Role of deformation in odd-even staggering in reaction cross sections for $^{30,31,32}$Ne and $^{36,37,38}$Mg isotopes

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We discuss the role of pairing anti-halo effect in the observed odd-even staggering in reaction cross sections for $^{30,31,32}$Ne and $^{36,37,38}$Mg isotopes by taking into account the ground state deformation of these nuclei. To this end, we construct the ground state density for the $^{30,31}$Ne and $^{36,37}$Mg nuclei based on a deformed Woods-Saxon potential, while for the $^{32}$Ne and $^{38}$Mg nuclei we also take into account the pairing correlation using the Hartree-Fock-Bogoliubov method. We demonstrate that, when the one-neutron separation energy is small for the odd-mass nuclei, a significant odd-even staggering still appears even with finite deformation, although the degree of staggering is somewhat reduced compared to the spherical case. This implies that the pairing anti-halo effect in general plays an important role in generating the odd-even staggering in reaction cross sections for weakly bound nuclei.

I. INTRODUCTION

The halo structure is one of the most important phenomena in neutron-rich nuclei.$^1, 2$ This phenomenon is characterized by a spatially extended density distribution originated from weakly bound valence neutron(s). This was first discovered by Tanihata et al., who observed considerably large interaction cross sections for $^{11}$Li, $^{11}$Be, and $^{14}$Be$^3, 4$. Subsequently, a narrow momentum distribution was also discovered$^5$ for a weakly bound nucleus, $^{11}$Li, establishing the concept of the halo structure. The heaviest halo nucleus discovered so far is $^{37}$Mg.$^5, 7$

For weakly bound nuclei with two valence neutrons, the pairing correlation between the valence neutrons may quench the halo structure$^8$. That is, the pairing correlation alters the asymptotic behavior of the wave function for the valence neutrons, reducing the divergence feature of nuclear radii for s and p waves at zero separation energy$^8, 12$. This effect is referred to as the pairing anti-halo effect, which can also be viewed as a generation of a spatially localized wave packet of quasi-particles originated from a coherent scattering of the valence neutrons to the continuum spectrum caused by the pairing interaction$^{13}$.

In the previous publications, we have argued that the pairing anti-halo effect plays an important role in the odd-even staggering observed in interaction cross sections$^{14, 16}$. That is, the experimental data have often shown a large odd-even staggering in interaction and reaction cross sections, in which cross sections for odd-mass nuclei are systematically larger than those for the neighboring even-mass nuclei$^{14, 17}$. Using the Hartree-Fock-Bogoliubov (HFB) method with spherical symmetry, we have shown that the observed odd-even staggering can be largely accounted for in terms of the pairing anti-halo effect (see also Refs. $^{18, 19}$).

In this paper, we extend our previous analyses by taking into account the ground state deformation of weakly bound nuclei. To this end, we study the odd-even staggering in the $^{30,31,32}$Ne and $^{36,37,38}$Mg isotopes, for which the $^{31}$Ne and $^{37}$Mg nuclei have been suggested to have a deformed halo structure with p wave$^3, 7, 17, 20, 28$.

There are two possible effects of nuclear deformation on the odd-even staggering. Firstly, several angular momentum components are mixed in a deformed single-particle wave function for a valence neutron, reducing the s and p wave components in the wave function. This will reduce the radius of the $^{31}$Ne and $^{37}$Mg nuclei, somewhat quenching the odd-even staggering in the interaction and reaction cross sections. Secondly, the deformation may change the level density around the Fermi level, which would result in either an enhancement or a decrease of the pairing correlation, depending on the position of the Fermi surface. This would eventually influence the magnitude of the pairing anti-halo effect, thus the cross sections for $^{32}$Ne and $^{38}$Mg. The primary aim of this paper is to investigate how these two effects interplay with each other in actual cases and how the conclusion obtained in our previous analyses based on spherical symmetry is altered if the deformation is explicitly taken into account.

The paper is organized as follows. In Sec. II, we briefly summarize the theoretical frameworks. In our calculations, we first generate the deformed ground state density using the HFB method with Woods-Saxon potentials, which is then used as an input to the Glauber theory in order to compute reactions cross sections. In Sec. III, we apply these frameworks to reaction cross sections for the $^{30,31,32}$Ne and $^{36,37,38}$Mg nuclei and discuss the role of deformation in the odd-even staggering in the reaction cross sections. We then summarize the paper in Sec. IV.
II. THEORETICAL FRAMEWORKS

A. Deformed density

We analyze the reaction cross sections for the $^{30,31,32}$Ne and $^{36,37,38}$Mg nuclei. Symbolically, we denote the three isotopes in each element as $A$, $A+1$, and $A+2$ systems, respectively. Our first task is to construct the ground state density of each nucleus by taking into account the deformation. For simplicity, we ignore the pairing correlation in the $A$ and $A+1$ systems, and construct the density distribution by putting the nucleons into the lowest $A$ and $A+1$ single-particle orbits in a deformed Woods-Saxon potential, respectively (we have confirmed that the reaction cross section for the $^{30}$Ne and $^{36}$Mg nuclei does not significantly change even if the pairing correlation is taken into account).

For the $A+2$ systems, on the other hand, we take into account the pairing correlation with the HFB method. In the coordinate space representation, the HFB equations read [29–31]

$$\begin{pmatrix} \hat{h} - \lambda + \Delta(r) \\ \Delta(r) - \hbar^2/2m \nabla^2 + V(r) \end{pmatrix} \begin{pmatrix} U_i(r) \\ V_i(r) \end{pmatrix} = E_i \begin{pmatrix} U_i(r) \\ V_i(r) \end{pmatrix},$$

where $\Delta(r)$ is the pairing potential, and $\lambda$ and $E_i$ are the chemical potential and a quasi-particle energy, respectively. $\hat{h}$ is a mean-field Hamiltonian given by

$$\hat{h} = -\frac{\hbar^2}{2m} \nabla^2 + V(r),$$

where $V(r)$ is a mean-field potential and $m$ is the nucleon mass. Here, we have assumed that the nucleon-nucleon interaction is zero range, so that the mean-field and the pairing potentials are both local. In this framework, the density distribution is given by

$$\rho(r) = \sum_i |V_i(r)|^2.$$ 

The chemical potential $\lambda$ is determined so that the average particle number is $A+2$, that is,

$$\int dr \rho(r) = A+2.$$ 

For simplicity, we use a deformed Woods-Saxon potential for $V(r)$ and $\Delta(r)$ given by,

$$V(r) = V_0 \left( f(r) - R_0 \beta_2 \frac{df(r)}{dr} Y_{20}(\theta_{rd}) \right) + V_h r \frac{df(r)}{dr} (l \cdot s),$$

$$\Delta(r) = \Delta_0 \left( f(r) - R_0 \beta_2 \frac{df(r)}{dr} Y_{20}(\theta_{rd}) \right)$$

with

$$f(r) = \frac{1}{1 + \exp[(r - R_0)/a]}.$$ 

Here, we have assumed axially symmetric quadrupole deformation with the deformation parameter of $\beta_2$, and denoted the angle between the vector $r$ and the symmetry axis as $\theta_{rd}$. In the single-particle potential, we take into account only the spherical part of the spin-orbit potential. For the protons, we also add the spherical Coulomb interaction to the mean-field potential, $E_i$, with the radius of $R_0$.

We use the same values for the parameters for the single-particle potential as those listed in Table I in Ref. [32], except for the depth parameter $V_0$ for the configuration for the valence orbit, for which we adjust the value of $V_0$ so that the neutron separation energy for the $A+1$ nuclei is reproduced. For simplicity, we use the same value for the deformation parameter for all the three isotopes, $A$, $A+1$, and $A+2$. For the strength $\Delta_0$ for the neutron pairing potential, $\Delta_0$, we determine it according to

$$\bar{\Delta} \equiv \frac{\int dr \Delta(r)f(r)}{\int dr f(r)},$$

with the average pairing gap of $\bar{\Delta} = 12/\sqrt{A+2}$ MeV. For the protons, we ignore the pairing correlation as they are deeply bound in the nuclei considered in this paper, and thus the effect of pairing correlation on the nuclear radius is expected to be small.

We solve the HFB equations, (1), by expanding the upper and the lower components of the quasi-particle wave functions as,

$$U_K(r) = \sum_{n,l,j} u_{nlj} \psi_{nljK}(r),$$

$$V_K(r) = \sum_{n,l,j} v_{nlj} \psi_{nljK}(r),$$

where $\{\psi_{nljK}(r)\}$ are eigen-functions of the single-particle Hamiltonian $\hat{h}$ when the deformation parameter $\beta_2$ is set to zero. This wave function is characterized by the radial quantum number $n$, the orbital angular momentum $l$, the total single-particle angular momentum $j$, and its projection onto the symmetry axis, $K$. Notice that $K$ is conserved in the quasi-particle wave functions, since we assume the axial symmetry. In Eqs. (9) and (10), we include the continuum states up to 30 MeV above the Fermi energy, $\lambda$, by discretizing them with the box boundary condition at $R_{box} = 60$ fm.

B. Reaction cross sections

We next consider collision of a deformed projectile nucleus with a spherical target nucleus, and compute the reaction cross sections, $\sigma_R$. To this end, we employ the Glauber theory, which is based on the eikonal approximation and the adiabatic approximation to the nucleonic motions [33, 34]. In order to calculate reaction cross sections, we also apply the adiabatic approximation to the
rotational motion of a deformed nucleus. That is, we first fix the orientation angle of the deformed nucleus and then take an average of the resultant cross section over all the orientation angles. The reaction cross sections are thus expressed as

$$\sigma_R = \frac{1}{4\pi} \int d\Omega \sigma_R(\Omega),$$

(11)

where $\Omega$ is the angle of the symmetric axis of the deformed nucleus in the laboratory frame, and $\sigma_R(\Omega)$ is the reaction cross section for fixed $\Omega$.

In the Glauber theory, the reaction cross section is computed as

$$\sigma_R(\Omega) = \int db \left( 1 - |e^{i\chi(b,\Omega)}|^2 \right),$$

(12)

where $b$ is the impact parameter and the phase shift function $\chi$ is given by

$$i\chi(b,\Omega) = -\int dr \rho_P(r;\Omega) \times \left[ 1 - \exp\left( -\int dr' \rho_T(r') \Gamma_{NN}(s-s'+b) \right) \right].$$

(13)

Here, $s$ and $s'$ are the transverse component (that is, the component which is parallel to $b$) of $r$ and $r'$, respectively. $\Gamma_{NN}$ is the profile function for the $NN$ scattering, which we assume to be

$$\Gamma_{NN}(b) = \frac{1 - i\alpha}{4\pi\beta} \sigma_{NN} \exp\left( -\frac{b^2}{2\beta} \right),$$

(14)

where $\sigma_{NN}$ is the total $NN$ cross section. In Eq. 13, $\rho_P$ and $\rho_T$ are the density distribution for the projectile and the target nuclei, respectively. We assume that the target density is spherical, $\rho_T(r') = \rho_T(r)$, while the projectile density has axial symmetry, that is,

$$\rho_P(r;\Omega) = \rho_P(r,\theta_{rd}).$$

(15)

The projectile density can be expanded into multipoles as

$$\rho_P(r,\theta_{rd}) = \sum_{\lambda} \rho^{(P)}_\lambda(r) Y_{\lambda 0}(\theta_{rd});$$

$$= \sum_{\lambda=0}^\infty \sqrt{\frac{4\pi}{2\lambda+1}} \rho^{(P)}_\lambda(r) Y_{\lambda 0}(\Omega).$$

(17)

Notice that the phase shift function given by Eq. 13 takes into account the effect beyond the optical limit approximation following the prescription proposed in Ref. 43. We evaluate it using the Fourier transform method.

In this paper, we analyze the experimental data at incident energy $E = 240$ MeV/nucleon with $^{12}$C target. We use the same density for $^{12}$C as that given in Ref. 33, while we use the same parameters given in Ref. 45 for the profile function, $\Gamma_{NN}$.

III. ODD-EVEN STAGGERING IN REACTION CROSS SECTIONS

A. $^{30,31,32}$Ne isotopes

Let us now numerically evaluate the reaction cross section for deformed nuclei and discuss the role of deformation in the odd-even staggering in the reaction cross sections. We first consider the $^{30,31,32}$Ne isotopes, for which the odd-even staggering has been investigated assuming spherical symmetry.

In Ref. 24, we have shown that the measured reaction cross section for the $^{31}$Ne nucleus can be reproduced both with the particle-rotor model and with the Nilsson model with a deformed Woods-Saxon potential when the quadrupole deformation parameter is in the range of $0.17 \leq \beta_2 \leq 0.33$. In this case, the valence neutron in $^{31}$Ne occupies the $[330 1/2^+ \ K^\pi = 1/2^+]$ orbit, which is connected to the $1f_{1/2}$ level in the spherical limit.

In this paper, we therefore choose $\beta_2 = 0.3$. As has been demonstrated in Ref. 24, the dependence of reaction cross sections on the deformation parameter is weak once the configuration of the valence orbit is fixed.

Figure 1 shows the root mean square radii so obtained for the $^{30,31,32}$Ne nuclei as a function of the energy of the valence orbit for the $^{31}$Ne nucleus, $\epsilon_n$. To draw this figure, we vary the depth parameter, $V_0$, in the Woods-Saxon potential for the $K^\pi = 1/2^+$ configuration. The dotted, the dashed, and the solid lines show the radius for the $^{30}$Ne, $^{31}$Ne, and $^{32}$Ne nuclei, respectively. One can
see that the radius of the $^{31}$Ne increases rapidly as the one neutron separation energy, $S_n = -\epsilon_n$, approaches to zero, due to the $p$ wave component in the wave function for the valence neutron. On the other hand, the radius of the $^{32}$Ne nucleus varies slowly as a function of the one neutron separation energy and becomes smaller than that of the $^{31}$Ne nucleus for $\epsilon_n \leq -0.42$ MeV due to the pairing anti-halo effect. This behavior is qualitatively the same as in the previous analysis shown in the middle panel of Fig. 2 in Ref. [14], which was based on the spherical symmetry of the Ne isotopes.

The reaction cross sections for the $^{30,31,32}$Ne nuclei evaluated at $S_n^{(31Ne)} = 0.3$ MeV are shown in Fig. 2. These are compared with the experimental interaction cross sections [17], which are expected to be close to the reaction cross sections for neutron-rich nuclei [7, 38, 46]. For comparison, we also show the result of the previous analysis [14] at a similar one neutron separation energy, the staggering parameter $\gamma_3$ defined in Eq. (18) for the $^{30,31,32}$Ne + $^{12}$C reactions at $E = 240$ MeV/nucleon. It is plotted as a function of the single-particle energy for the valence orbit for the $^{31}$Ne nucleus. The dashed and the solid lines are obtained using the spherical and the deformed density distributions, respectively. The experimental data is evaluated using the measured interaction cross sections shown in Ref. [15].

![FIG. 2: The reaction cross sections for the $^{30,31,32}$Ne + $^{12}$C reaction at $E = 240$ MeV/nucleon. These are evaluated at the one neutron separation energy of $S_n = 0.3$ MeV. The filled circles with error bars indicate the experimental interaction cross sections taken from Ref. [17]. For comparison, the result of the previous analysis based on the spherical density distributions is also shown by the dashed line.](image1)

The staggering parameter $\gamma_3$ defined as [15],

$$\gamma_3 \equiv \frac{(-1)^A}{2} \left[ \sigma_R(A + 1) - 2\sigma_R(A) + \sigma_R(A - 1) \right],$$

where $\sigma_R(A)$ is the reaction cross section of a nucleus with mass number $A$. The dashed and the solid lines in the figure show the staggering parameter for the spherical and the deformed cases, respectively. For a fixed value of separation energy, the staggering parameter $\gamma_3$ is smaller in the deformed case as compared to the spherical case, which is consistent with Fig. 2. However, the staggering parameter increases as the separation energy decreases, and eventually comes closer to the central value of the experimental data when $\epsilon_n$ is around the threshold.

In Sec. I, we have conjectured that the nuclear deformation may lead to two effects on reaction cross sections. One is to decrease the cross section for $^{31}$Ne due to the admixture of several angular momentum components in the single-particle wave function for the valence orbit. The other is to change the cross section for $^{32}$Ne because of a change in the degree of pairing anti-halo effect. The former effect would reduce the staggering, while the latter effect either enhances or reduces it depending on the level density around the Fermi surface. Our calculation shown in Fig. 2 indicates that the former effect indeed exists, while the latter effect is much less clear. In order to shed light on the latter effect, Fig. 4 shows the dependence of the reaction cross section on the strength of the pairing correlation. To this end, we vary the average pairing gap, $\Delta$, defined by Eq. (5). The filled squares with the dashed, the solid, and the dotted lines are obtained by setting the average pairing gap to be $1, 2.1 \left(=\frac{12}{\sqrt{32}}\right)$, $3 \left(=\frac{18}{\sqrt{32}}\right)$, and $4 \left(=\frac{24}{\sqrt{32}}\right)$, respectively. The results are shown in Fig. 4. In the figure, we compare the calculated cross sections with the experimental data [17] for the $^{30,31,32}$Ne + $^{12}$C reactions at $E = 240$ MeV/nucleon. The dashed and the solid lines in the figure show the calculated cross sections for the spherical and the deformed cases, respectively. The experimental data is obtained using the measured interaction cross sections shown in Ref. [15].

![FIG. 3: The staggering parameter $\gamma_3$ defined by Eq. (18) for the $^{30,31,32}$Ne + $^{12}$C reactions at $E = 240$ MeV/nucleon.](image2)
and 3 MeV, respectively. Notice that the solid line is the same as that in Fig. 2. For comparison, the open square with the dashed line shows the result without the pairing correlation. One can clearly see that the reaction cross section for $^{32}$Ne is not sensitive to the value of average pairing gap as long as it is large enough. This would be correlated to the fact that the root-mean-square radius for $^{32}$Ne is not sensitive to the single-particle energy for the valence orbit of $^{31}$Ne, as has been shown in Fig. 1. Even though the occupation probability for the $p$-wave orbital may depend largely on the average pairing gap, the root-mean-square radius does not change much once the radius is significantly shrunk due to the pairing anti-halo effect so that the $s$ and $p$-wave states do not behave abnormally. This indicates that the main effect of nuclear deformation is simply to decrease the odd-even staggering in reaction cross sections, at least for the Ne isotopes.

B. $^{36,37,38}$Mg isotopes

Let us next discuss the $^{36,37,38}$Mg isotopes. For $^{37}$Mg in these isotopes, the $p$-wave halo structure has been suggested from a measurement of the one neutron removal reaction on C and Pb targets, with a small one neutron separation energy of $0.22^{+0.12}_{-0.09}$ MeV. Moreover, the experimental reaction cross sections indicate a large odd-even staggering for $^{36,37,38}$Mg.

We first determine the deformation parameter $\beta_2$ for these isotopes. With the deformed Woods-Saxon potential shown in Sec. II-A, together with the parameters listed in Table I in Ref. [32], we find that the valence neutron in $^{37}$Mg occupies the $[312 5/2]$ orbit for $\beta_2 \leq 0.4$, while it occupies the $[321 1/2]$ orbit for $0.4 \leq \beta_2 \leq 0.6$. The former is connected to the $1f_{7/2}$ level while the latter to the $2p_{3/2}$ level in the spherical limit. The former state has $K^\pi = 5/2^-$, and thus contains angular momenta larger than $l = 3$, which do not form a halo structure. In contrast, the latter state contains a large $p$-wave component, being consistent with the suggested halo structure for $^{37}$Mg. We therefore choose $\beta_2 = 0.5$ in the analysis shown below.

Figure 5 shows the root-mean-square radii for the $^{36,37,38}$Mg nuclei as a function of the single-particle energy for the valence orbit of $^{37}$Mg. The radii behave qualitatively the same as those for the Ne isotopes shown in Fig. 1. That is, the radius of $^{37}$Mg diverges in the limit of vanishing single-particle energy, while that of $^{38}$Mg varies much more slowly due to the pairing anti-halo effect.

The reaction cross sections for the $^{36,37,38}$Mg isotopes are shown in Fig. 6 for two different values of the one neutron separation energy, $S_n$, for $^{37}$Mg. The solid line is obtained with $S_n = 0.32$ MeV, while the dashed line with $S_n = 1.5$ MeV. The experimental odd-even staggering can be well reproduced with $S_n = 0.32$ MeV. On the other hand, for $S_n = 1.52$ MeV, the reaction cross section increases monotonically as a function of mass number, that is inconsistent with the experimental data. Notice that this behavior is qualitatively similar to the odd-even staggering for the Ne isotopes shown in Fig.3 in Ref. [14] obtained with the spherical densities. Therefore it is evident that the pairing anti-halo effect plays an important role in the odd-even staggering of deformed neutron-rich nuclei, such as $^{30,31,32}$Ne and $^{36,37,38}$Mg isotopes.

IV. SUMMARY

We have investigated the role of nuclear deformation in the odd-even staggering observed in reaction cross sections for several systems. To this end, we have used
FIG. 6: The reaction cross sections for the $^{36,37,38}\text{Mg} + ^{12}\text{C}$ reaction at $E = 240\text{ MeV/nucleon}$. These are evaluated for the quadrupole deformation parameter of $\beta_2 = 0.5$, at which the valence neutron in $^{37}\text{Mg}$ occupies the $[321,1/2]$ orbital. The solid and the dashed lines denote the results for $S_n = 0.32$ and 1.5 MeV, respectively, where $S_n$ is the one neutron separation energy of $^{37}\text{Mg}$. The experimental reaction cross sections are taken from Ref. [7].

the deformed Hartree-Fock-Bogoliubov method to take into account both the deformation and the pairing effects. We have applied this method to the $^{30,31,32}\text{Ne}$ and $^{36,37,38}\text{Mg}$ isotopes and have shown that the deformation mainly decreases the degree of odd-even staggering due to the admixture of several angular momentum states in a deformed single-particle wave function. Despite this, the odd-even staggering persists even with finite deformation, when the one neutron separation energy is small enough. In particular, we have successfully accounted for the experimental odd-even staggering both for the $^{30,31,32}\text{Ne}$ and $^{36,37,38}\text{Mg}$ isotopes within the unified theoretical framework. This strongly indicates that the pairing anti-halo effect indeed has a responsibility to the observed odd-even staggering in reaction cross sections.

Our calculation can be improved in several ways. For instance, in this paper, we have assumed that the deformation parameter is the same for the three nuclei within the same element. It might be important to take into account an iso-
topic dependent deformation, as has been predicted e.g., by the anti-symmetrized molecular dynamics (AMD) [24,25], although the dependence of the reaction cross sections on the deformation would not be large once the single-particle configuration is fixed. Another issue is the treatment of pairing for the odd-mass nuclei. For simplicity, in this paper we have neglected the pairing correlation in $^{30,31}\text{Ne}$ and $^{36,37}\text{Mg}$, because the effect of the pairing on the radius of $^{30}\text{Ne}$ and $^{36}\text{Mg}$ had turned out to be small. However, if one regards $^{31}\text{Ne}$ and $^{37}\text{Mg}$ as one quasi-particle excitation on top of $^{30}\text{Ne}$ and $^{36}\text{Mg}$, the pairing might play some role in these nuclei as well. A more consistent way towards this end would be to treat the odd-mass nuclei using the blocked HFB method [26,48], that would be an interesting future work.

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[1] I. Tanihata, H. Savajols, and R. Kanungo, Prog. Part. Nucl. Phys. 68, 215 (2013).
[2] K. Hagino, I. Tanihata, and H. Sagawa, in 100 Years of Subatomic Physics, ed. by E.M. Henley and S.D. Ellis (World Scientific, Singapore, 2013), p. 231.
[3] I. Tanihata et al., Phys. Rev. Lett. 55, 2676 (1985).
[4] I. Tanihata et al., Phys. Lett. B206, 592 (1988).
[5] T. Kobayashi et al., Phys. Rev. Lett. 60, 2599 (1988).
[6] N. Kobayashi et al., Phys. Rev. Lett. 112, 242501 (2014).
[7] M. Takechi et al., Phys. Rev. C90, 061305(R) (2014).
[8] K. Bennaceur, J. Dobaczewski, and M. Plусzajczak, Phys. Lett. B496, 154 (2000).
[9] K. Riisager, A.S. Jensen, and P. Møller, Nucl. Phys. A548, 393 (1992).
[10] A.S. Jensen, K. Riisager, D.V. Fedorov, and E. Garrido, Rev. Mod. Phys. 76, 215 (2004).
[11] H. Sagawa, Phys. Lett. B286, 7 (1992).
[12] H. Sagawa and K. Hagino, Eur. Phys. J. A51, 102 (2015).
[13] K. Hagino and H. Sagawa, Phys. Rev. C95, 024304 (2017).
[14] K. Hagino and H. Sagawa, Phys. Rev. C84, 011303(R) (2011).
[15] K. Hagino and H. Sagawa, Phys. Rev. C85, 014303 (2012).
[16] K. Hagino and H. Sagawa, Phys. Rev. C85, 037604 (2012).
[17] M. Takechi et al., Phys. Lett. B707, 357 (2012).
[18] S. Sasabe, T. Matsumoto, S. Tagami, N. Furutachi, K. Minomo, Y.R. Shimizu, and M. Yahiro, Phys. Rev. C88, 037602 (2013).
[19] T. Matsumoto and M. Yahiro, Phys. Rev. C90, 041602(R) (2014).
[20] T. Nakamura et al., Phys. Rev. Lett. 103, 262501 (2009).
[21] T. Nakamura et al., Phys. Rev. Lett. 112, 142501 (2014).
[22] I. Hamamoto, Phys. Rev. C81, 021304(R) (2010).
[23] Y. Urata, K. Hagino, and H. Sagawa, Phys. Rev. C83, 041303(R) (2011).
[24] Y. Urata, K. Hagino, and H. Sagawa, Phys. Rev. C86, 044613 (2012).
[25] K. Minomo, T. Sumi, M. Kimura, K. Ogata, Y.R. Shimizu, and M. Yahiro, Phys. Rev. C84, 034602 (2011).
[26] K. Minomo, T. Sumi, M. Kimura, K. Ogata, Y.R. Shimizu, and M. Yahiro, Phys. Rev. Lett. 108, 052503 (2012).
[27] T. Sumi, K. Minomo, S. Tagami, M. Kimura, T. Matsumoto, K. Ogata, Y.R. Shimizu, and M. Yahiro, Phys. Rev. C85, 064613 (2012).
[28] S. Watanabe, K. Minomo, M. Shimada, S. Tagami, M. Kimura, M. Takechi, M. Fukuda, D. Nishimura, T. Suzuki, T. Matsumoto, Y.R. Shimizu, and M. Yahiro, Phys. Rev. C89, 044610 (2014).
[29] J. Dobaczewski, W. Nazarewicz, T.R. Werner, J.F. Berger, C.R. Chinn, and J. Dechargé, Phys. Rev. C53, 2809 (1996).
[30] J. Dobaczewski, H. Flocard, and J. Treiner, Nucl. Phys. A422, 103 (1984).
[31] A. Bulgac, [arXiv:nucl-th/9907088].
[32] T. Shoji and Y.R. Shimizu, Prog. Theo. Phys. 121, 319 (2009).
[33] R.J. Glauber, in Lectures in Theoretical Physics, edited by W.E. Brittin (Interscience, New York, 1959) Vol. 1, p. 315.
[34] C.A. Bertulani and P. Danielewicz, Introduction to Nuclear Reactions (IOP Publishing, Bristol, UK, 2004).
[35] J.A. Christley and J.A. Tostevin, Phys. Rev. C59, 2309 (1999).
[36] K. Hagino and N. Takigawa, Prog. Theor. Phys. 128, 1061 (2012).
[37] Y. Ogawa, T. Kido, K. Yabana, and Y. Suzuki, Prog. Theo. Phys. Suppl. 142, 157 (2001).
[38] Y. Ogawa, K. Yabana, and Y. Suzuki, Nucl. Phys. A543, 722 (1992).
[39] W. Horiuchi, Y. Suzuki, B. Abu-Ibrahim, and A. Kohama, Phys. Rev. C75, 044607 (2007).
[40] W. Horiuchi, Y. Suzuki, P. Capel, and D. Baye, Phys. Rev. C81, 024606 (2010).
[41] W. Horiuchi, T. Inakura, T. Nakatsukasa, and Y. Suzuki, Phys. Rev. C86, 024614 (2012).
[42] W. Horiuchi, S. Hatakeyama, S. Ebata, and Y. Suzuki, Phys. Rev. C93, 044611 (2016).
[43] B. Abu-Ibrahim and Y. Suzuki, Phys. Rev. C61, 051601(R) (2000); C62, 034608 (2000).
[44] C.A. Bertulani and H. Sagawa, Nucl. Phys. A588, 667 (1995).
[45] B. Abu-Ibrahim, W. Horiuchi, A. Kohama, and Y. Suzuki, Phys. Rev. C77, 034607 (2008).
[46] A. Kohama, K. Iida, and K. Oyamatsu, Phys. Rev. C78, 061601(R) (2008).
[47] B. Bally, B. Avez, M. Bender, and P.-H. Heenen, Phys. Rev. Lett. 113, 162501 (2014).
[48] T.T. Sun, M. Matsuo, Y. Zhang, and J. Meng, [arXiv:1310.1661 [nucl-th]].