Antisymmetrized Molecular Dynamics with Coherent State Pion and Its Application to Excited Spectrum of $^{12}$C

Akinori ISSHIKI, Kenichi NAITO* and Akira OHNISHI

Division of Physics, Graduate School of Science, Hokkaido University, Sapporo 060-0810, Japan

* Meme Media Laboratory, Hokkaido University, Sapporo 060-8628, Japan

(Received March 30, 2022)

We have introduced coherent state neutral pion into Antisymmetrized Molecular Dynamics. With the aid of coherent state technique, it becomes possible to calculate transition matrix elements of the pion field operator and to study excited states containing pions. For large pion-nucleon coupling $f_{\pi N} \gtrsim 1.6$, pions have a finite expectation value and bring large energy gain in $^{12}$C. We discuss two aspects of pionic effects in spectroscopy; the $LS$ interaction like effect and the mixing of different nucleon parity states, which would modify low energy nuclear levels.

§1. Introduction

Nuclei have been basically understood as nucleon many-body systems in which nucleons move in a mean field and interact via small residual interactions. In shell models, each nucleon single particle wave function is assumed to have its own orbital angular momentum, spin and parity ($lj\pi$), provided that the spherically symmetric mean field consists of central and spin-orbit ($LS$) parts. This basic picture of nuclei has been successful in describing low-lying states of most of medium to heavy nuclei, while there are several exceptions such as clustering states in light nuclei.

On the other hand, the long range part ($r \gtrsim 2$ fm) of the bare nucleon-nucleon potential is described by the one-pion exchange potential (OPEP) having strong tensor part, which mixes different partial waves. This mixing is essential for the binding of deuteron, but makes it difficult to treat exactly in many-body systems. Thus tensor force has been usually treated in the form of effective central and $LS$ forces in solving nuclear many-body problems in spite of its importance in bare nuclear force, hoping that the explicit role of the tensor force is not large in nuclei. Actually, the first order effects of tensor force vanish in the Fermi gas state, due to the cancellation in $(2J + 1)$-weighted sum.

If pions have expectation values in nuclei or in nuclear matter, the cancellation does not work and tensor force may play a dominant role. There have been a lot of discussions on the possibilities of pion condensation in nuclear matter. For some time, it was considered that strong nucleon-$\Delta$ short range repulsion might suppress pion condensation to emerge, provided that the Landau-Migdal parameter follows the "universality", $g'_{N\Delta} = g'_{NN} \simeq 0.6$, and that it is density-independent. Recent observations of the non-quenching in Gamow-Teller giant resonance sum rule clarified that the short range repulsion between nucleons and $\Delta$ resonances is
not very strong, $g'_{N\Delta} < 0.25$, suggesting that the pion condensation would come in exist at least in high-density nuclear matter. Also in recent ab-initio calculations of light nuclei ($A \leq 12$) with realistic bare $NN$ potentials, it has been shown that the OPEP contribution is dominant in the total potential energy. This result suggests that the cancellation in the Fermi gas is not working well in actual nuclei and that it would be necessary to consider the explicit role of pions more seriously.

On these backgrounds, pion condensation and explicit role of tensor force in nuclei has now attracting renewal interests. In a relativistic mean field framework, it is demonstrated that neutral pions can condensate in the surface region of nuclei, and that this condensation enhances the binding energy of $jj$ closed nuclei, such as $^{12}$C. In their model, single particle states are first prepared to have fixed $lj\pi$, and those states having the same $j$ are mixed to gain the potential energy from pions. This mixture also plays a role of $LS$ like potential. Since pions mix different parity (but same $j$) states, the yrast single particle states (the lowest energy single particle states for a given $j$, $s_{1/2}, p_{3/2}, d_{5/2}, f_{7/2}, g_{9/2}, h_{11/2}, i_{13/2}$) will have the largest energy gain. On the other hand, those states having smaller $j$ at around the Fermi energy will be pushed up from the mixing with the lower energy single particle states. It is interesting to note that the last four yrast single particle states are in charge of the nuclear magic numbers of 28, 50, 82 and 126. More recently, Charge and Parity Projected Hartree-Fock (CPPHF) method has been developed in order to take account of the coupling of proton and neutron single particle states generated by OPEP. It has been shown that the charge projection enhances the tensor contribution by around three times in the case of $^4$He nuclei.

At this stage, it would be desirable to extend the scope of pion and tensor force study from the ground state and single particle states to nuclear excited level spectroscopy with specified $J\pi$, which has richer information on wave functions. Specifically, we are more interested in constructing a framework in which explicit pionic degrees of freedom are incorporated, rather than introducing tensor interaction, since we believe that it is more fundamental to describe nuclear many-body system with pions.

In this work, we introduce coherent state pions into Antisymmetrized Molecular Dynamics (AMD) and discuss the pionic effects on excited states of $^{12}$C. Pion coherent state enables us to calculate matrix elements of the pion operator with different states as well as the expectation value with a given state. In AMD, nuclear wave function is represented by the Slater determinant of nucleon Gaussian wave packets, which is wide enough to describe clustering states as well as shell model states. By using the product of nucleon AMD state and pion coherent state, we can evaluate the transition matrix element of the Hamiltonian operator containing nucleon and pion operators. Therefore, it becomes possible to project the wave function to the eigen state of given $J\pi$ and to diagonalize the Hamiltonian matrix consisting of wave functions with different nucleon and pion configurations.
§2. AMD with Coherent State Pion

The nucleon-pion basis state is assumed to be the product of the nucleon AMD state\textsuperscript{10,11} and the pionic coherent state\textsuperscript{9}

\[ |\Psi(Z, f)\rangle = |\Psi_{\text{AMD}}(Z)\rangle \otimes |\Phi_{\pi}(f)\rangle. \]

AMD wave function is a Slater determinant of nucleon Gaussian wave packets,

\[ |\Psi_{\text{AMD}}(Z)\rangle = A \prod_i |\psi_{z_i}\rangle |\chi_i^\sigma \chi_i^\tau\rangle \quad (Z = \{z_i \mid i = 1, 2, \ldots A\}), \]

\[ \langle r|\psi_z\rangle = \left(\frac{2\nu}{\pi}\right)^{3/4} \exp\left[-\nu(r - z/\sqrt{\nu})^2 + z^2/2\right], \]

where \( |\chi_i^\sigma \chi_i^\tau\rangle \) represents spin-isospin wave function.

Pion coherent state introduced by Amado et al.\textsuperscript{9} is represented as,

\[ |\Phi_{\pi}(f)\rangle = \exp\left[\int dk f(k)\hat{\alpha}(k)\right]|0\rangle. \]

By setting the commutation relation of the annihilation and creation operator, \( \hat{a}(k) \) and \( \hat{a}^\dagger(k') \), as \( [\hat{a}(k), \hat{a}^\dagger(k')] = \delta(k - k') \), we can show that the above pion coherent state is an eigen state of the positive frequency operator \( \hat{\phi}^{(+)} \),

\[ \hat{\phi}^{(+)}(r, t) = \int \frac{hc}{\sqrt{(2\pi)^3 2\omega_k / c}} \hat{a}(k)e^{ik \cdot r - i\omega_k t}, \]

\[ \hat{\phi}^{(+)}(r, t)|\Phi_{\pi}(f)\rangle = \varphi_f(r, t)|\Phi_{\pi}(f)\rangle, \]

\[ \varphi_f(r, t) = \int \frac{hc}{\sqrt{(2\pi)^3 2\omega_k / c}} f(k)e^{ik \cdot r - i\omega_k t}. \]

Since the bra state is an eigen state of the negative frequency operator, \( \hat{\phi}^{(-)}(r, t) \), and the pion operator is a sum of \( \hat{\phi}^{(+)} \) and \( \hat{\phi}^{(-)} \), we can easily calculate the transition matrix element of the pion operator as follows,

\[ \hat{\phi}(r, t) = \hat{\phi}^{(+)}(r, t) + \hat{\phi}^{(-)}(r, t), \quad \hat{\phi}^{(-)}(r, t) = \left(\hat{\phi}^{(+)}(r, t)\right)^\dagger, \]

\[ \langle \Phi_{\pi}(f)|\hat{\phi}(r, t)|\Phi_{\pi}(g)\rangle = N_{\pi}(f, g) \times (\varphi_f(r, t) + \varphi_g(r, t)), \]

\[ N_{\pi}(f, g) \equiv \langle \Phi_{\pi}(f)|\Phi_{\pi}(g)\rangle = \exp\left[\int dk f(k)g(k)\right]. \]

Now we consider the following Hamiltonian of \( N \)-body nucleon and pion system containing the second quantized pion operator \( \hat{\phi} \) in the axial vector \( P \)-wave pion-nucleon coupling,

\[ H = \sum_{i=1}^N \frac{p_i^2}{2m} + \sum_{i<j} V(r_{ij}) + \frac{1}{2hc} \int dr \left[ \nabla \hat{\phi}(r) \cdot \nabla \hat{\phi}(r) + \mu_{\pi}^2 \hat{\phi}^2(r) \right]. \]
\[ + \sum_{i=1}^{N} \frac{f_{\pi N}}{\mu_{\pi}} \tau_0 (\sigma_i \cdot \nabla_i) \hat{\phi}(r_i) , \]  

(2.11)

where \( \mu_{\pi} = m_{\pi} c / \hbar \), and we have omitted the time dependence in the pion part.

The matrix element of this Hamiltonian is evaluated as

\[ \mathcal{H} = \langle \Psi(Z, f) | \hat{H} | \Psi'(Z', g) \rangle = \mathcal{H}_N(Z, Z') + \mathcal{H}_\pi(f, g) + \mathcal{H}_\pi N(Z, Z', f, g) , \]  

(2.12)

\[ \mathcal{H}_\pi(f, g) = \int \frac{dr}{2\hbar c} \left[ \{ \nabla (\varphi_f(r) + \varphi_g(r)) \}^2 + \mu_{\pi}^2 \{ \varphi_f(r) + \varphi_g(r) \}^2 \right] , \]  

(2.13)

\[ \mathcal{H}_\pi N(Z, Z', f, g) = \frac{f_{\pi N}}{\mu_{\pi}} \int dr S(r) \cdot \nabla (\varphi_f(r) + \varphi_g(r)) , \]  

(2.14)

\[ S(r) = \langle \Psi_{\text{AMD}}(Z) | \sum_i \sigma_i \tau_0 \delta(r - r_i) | \Psi_{\text{AMD}}(Z') \rangle / \langle \Psi_{\text{AMD}}(Z) | \Psi_{\text{AMD}}(Z') \rangle . \]  

(2.15)

Here \( \mathcal{H}_N \) is the usual AMD Hamiltonian matrix element including \( NN \) interaction.

In the actual calculation, we have expanded the pion eigen function \( \varphi(r) \) in local Gaussians, whose centers and amplitudes are the variation parameters. Thus we can apply the cooling equations for these pion parameters and nucleon phase space parameters \( z_i \)'s. We have made the non-relativistic approximation in the calculation of pionic state norm,

\[ N_{\pi}(f, g) \simeq \exp \left[ \frac{2m_{\pi} c^2}{\hbar^3} \int dr \hat{\varphi}_f(r) \hat{\varphi}_g(r) \right] . \]  

(2.16)

The Hamiltonian form of Eq. (2.11) with pion-nucleon \( P \)-wave interaction is the simplest one. In addition to coupling with charged pions, higher dimension terms such as pion-nucleon \( S \)-wave interaction coming from \( \bar{N} \phi^2 N \) coupling would have visible contributions when pions have large expectation values. Charged pions should give similar energy gains to neutral pions, but coherent state treatment of charged pions mixes different charge states, and this charge mixing may lead to serious problems in the discussion of excited levels. In order to overcome this problem, it is necessary to perform the coupled channel calculation of different nucleon and pion charge states or to perform isospin projection. These are beyond the scope of this paper, and will be discussed elsewhere. On the other hand, higher dimension terms such as the pion-nucleon \( S \)-wave interaction are not expected to give large contributions in energy, since the number of pions is around 0.1 in \( ^{12}\text{C} \) nuclei in the present framework as is shown later.

§ 3. An Example of Application — \( ^{12}\text{C} \) Nucleus —

We have applied AMD with coherent state pion to the study of \( ^{12}\text{C} \) nuclei. In the ground state of \( ^{12}\text{C} \), nucleons occupy the single particle states of \( 0s_{1/2} \) and \( 0p_{3/2} \), both of which are the yrast single particle states, then the pionic effects are expected
to be large. While the 3α cluster model generally describes the excited levels of this nucleus very well, the first excited state ($2^+_1$) is calculated to be too low. Since the "spin-orbit" splitting of the $0p_{3/2}$ and $0p_{1/2}$ single particle states is responsible to the $0^+_1-2^+_1$ level spacing, pionic effect to push up the $0p_{1/2}$ level may appear as the increase of $E^*(2^+_1)$ in $^{12}$C. It would be also interesting to study un-natural parity levels such as $0^-, 1^+, 2^-, ...$, whose excitation energy might decrease due to the coupling to the natural parity nucleon state with $0^-$ pionic state.

In this paper, we apply the simplest model of AMD with coherent state pions as the first step. We show the results with projection after variation (PAV); we first construct the intrinsic state by using the cooling variational method for the parametrized wave function of Eq. (2.1), and projection to specified $J^\pi$ has been carried out from the prepared intrinsic state. We find that the effects of parity projection before variation in $^{12}$C nuclei are not large when pions are included explicitly, while it has been found to be important for spectroscopic studies of light nuclei without pions.\textsuperscript{5,11} In the cooling stage of intrinsic energy, the imaginary part of $\varphi_f(r)$ is a redundant degree of freedom as is clear from Eqs. (2.13) and (2.14) with $\varphi_f = \varphi_g$, then we cannot control the imaginary part which is given randomly in the initial state of variation. Thus we have limited the pion eigen function $\varphi(r)$ to be real. We find that there are many local minima especially for small $f_{\pi N}$ values, then we have selected the lowest energy wave functions from several candidates obtained from different random seeds. Since we do not take account of the isospin projection which enhances the tensor force effect by around three times, we use larger $f_{\pi N}$ value in the range $f_{\pi N} = 1 \sim \sqrt{3}$. For finite nuclei, we use the scaled coupling constant $f_{\pi N}(A) = \sqrt{(A-1)/Af_{\pi N}}$ in order to include approximately the effects of the Fock (exchange) term of OPEP, which requires quantum corrections in a field description.

As the effective nucleon-nucleon interaction, we start from Brink–Boeker type two range Gaussian interactions, which approximately reproduces the binding energies of $^4$He, $^{16}$O (28.4 and 128.9 MeV, BBO\textsuperscript{13}) or the binding energy of $^4$He and nuclear matter saturation (28.1 MeV and $E/A = -16.8$ MeV at $\rho = 0.165\text{fm}^{-3}$, BBO2). Since the potential energy from pions is very large, it would be important for nucleon-nucleon interaction to have saturation property in order to avoid collapsing. The interaction range is chosen to be shorter than that of, for example, the Volkov interaction. When we include pions, we employ the nucleon-nucleon interaction BBO$\pi$, which is a little modified from BBO2 to reproduce the ground state energy and the first excited state energy of $^{12}$C. Parameters of these interactions are summarized in Table I.

| $\mu_1$ (fm) | $v_1$ (MeV) | $M_1$ | $\mu_2$ (fm) | $v_2$ (MeV) | $M_2$ | $f_{\pi N}$ |
|-------------|-------------|-------|-------------|-------------|-------|-------------|
| Volkov      | 1.6         | -83.34| 0.575       | 0.82        | 144.86| 0.575       |
| BBO\textsuperscript{13} | 1.2         | -253.798| 0.2186     | 0.6         | 924.631| -1.551     |
| BBO2        | 1.2         | -258.3 | 0.25        | 0.6         | 950.00| -1.658     |
| BBO$\pi$    | 1.2         | -256.0 | 0.25        | 0.6         | 950.00| 1.63       |
In Fig. 1 we show the intrinsic state energy as a function of the pion number expectation value,

\[
n_\pi = \frac{\langle \Phi_\pi(f) | \int dk \hat{a}^\dagger(k) \hat{a}(k) | \Phi_\pi(f) \rangle}{\langle \Phi_\pi(f) | \Phi_\pi(f) \rangle} \approx \frac{2m_\pi c^2}{\hbar^3} \int dr \tilde{\varphi}_f(r) \varphi_f(r). \tag{3.1}
\]
At small $f_{\pi N}$ values around one, pure nucleon state is energetically favored. When we increase $f_{\pi N}$, pion-nucleon interaction gives very large binding, and the optimal state has finite pions for $f_{\pi N} \gtrsim 1.6$. The total pionic energy amounts to be around $-90$ MeV in the case $f_{\pi N} = 1.63$ as shown in Fig. 2. In well developed pionic states, the nucleus loses energy in the nucleon part $\mathcal{H}_{NN}$ instead of gaining pion-nucleon interaction energy efficiently. This feature is similar to the case of pion condensation in high density nuclear matter.\(^{14}\)

\[\text{Fig. 3. Pion number dependence of the total energy (top), energy difference from } 0^+1 \text{ state (middle), and the nucleonic abnormal parity probability (bottom) in } ^{12}\text{C. Filled circles show the energy minimum points for each } J^\pi.\]

In the upper panel of Fig. 3, we show the results of total energy after $J^\pi$ projection from the cooled intrinsic wave functions under $n_\pi$ constraint. All the $J^\pi$ states have their minima at finite $n_\pi$ when we adopt $f_{\pi N} = 1.63$. It is interesting to find that natural parity states favor smaller $n_\pi$, and un-natural parity states favor larger
Finite number pion is expected to act as the LS-like interaction and to increase the excitation energy of \(2^+_1\). In the middle panel of Fig. 8 we show the energy difference \(E(J^\pi) - E(0^+_1)\) as functions of \(n_\pi\) for \(f_{\pi N} = 1.63\). At \(n_\pi = 0\) where the present model is equivalent to the normal AMD, \(2^+_1\) has small excitation energies, which is a feature of \(\alpha\) cluster models. At around \(n_\pi \simeq 0.05\), \(2^+_1\) state starts to go up. This increase of energy difference at \(n_\pi > 0.1\) is not a consequence of the nuclear shrinkage, but the result of pionic LS-like effect. Actually we find that the calculated rms radius grows in the region \(n_\pi \gtrsim 0.07\).

Contrary to the positive parity rotational states, \(0^-_1\) and \(1^+_1\) states go down as the pion number increases. This is due to the coupling to the nucleonic abnormal parity states such as,

\[
|\Psi(0^-)⟩ = |Ψ_N(0^-)⟩ + |Ψ_N(0^+)⟩ \otimes |Φ_\pi(0^-)⟩,
\]

for the \(0^-\) state. In order to demonstrate this point, we show the nucleonic abnormal parity probability \(P_{N}^{Abn}\), in the bottom panel of Fig. 8. At zero pion number \(n_\pi = 0\), all the states should be purely described in nucleonic state \((P_{N}^{Abn} = 0)\), but the probability goes up to around 5\% and 3\% at the projected energy minima for \(0^-\) and \(1^+\) states, respectively. Other rotational levels are also calculated to contain the abnormal nucleonic parity component of around 1\%. If these are true, it may be interesting to observe pion knock-out reaction, which leaves the nucleus in the un-natural parity states.

Table II. Energy components in \(^{12}\)C levels. For \(0^+\) state, calculated total energies are shown in the parentheses. All the energies are shown in MeV.

| \(J^\pi\) | \(E^*\) | \(T_N\) | \(V_c\) | \(V_{Coul}\) | \(V_{LS}\) | \(E_\pi\) |
|---------|---------|---------|--------|-----------|----------|--------|
| Volkov  |         |         |        |           |          |        |
| 0^+     | (-92.4) | 234.3   | -320.6 | 8.9       | -15.0    | -       |
| 2^+     | 4.2     | 234.5   | -320.6 | 8.8       | -11.1    | -       |
| 4^+     | 12.0    | 235.4   | -317.5 | 8.8       | -7.2     | -       |
| 3^-     | 19.5    | 244.5   | -314.1 | 8.7       | -12.0    | -       |
| BBO\pi  |         |         |        |           |          |        |
| 0^+     | (-92.4) | 187.5   | -228.8 | 7.2       | 0.01     | -58.2  |
| 2^+     | 4.6     | 190.9   | -233.3 | 7.2       | -0.03    | -52.5  |
| 4^+     | 10.6    | 196.3   | -224.9 | 7.1       | -0.15    | -60.1  |
| 1^+     | 17.0    | 198.1   | -195.3 | 7.0       | 0.25     | -89.6  |
| 0^-     | 27.0    | 204.0   | -183.5 | 7.0       | 0.48     | -89.1  |
| 3^-     | 8.7     | 194.7   | -228.9 | 7.2       | -0.04    | -56.5  |

The ground state and the first \(2^+\) state energy can be reproduced in AMD without pion effects when we adopt strong LS interactions,\(^{15}\) but the wave functions in these two descriptions are very different. In Table 11 we compare the energy components for \(^{12}\)C levels in AMD with Volkov interaction with strong LS interaction \((V_{LS} = 1800\text{ MeV})\) and in the present model with moderate LS interaction \((V_{LS} = 900\text{ MeV})\). The energy difference of \(0^+\) and \(2^+\) mainly comes from the LS interaction in the case without pions, while the pionic energy is the main source of the energy difference when pions are included. In addition, it is interesting to find that the LS interaction acts in the reverse way — LS interaction is more attractive for \(0^+\)
without pions, but it weakly acts repulsively for $0^+$ with pions.

§4. Summary

In this paper, we have developed a new framework to include pionic degrees of freedom in the nuclear many-body systems, Antisymmetrized Molecular Dynamics (AMD) with coherent state pion. Compared to the mean field treatment of pions, the present model has a merit that we can evaluate the transition matrix elements of the pion-nucleon coupling term containing the second quantized pion operator $\hat{\phi}$. This enables us to calculate the excitation energies for specified $J^P$ with explicit pion degrees of freedom, through the parity and angular momentum projection from the intrinsic state. It is also interesting to find that the pion coherent state has a norm, and the pionic state overlap $\langle \Phi_\pi(f) | \Phi_\pi(g) \rangle$ reduces the total state overlap, $\langle \Psi(Z,f) | \Psi(Z',g) \rangle$.

We have applied this model to the study of $^{12}$C structure. The LS-like effects of pions can be seen as the increase of the $2^+_1$ state excitation energy. It is also suggested that explicit pionic state $| \Psi_N(0^+) \rangle \otimes | \Phi_\pi(0^-) \rangle$ can admix to the $0^-$ state with around 5% when we adopt $f_{\pi N} = 1.63$.

In order to obtain firm conclusions on the explicit pionic effects in nuclear structure, further theoretical and experimental developments are mandatory. First, charged pions should be included in the framework. Combined with the charge (isospin) projection, it is expected to enhance the tensor effect by three times. We have simulated this enhancement by increasing the pion-nucleon coupling $f_{\pi N}$, but the exchange of proton and neutron may lead to non-trivial effects which cannot be mimicked by increasing $f_{\pi N}$. Next, we should take care of the exchange term (Fock term) and the zero range ($\delta$ type) part of the OPEP in a more reliable manner than to scale the coupling constant by $\sqrt{(A-1)/A}$ times. Including the Landau-Migdal interaction may be an efficient way for this problem. Extension of the wave function is also an important direction. Since we have assumed that the total wave function is a product of nucleonic state and pionic coherent state, nucleonic part of the wave function is common to the zero and one pion states. Thus the nucleonic part of the wave function has to contain both of the $T = 0$ and $T = 1$ states, which may lead to the underestimate of the binding energy. Works in these directions are in progress.

Acknowledgements

We would like to thank Prof. S. Okabe, Prof. K. Katō, Prof. K. Ikeda, Prof. H. Toki and Dr. T. Myo for useful discussions. This work is supported in part by the Ministry of Education, Science, Sports and Culture, Grand-in-Aid for Scientific Research (C)(2), No. 15540243, 2003.

References

1) T. Kunihiro, T. Takatsuka, R. Tamagaki, T. Tatsumi, Prog. Theor. Phys. Suppl. 112 (1993), 123.
2) T. Wakasa et al., Phys. Rev. C 55 (1997), 2909.
3) T. Suzuki and H. Sakai, Phys. Lett. B 455 (1999), 25.
4) S. C. Pieper, K. Varga and R. B. Wiringa, Phys. Rev. C 66 (2002), 044310.
5) H. Toki, S. Sugimoto and K. Ikeda, Prog. Theor. Phys. 108 (2002), 903.
6) Y. Ogawa, H. Toki, S. Tamenaga, H. Shen, A. Hosaka, S. Sugimoto and K. Ikeda, Prog. Theor. Phys. 111 (2004), 75.
7) T. Myo, K. Katō, K. Ikeda, Prog. Theor. Phys. 113 (2005), 763.
8) S. Sugimoto, K. Ikeda and H. Toki, Nucl. Phys. A 740 (2004), 77.
9) R.D. Amado et al., Phys. Rev. C 50 (1994), 640.
10) A. Ono et al., Prog. Theor. Phys. 87 (1992), 1185.
11) Y. Kanada-En’yo, H. Horiuchi, and A. Ono, Phys. Rev. C 52 (1995), 628.
12) A. Dote, H. Horiuchi, Y. Akaishi, and T. Yamazaki, Phys. Rev. C 70 (2004), 044313.
13) S. Okabe, private communication.
14) T. Takatsuka, R. Tamagaki and T. Tatsumi, Prog. Theor. Phys. Suppl. 112 (1993), 67.
15) Y. Kanada-En’yo, Phys. Rev. Lett. 81 (1998), 5291.