

where $q_l$ (or $q_r$) is the left- (or right-) handed component of the quark field in the sense of Lorentz group, and $L$ and $R$ are independent $SU(N_f)$ rotations for the $N_f$ flavors.

On the other hand, the realization of ChS in hadrons is not a trivial issue as the reflections of their quark composite structure and the spontaneous breaking of ChS. There are two possible ways of the realization of ChS for hadrons, non-linear and linear realizations. They are not conflicting concepts but compliment each other.

In the nonlinear realization, the effective Lagrangians are constructed on the...
premise that ChS is spontaneously broken. The manifestation of ChS in Lagrangian is accomplished by giving the special role in the axial transformation to the Nambu-Goldstone boson. The transformation rule of the other hadron under ChS is uniquely fixed, once the transformation rule under the vector rotation is given.\(^8\). Dynamical properties are determined according to an expansion in powers of momenta of Nambu-Goldstone bosons. This leads us to obtain the most general Lagrangian at low energies, which is summarized in chiral perturbation theory.\(^9\),\(^10\).

On the other hand, in the linear representation, based on the fact that all hadrons are in principle classified into some representations of the chiral group \(SU(N_f)_L \otimes SU(N_f)_R\), the effective Lagrangians are constructed with assigning irreducible representations to the hadrons. The linear realization has a connection between chiral symmetric and broken phase, and gives constraints on intrinsic properties of hadron appearing in the broken phase, such as the mass of the nucleon.

In this paper, we would like to discuss the nucleon and its excited state with odd parity in a unified formulation based on the linear realization, considering the fact that the parity degeneracy appears in the restoration limit of ChS, where the non-linear realization breaks down. We discuss the lowest-lying \(N(1535)\) as the excited state with odd parity considered here, although it can be, in principle, assumed to be any nucleon excited state with \(J^P = (1/2)^-\).

In Sec. 2, we discuss how to realize chiral symmetry for \(N\) and \(N^*\) in the linear representation and see that there are two possible ways to assign the chiral transformation to \(N^*\). According to these assignments, in Sec. 3, we construct the two chiral doublet models and discuss their physical consequences. In Sec. 4, we apply the chiral double model to the study of the in-medium properties of \(N(1535)\) probed by \(\eta\) mesic nuclei. Summary is presented in Sec. 5.

\section{Chiral symmetry for baryons}

In order to construct an effective Lagrangian in the linear representation, it is necessary to assign an appropriate irreducible representation of the chiral group to the nucleon \(N\) and its odd parity excited state \(N^*\). The chiral multiplet is a good quantum number in the chiral symmetric limit and represents the quark configuration inside the nucleons. The most suitable combinations of the chiral multiplets for the nucleon should be in principle determined by the dynamics of the quarks and gluons. Now let us concentrate the flavor two case \((N_f = 2)\) and the chiral limit \(m_q = 0\), and we do not consider the possible mixing to the other representations.

As mentioned in introduction, ChS in QCD is defined in terms of the quark field, and the \(u\)- and \(d\)-quark fields belong to the fundamental representation \((\frac{1}{2}, 0) \oplus (0, \frac{1}{2})\) where the first and second numbers in the parenthesis expresses the irreducible representations of \(SU(2)_L\) and \(SU(2)_R\), respectively. Considering that the baryons consist of three valence quarks, possible candidates of the chiral multiplet for the baryons may be given by the following three multiplets:

\[
\left[\left(\frac{1}{2}, 0\right) \oplus \left(0, \frac{1}{2}\right)\right]^3 = 5 \left[\left(\frac{1}{2}, 0\right) \oplus \left(0, \frac{1}{2}\right)\right] \oplus 3 \left[\left(\frac{1}{2}, 1\right) \oplus \left(1, \frac{1}{2}\right)\right] \oplus \left[\left(\frac{3}{2}, 0\right) \oplus \left(0, \frac{3}{2}\right)\right]
\]

\section{(2.1)
Chiral doublet model for the positive and negative parity nucleons

The terms in the first and third parentheses in r.h.s. have purely isospin 1/2 and 3/2, respectively, while the terms in the second parenthesis is the mixture representation of the isospin 1/2 and 3/2. Here we take for $N$ the first representation $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$, which has the isospin 1/2. This representation has been used in the linear $\sigma$ model by Gell-Mann and Lévy, the QCD sum rules$^{12)}$ and lattice QCD calculations. We assume that the equivalent multiplet to the nucleon is assigned to $N^*$.

After the choice of the irreducible representation for both nucleons, there are two possible way to assign the chiral multiplet to $N$ and $N^*$, depending on how to introduce the mass terms of the nucleons consistently with ChS.

In the first case, the transformation rules for $N$ and $N^*$ are given by

\[
N_l \to LN_l \quad N_r \to RN_r
\]

\[
N^*_l \to LN^*_l \quad N^*_r \to RN^*_r
\]

(2.2)

which is called as the naive model$^4)$. The matrices $L$ and $R$ in (2.2) represent the rotations of $SU(2)_L$ and $SU(2)_R$, respectively, and, $N_{l,r}$ and $N^*_{l,r}$ are their doublets. In this case, the explicit introduction of the nucleon mass terms breaks invariance under the axial transformation of ChS. The only prescription to make models invariant under ChS is that the nucleons is introduced as massless Dirac particles coupling to a scalar field and condensate of the scalar field induces generation of the nucleon masses as well as the spontaneous symmetry breaking. However, in this case, the two nucleons belong to the completely separate multiplet and transform independently under ChS. Therefore there are no connections between them in the group theoretical point of view. If $N^*$ were assigned to even parity, there would be no changes in the argument presented here for the naive model because of no connection between $N$ and $N^*$ in terms of ChS.

Alternatively it is possible to keep the invariance of the mass term under the linear transformation when we introduce the chiral partner of the nucleon, which is the particle to form the parity degeneracy in the restoration limit of ChS. Let us consider the nucleons $N$ and $N^*$ which are the chiral partners and transform each other under the axial transformation of ChS, similarly to the $\sigma$ and $\pi$ fields. The nucleon mass term is written with a common mass $m_0$ in a chiral invariant way as

\[
m_0(\bar{N}N + \bar{N}^*N^*) ,
\]

(2.3)

The physical basis $(N, N^*)$ is different form the basis of ChS, which is defined in the transformation rule under $SU(2)_L \otimes SU(2)_R$ as

\[
N_{1r} \to RN_{1r} \quad N_{1l} \to LN_{1l}
\]

\[
N_{2r} \to LN_{2r} \quad N_{2l} \to RN_{2l}
\]

(2.4)

where $N_1$ and $N_2$ are assumed to have even and odd parity, respectively. Note that $N_1$ and $N_2$ transform in the opposite way under the axial transformation. This is called as the mirror model$^4)$. It is possible to introduce the following mass term without any contradictions with ChS$^2$:

\[
m_0 (\bar{N_1}\gamma_5N_2 - \bar{N_2}\gamma_5N_1).
\]

(2.5)
The physical basis is obtained so as to diagonalize the mass term (2.5):

\[ N = \frac{1}{\sqrt{2}} (N_1 + \gamma_5 N_2) \quad N^* = \frac{1}{\sqrt{2}} (\gamma_5 N_1 - N_2). \]  

(2.6)

It is shown that \( N \) transforms to \( N^* \) under the axial transformation (2.4).

Here let us make a remark regarding to the axial \( U(1) \) symmetry. The QCD Lagrangian is invariant under the global \( U(1) \) axial transformation, and the \( U(1)_A \) symmetry is anomalously broken due to the quantum effect. In the present work, although we are constructing the effective models for hadrons which have the same symmetries as QCD, we do not assume the \( U(1)_A \) symmetry in the effective models, since the effective models emerge after the quark loops are integrated out. Nevertheless, the \( U(1)_A \) charge of hadron gives further constraints on the quark structure of the hadron. For instance, the chiral multiplet \( (\frac{1}{2}, 0) \oplus (0, \frac{1}{2}) \) considered here is composed in the two different ways \( (\bar{q}_l q_i \gamma_5 \vec{\tau} \cdot \vec{\pi})_{I=0} \oplus (\bar{q}_r q_i)_{I=0} \) and \( (\bar{q}_l q_i \gamma_5 \vec{tau})_{I=0} \oplus (\bar{q}_l q_i)_{I=0} \). Both are the same multiplet in the \( SU(2)_L \otimes SU(2)_R \) group, but have the different charges of \( U(1)_A \).

§3. Chiral double models

In this section we briefly discuss the phenomenological consequences of the two assignments introduced in the preceding section, constructing the linear \( \sigma \) models according to their transformation rules. The detailed discussion has been shown in Refs. 4, 5. The important consequences are summarized in Table 1.

Considering the chiral transformation rule for the scalar and pseudo-scalar fields \( M \equiv \sigma + i\vec{\tau} \cdot \vec{\pi} \rightarrow LMR^1 \), we obtain the linear \( \sigma \) model with the naive assignment as

\[ \mathcal{L}_{\text{naive}} = \sum_{j=1,2} [\bar{N}_j i\gamma_5 \sigma N_j - a_j \bar{N}_j (\sigma + i\gamma_5 \vec{\tau} \cdot \vec{\pi}) N_j] + \mathcal{L}_{\text{meson}}, \]  

(3.1)

where \( N_1 \equiv N \) and \( N_2 \equiv N^* \), and \( a_1 \) and \( a_2 \) are free parameters independent of ChS. The Lagrangian \( \mathcal{L}_{\text{meson}} \) is for the \( \sigma \) and \( \pi \) fields, and its explicit form is irrelevant for the present argument as long as it causes the spontaneous chiral symmetry breaking. The transition coupling of \( N \) and \( N^* \) with the meson field is also possible to be introduced in the chiral invariant way \( 4, 5, \) but the coupling term is always diagonalized by a suitable linear combination of \( N \) and \( N^* \) without any contradiction with the chiral transformation rule (2.2). This Lagrangian has the minimal terms invariant under ChS and it is allowed to add more terms with derivatives and multiple \( M \) fields. These terms give corrections to the axial charges and the masses of the nucleons in powers of the sigma condensation \( \sigma_0 \).

The Lagrangian of the naive model is just a sum of two independent linear \( \sigma \) models. Therefore the phenomenological consequences are followed by those of the usual linear \( \sigma \) model for the single nucleon. The masses of \( N \) and \( N^* \) are calculated with the finite condensate of the scalar field as

\[ m_N = a_1 \sigma_0, \quad m_{N^*} = a_2 \sigma_0. \]  

(3.2)

The isovector axial charges of \( N \) and \( N^* \) are unities independently of \( \sigma_0 \) at tree level. There are no transitions between \( N \) and \( N^* \) with pion, which is qualitatively...
consistent with the empirically small value of the $\pi NN^*$ coupling $g_{\pi NN^*} \simeq 0.7$ compared to the strong $\pi NN$ coupling $g_{\pi NN} \simeq 13$.

Another interesting consequence is that the values of the masses and the axial charges in the chiral restoration limit; $m_N = m_{N^*} = 0$, $g_A^N = g_A^{N^*} = 1$ and $g_A^{NN^*} = 0$. This is a quite general conclusion. Even if we add more terms invariant under ChS to Lagrangian (3.1), their contributions to the masses and the axial charges are written in powers of $\sigma_0$, and, therefore, the masses of $N$ and $N^*$ are decreasing to zero and their axial charges are approaching to unity and zero at least close to the restoration limit [14].

Now let us turn to the mirror model. The Lagrangian is given by

$$\mathcal{L}_{\text{mirror}} = \sum_{j=1,2} \left[ N_j i\bar{\phi}N_j - g_j \bar{N}_j (\sigma + (-)^{j+1}i\gamma_5 \vec{\tau} \cdot \vec{\pi}) N_j \right] - m_0 (\bar{N}_1 \gamma_5 N_2 - \bar{N}_2 \gamma_5 N_1) + \mathcal{L}_{\text{meson}}, \quad (3.3)$$

where $g_1$, $g_2$ and $m_0$ are free parameters, and we assume to truncate the terms with derivatives and multi mesons again. The Yukawa term mixing $N_1$ and $N_2$ is not invariant under the transformation (2.4). This Lagrangian was first formulated and investigated by DeTar and Kunihiro [2]. Historically a similar Lagrangian to the present one was considered before by B. Lee in [15], but symmetry between $N$ and $N^*$ under the axial $U(1)$ transformation was implicitly assumed in his Lagrangian and he got physically uninteresting results.

As mentioned in the previous section, the physical nucleons $N$ and $N^*$ diagonalizing the mass term are given by a linear combination of $N_1$ and $N_2$. After breaking ChS spontaneously, $N$ and $N^*$ are given with a mixing angle $\theta$ as

$$N = \cos \theta N_1 + \sin \theta \gamma_5 N_2, \quad N^* = -\sin \theta \gamma_5 N_1 + \cos \theta N_2. \quad (3.4)$$

The mixing angle depends on the sigma condensate: $\tan 2\theta = 2m_0/\sigma_0(g_1 + g_2)$. The corresponding masses are calculated as

$$m_{N,N^*} = \frac{1}{2} \left( \sqrt{(g_1 + g_2)^2\sigma_0^2 + 4m_0^2} + (g_2 - g_1)\sigma_0 \right). \quad (3.5)$$

In this model, mass degeneracy of $N$ and $N^*$ takes place with a finite mass $m_0$ in the chiral restoration limit, and the spontaneous breaking of ChS causes the mass splitting. The model parameters are fitted so as to reproduce the $N$ and $N^*$ masses and the $\pi NN^*$ coupling: $g_1 = 9.8$, $g_2 = 16$ and $m_0 = 270$ MeV [2,4,5]. The mixing angle is calculated as $\theta = 6.3^\circ$. The similar mass formula to (3.5) has been obtained for the isospin 1/2 and 3/2 nucleon resonances, such as $\Delta$, $N(1520)$ and their parity partners, in the same approach presented here with the multiplet $((1/2, 1) \oplus (1, 1/2))$, and it is consistent with the observed pattern of the mass spectra [16].

One of the phenomenological significances of this model is that the axial charges of $N$ and $N^*$ have the opposite sign to each other. The isovector axial charges are calculated in a function of the mixing angle as

$$g_A = \begin{pmatrix} \cos 2\theta & -\sin 2\theta \\ -\sin 2\theta & \cos 2\theta \end{pmatrix} \quad (3.6)$$
This shows that the relative sign of the axial charges of $N$ and $N^*$ is negative independently of the mixing angle and that the off-diagonal component does not necessarily vanish. It is also shown that in the chiral restoration limit ($\theta = \pi/4$) the diagonal components are zero, while the off-diagonal component is unity. Considering the empirical value of the transition axial charge $g_{NN^*}^A \approx 0.2$ obtained by the generalized Goldberger-Treiman relation with $g_{\pi NN^*} \approx 0.7$, the transition axial charge $g_{NN^*}^A$ is enhanced to unity as the sigma condensate decreases\(^{14}\).

The parameter $m_0$ introduced here is a new parameter not constrained by ChS, and it gives the nucleon mass in the chiral restoration limit. If the mirror model is realized in the physical nucleon, the origin of $m_0$ in QCD is important to understand the mirror nucleons in the context of non perturbative QCD\(^{5}\).

| Table I. Summary of the phenomenological consequences of the naive and mirror models. |
|-----------------------------------------------|
| Naive model | Mirror model |
| Definition | $N_{1r} \rightarrow RN_{1r}, \ N_{1l} \rightarrow LN_{1l}$ | $N_{1r} \rightarrow RN_{1r}, \ N_{1l} \rightarrow LN_{1l}$ |
| Nucleon mass term | $N_{2r} \rightarrow RN_{2r}, \ N_{2l} \rightarrow LN_{2l}$ | $N_{2r} \rightarrow LN_{2r}, \ N_{2l} \rightarrow RN_{2l}$ |
| Nucleon in Wigner phase generated by scalar field | massless | mass generation |
| Role of $\sigma_0$ | introduced with chiral partner | mass splitting |
| Chiral partner | $N \leftrightarrow \gamma_5 N$ | $N \leftrightarrow N^*$ |
| Sign of $\pi N^* N^*$ coupling | positive | negative |
| In-medium $\pi NN^*$ coupling | suppressed | enhanced |

§4. Application: in-medium properties of $N(1535)$ probed by $\eta$ mesic nuclei

In this section, we show an application of the parity doublet model to the investigation of the in-medium properties of $N(1535)$. It has been pointed out in Ref. 17) that formation experiments of $\eta$ mesic nuclei, such as $(d,^3\text{He})$ reactions with nuclear targets, could be good tools to observe the in-medium effect on the mass of $N^*$.

The basic idea is that the $\eta$ optical potential in the nucleus is expected to be very sensitive to medium modifications of the $N^*$ mass. Assuming the $N^*$ dominance hypothesis and the heavy baryon limit, we obtain the $\eta$ optical potential as

$$V_\eta = \frac{g_\eta^2}{2\mu} \omega + m_N^*(\rho) - m_{N^*}^*(\rho) + i\Gamma_{N^*}(s;\rho)/2$$

(4.1)

with $g_\eta$ the $\eta NN^*$ coupling, $\mu$ the reduced mass of the $\eta$-nucleus system and $\rho(r)$ the nuclear density. Considering the fact that the $N^*$ mass in free space lies only 50 MeV above the $\eta N$ threshold, we conclude that the $\eta$-nucleus potential turns to be repulsive at the nuclear center, if the in-medium effect leads to a significant reduction of the mass difference of $N$ and $N^*$\(^{17}-19\). We expect to observe the repulsive nature in the formation experiment of $\eta$ mesic nuclei with the $(d,^3\text{He})$ reactions.

We use the parity doublet model to calculate the $N$ and $N^*$ masses and the $N^*$ decay width in medium. We assume the partial restoration of ChS in nuclei with a linear parametrization of the density dependence of the sigma condensate\(^{20},21\):

$$\langle \sigma \rangle = (1 - C\rho/\rho_0)\sigma_0$$

(4.2)
Chiral doublet model for the positive and negative parity nucleons

with $C = 0.1 - 0.3$ and $\rho_0$ the saturation density. Since the mass difference of $N$ and $N^*$ is proportional to the sigma condensate in both chiral doublet models as seen in eqs. (3-2) and (3-5), the in-medium mass difference is obtained as

$$m_{N^*}(\rho) - m_N(\rho) = (1 - C \rho/\rho_0)(m_{N^*} - m_N)$$  \hspace{1cm} (4.3)

with the density dependent sigma condensate (4.2). Therefore, as long as partial restoration of ChS is assumed in the nuclear medium, the mass difference is reduced as the density increases and the $\eta$ optical potential in nucleus has possibility to turn to be repulsive at the center.

We show in Fig. 1 the spectra of the $(d,^3\text{He})$ reaction with $^{12}\text{C}$ target calculated in the various cases of the $\eta$-nucleus interaction. In Fig. 1 (a), the medium modifications of the $N$ and $N^*$ masses are not assumed and the $\eta$ optical potential is attractive in the nucleus independently of the density. This case is equivalent to the so-called $T-\rho$ approximation. Shown in Fig. 1 (b) is the spectrum calculated in the mirror model with the partial restoration of ChS and its strength parameter $C = 0.2$. Here we find the significant difference in these two plots. In the mirror model, as a result of the repulsive nature of the $\eta$ optical potential at the center of nucleus, the spectrum is shifted to the higher energies. Note that the peak structure shown in the plots is not responsible for the formation of the $\eta$ bound state but just the threshold effects.

It is also interesting to compare the above results with the spectrum calculated by the chiral model based on the different picture of $N^*$, such as a dynamically generated object in meson-baryon scattering. This model was formulated first in Ref. 6. There $N^*$ is calculated in the coupled channels of $\pi N$, $\eta N$, $K \Lambda$ and $K \Sigma$, and the $N^*$ is found to be formed dominantly as a $K \Sigma$ state. Since the in-medium modification of $N^*$ is insignificant in this model, the optical potential of $\eta$ in nuclei is basically attractive inside of the nuclei. Here we directly take the in-medium $\eta$ optical potential shown in Ref. 23 to calculate the $(d,^3\text{He})$ spectrum. As shown in Fig. 1 (c), as a result of the irrelevance of the in-medium modification of $N^*$, the spectrum has the similar shape to Fig. 1 (a). We would expect that the spectra obtained with the different pictures of $N^*$ are distinguished in experiment.

§5. Summary

We have investigated the properties of the nucleon and its excited state with odd parity in the effective models which are strongly constrained by chiral symmetry. There are two possible ways to assign the chiral transformation to $N^*$ in the linear realization of chiral symmetry. So far we do not know which models is realized in the physical nucleon and excited state. To confirm it experimentally, the important observable is the sign of the isovector axial charge of $N^*$.

We have also discussed the in-medium properties of $N(1535)$ probed by $\eta$ mesic nuclei based on the two distinct physical pictures of the structure of $N^*$. We have found that the models based on these pictures produce quantitatively different consequences and they would be distinguishable in formation experiments of $\eta$-mesic nuclei, for instance, the $(d,^3\text{He})$ reaction with nuclear target.
Fig. 1. The calculated spectra of $^{12}$C($d,^3$He)$^{11}$B$\otimes\eta$ reaction at $T_d=3.5$ GeV are shown as functions of the excited energy $E_{ex}$. $E_0$ is the $\eta$ production threshold energy. The $\eta$-nucleus interaction is calculated by (a) the $t\rho$ approximation, (b) the chiral doublet model with $C=0.2$ and (c) the chiral unitary approach. The total spectra are shown by the thick solid lines, and the dominant contributions from the $(0s_{1/2})^{-1}p\otimes s_\eta$ and the $(0p_{3/2})^{-1}p\otimes p_\eta$ configurations are shown by dashed lines and dash-dotted lines, respectively. Here the proton-hole states are indicated as $(n\ell_j)p^{-1}$ and the $\eta$ state as $\ell_\eta$.

Acknowledgements

I would like to express my sincere gratitude to Prof. M. Oka, Prof. A. Hosaka, Dr. Y. Nemoto, Dr. H. Kim, Prof. T. Hatsuda, Prof. T. Kunihiro, Prof. S. Hirenzaki and Dr. H. Nagahiro, who have participated in the works reported here.

References

1) See, for example, T. Kunihiro, the contribution to this conference proceedings.
2) C. DeTar and T. Kunihiro, Phys. Rev. D 39 (1989), 2805.
3) D. Jido, M. Oka and A. Hosaka, Phys. Rev. Lett. 80 (1998), 448.
4) D. Jido, Y. Nemoto, M. Oka and A. Hosaka, Nucl. Phys. A 671 (2000), 471.
5) D. Jido, M. Oka and A. Hosaka, Prog. Theor. Phys. 106 (2001), 873.
6) N. Kaiser, P.B. Siegel and W. Weise, Phys. Lett. B 362 (1995), 23.
7) T. Inoue, E. Oset and M.J. Vicente Vacas, Phys. Rev. C 65 (2002), 035204.
8) S. Weinberg, Phys. Rev. 166 (1968), 1568.
9) S. Weinberg, Physica A 96 (1979), 327.
10) J. Gasser and H. Leutwyler, Ann. Phys. 158 (1984), 142.
11) T.D. Cohen and X.D. Ji, Phys. Rev. D 55 (1997), 6870.
12) B.L. Ioffe, Nucl. Phys. B 188 (1981), 317.
13) G.A. Christos, Z. Phys. C 21 (1983), 83; ibid. 29 (1985), 361; Phys. Rev. D 35 (1987), 330.
14) H.e. Kim, D. Jido and M. Oka, Nucl. Phys. A 640 (1998), 77.
15) B.W. Lee, Chiral Dynamics (Gordon and Breach, 1972).
16) D. Jido, T. Hatsuda and T. Kunihiro, Phys. Rev. Lett. 84 (2000), 3252.
17) D. Jido, H. Nagahiro and S. Hirenzaki, Phys. Rev. C 66 (2002), 045202.
18) D. Jido, H. Nagahiro and S. Hirenzaki, Nucl. Phys. A 721 (2003), 665c.
19) H. Nagahiro, D. Jido and S. Hirenzaki, nucl-th/0304068, to be published in Phys. Rev. C.
20) T. Hatsuda, T. Kunihiro and H. Shimizu, Phys. Rev. Lett. 82 (1999), 2840.
21) D. Jido, T. Hatsuda and T. Kunihiro, Phys. Rev. D 63 (2001), 011901.
22) T. Waas and W. Weise, Nucl. Phys. A 625 (1997), 287.
23) T. Inoue and E. Oset, Nucl. Phys. A 710 (2002), 354.
24) D. Jido, M. Oka and A. Hosaka, Prog. Theor. Phys. 106 (2001), 823.