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Faddeev calculation of $3\alpha$ and $\alpha\alpha\Lambda$ systems using $\alpha\alpha$ resonating-group method kernels

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We carry out Faddeev calculations of three-alpha ($3\alpha$) and two-alpha plus $\Lambda$ ($\alpha\alpha\Lambda$) systems, using two-cluster resonating-group method kernels. The input includes an effective two-nucleon force for the $\alpha\alpha$ resonating-group method and a new effective $\Lambda N$ force for the $\Lambda\alpha$ interaction. The latter force is a simple two-range Gaussian potential for each spin-singlet and triplet state, generated from the phase-shift behavior of the quark-model hyperon-nucleon interaction, $fss2$, by using an inversion method based on supersymmetric quantum mechanics. Owing to the exact treatment of the Pauli-forbidden states between two $\alpha$ clusters, the present three-cluster Faddeev formalism can describe the mutually related, $\alpha\alpha$, $3\alpha$, and $\alpha\alpha\Lambda$ systems, in terms of a unique set of the baryon-baryon interactions. For the three-range Minnesota force which describes the $\alpha\alpha$ phase shifts quite accurately, the ground-state and excitation energies of $^9\Lambda$Be are reproduced within 100–200 keV accuracy.

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I. INTRODUCTION

In spite of much effort to incorporate microscopic features of the alpha-alpha ($\alpha\alpha$) interaction, a consistent description of the three-alpha ($3\alpha$) and two-alpha plus $\Lambda$ ($\alpha\alpha\Lambda$) systems has not yet been obtained in the Faddeev formalism. The most favorable description of the $\alpha\alpha$ system is the $\alpha\alpha$ resonating-group method (RGM) [1]. Although some of the previous works deal with the $\alpha\alpha$ RGM kernel explicitly in the $3\alpha$-cluster Faddeev formalism, they usually yield a large overbinding for the ground state and sometimes involve spurious states because of an incomplete treatment of the Pauli-forbidden states in the $3\alpha$ model space [2–5]. Various types of $3\alpha$ orthogonality condition models (OCM) [6–8] also yield a similar overbinding for the ground state, although the effect of the Pauli principle between $\alpha$ clusters is satisfactorily treated in each framework. Only one exception to this rule is the $3\alpha$ OCM in Refs. [9,10], in which the Pauli forbidden components described by the $\alpha\alpha$ bound-state solutions of the deep Buck, Friedrich, and Wheatley (BFW) potential [11] are completely eliminated. The result is rather similar to the traditional $3\alpha$ Faddeev calculation using Ali-Bodmer phenomenological $\alpha\alpha$ potential with a repulsive core [12]. In these calculations, the ground-state energy of the $3\alpha$ system is less than 1.5 MeV, and a simultaneous description of the compact shell-model-like ground state and the excited $0^+$ state with well-developed cluster structure is not possible. The origin of the different conclusions in Refs. [8–10] is spelled out in Ref. [13], in which the existence of almost forbidden Faddeev components inherent to this $3\alpha$ OCM using the bound-state Pauli-forbidden states of the BFW potential is essential.

A possible resolution of this overbinding problem of the $3\alpha$ model is found in our new three-cluster Faddeev formalism, which uses singularity-free $T$-matrices (RGM $T$-matrices) generated from the two-cluster RGM kernels [14]. In this formalism, solving the Faddeev equation automatically guarantees the elimination of the three-cluster redundant components from the total wave function. The explicit energy dependence inherent in the exchange RGM kernel is self-consistently treated. We first applied this formalism to the three-dineutron and $3\alpha$ systems, and obtained complete agreement between the Faddeev calculations and variational calculations using the translationally invariant harmonic-oscillator (h.o.) basis [14,15]. Next, this formalism was applied to a Faddeev calculation of the three-nucleon bound state [16], which employs complete off-shell $T$-matrices derived from the non-local and energy-dependent RGM kernels of the quark-model $NN$ interactions, FSS [17] and $fss2$ [18]. The $fss2$ model yields a triton binding energy $B_t=8.519$ MeV in the 50 channel calculation, when the $np$ interaction is employed for all the $NN$ pairs in the isospin basis [19]. The effect of the charge dependence of the two-body $NN$ interaction is estimated to be $-0.19$ MeV for the triton binding energy [20]. This implies that our result is not overbinding in comparison with the empirical value, $B_t^{\exp} =8.482$ MeV. If we attribute the difference, 0.15 MeV, to the effect of the three-nucleon force, it is by far smaller than the generally accepted values, 0.5–1 MeV [21], predicted by many Faddeev calculations employing modern realistic meson-theoretical $NN$ interactions. We have further applied this three-cluster Faddeev formalism to the hypertriton system [22], in which the quark-model hyperon-nucleon ($YN$) interactions of $fss2$ yield a reasonable result of the hypertriton properties similar to the Nijmegen soft-core potential NSC89 [23]. Most mathematical details for the Faddeev
equations, employed in this calculation, are given in the present paper.

Here we apply the present three-cluster Faddeev formalism to the $aa\Lambda$ model for $^9\text{Be}$. This hypernucleus plays an important role to study the $\Lambda N$ interaction in the $p$-shell $\Lambda$-hypernuclei. From the early time of the hypernuclear study, $^9\text{Be}$ is considered to be a prototype of $\alpha$-cluster structure, in which the two $\alpha$ clusters form a loosely bound subsystem by the effect of the extra $\Lambda$ hyperon [24]. Since the $YN$ interaction is usually weaker than the $NN$ interaction, this system is suitable for studying a subtle structure change of the two-$\alpha$ system from $^8\text{Be}$. In fact, in addition to the $1/2^+$ ground state [25–28] with the $\Lambda$-separation energy $B_{\Lambda}(^9\text{Be})=6.71\pm 0.04$ MeV [29], the recent $\gamma$-ray spectroscopy [30,31] has revealed the existence of two narrow resonances in the excited states, which are supposed to be $5/2^+$ and $3/2^+$ states generated from the small spin-orbit splitting in the weak coupling picture of $^9\text{Be}(2\pi^0)\Lambda$ (spin $S=1/2$). From a theoretical point of view, this is the simplest non-trivial system which requires the Faddeev formalism with two identical particles, involving three Pauli-forbidden states between two clusters. Several model calculations were already done with various frameworks and two-body potentials. Hiyama et al. [7] used the OCM for the $\alpha\alpha$, $3\alpha$ and $aa\Lambda$ systems and discussed not only the ground state of $^9\text{Be}$, but also the spin-orbit splitting of the $5/2^+$ and $3/2^+$ states [32]. They employed simple three-range Gaussian potentials for the $\Lambda N$ interaction based on $G$-matrix calculations [33] of various Nijmegen and Jülich $YN$ one-boson-exchange-potential (OBEP) models. The $\Lambda\alpha$ potentials are generated from these $\Lambda N$ effective potentials by the folding procedure with respect to the $(0s)^4$ h.o. wave function of the $\alpha$ cluster. They introduced a three-$\alpha$ force and adjusted the $YN$ parameters to reproduce the binding energies of the $^{12}\text{C}$ and $^9\text{Be}$ ground states. Filikhin and Gal [34] used the Faddeev and Faddeev-Yukovskov formalism to calculate the $^8\text{Be}$ and $^{10}\Lambda\text{Be}$ ground states. They used the Ali-Bodmer $\alpha\alpha$ potential [12] and the so-called Isle potential [35] for the $\Lambda\alpha$ interaction. They included only $S$ wave in the calculation, and reproduced the $^8\text{Be}$ ground-state energy correctly. However, if one includes higher partial waves the Ali-Bodmer $\alpha\alpha$ potential yields overbinding for $^9\text{Be}$ by more than 0.5 MeV. Oryu et al. [36] carried out an $aa\Lambda$ Faddeev calculation by using the $aa\Lambda$ RGM kernel and various types of $\Lambda\alpha$ potentials in the separable expansion method. Their energy spectrum of $^9\Lambda\text{Be}$ is reasonable, but the treatment of the two-$\alpha$ Pauli principle in the $aa\Lambda$ system is only approximate. Since they neglected the Coulomb force, a detailed comparison between their calculated results and experiment is not possible. Cravo, Fonseca, and Koike [37] performed $aa\Lambda$ and $a\alpha\Lambda$ Faddeev calculations by using many $\alpha\alpha$ and $\Lambda\alpha$ potentials with the Coulomb force included between the two $\alpha$ particles. From the comparison of the results for the $^8\text{Be}$ and $^9\text{Be}$ systems, they found an interesting sign change of the quadrupole moments and the magnetic moments for some excited states. They also pointed out a possibility of negative parity resonances with $^8\text{He}+\alpha$ structure in this threshold region.

Our purpose for the $aa\Lambda$ Faddeev calculations using $aa\Lambda$ RGM kernels is threefold. First, we develop a general three-cluster Faddeev formalism with two identical clusters, in order to apply it to more complex three-cluster systems like the hypertriton interacting via quark-model baryon-baryon interactions. In the hypertriton system, we have to deal with the $\Lambda N\Sigma NN$ coupled-channel system which involves a Pauli-forbidden state at the quark level in the $\Lambda N\Sigma N$ subsystem. Since the baryon-baryon interactions in the quark model are formulated in the two-cluster RGM formalism, the present three-cluster formalism is most appropriate to correlate the baryon-baryon interactions with the structure of few-baryon systems. The second purpose is to make a consistent description of the $\alpha\alpha$, $3\alpha$ and $aa\Lambda$ systems using effective $NN$ and $\Lambda N$ interactions. This attempt is beyond the scope of the usual OCM framework and the Faddeev formalism assuming only inter-cluster potentials. A comparison of the present $3\alpha$ results with the fully microscopic $3\alpha$ RGM or GCM [38–40] is useful to examine the approximations involved in the present three-cluster formalism. The third purpose is to present a tractable effective $\Lambda N$ force for cluster calculations of various $p$-shell hypernuclei, which is not purely phenomenological but derived microscopically from quark-model baryon-baryon interactions. In particular, this effective $\Lambda N$ force should be able to reproduce the correct $\Lambda$-separation energy of $^9\Lambda\text{He}$; $B_{\Lambda}(^9\Lambda\text{He})=3.12\pm 0.02$ MeV. Such a $\Lambda\alpha$ interaction is indispensable for, e.g., a $\Lambda\Lambda\alpha$ Faddeev calculation using the quark-model $\Lambda\alpha$ interaction [41]. In this paper, we derive an effective $\Lambda N$ force of two-range Gaussian form from the phase-shift behavior of the quark-model $YN$ interaction, $f_{\Lambda\alpha}$, by using an inversion method based on supersymmetric quantum mechanics [42].

This paper is organized as follows. In the next section, the three-cluster Faddeev formalism with two identical clusters is given, together with expressions to calculate the expectation values of the two-cluster Hamiltonian with respect to the solutions resulting from the Faddeev equations. The procedure to calculate the $\Lambda\alpha$ and $aa\Lambda$ $T$-matrices is also discussed, as well as the treatment of the cut-off Coulomb force employed in this paper. In the third section, we first briefly discuss the results of the $3\alpha$ Faddeev calculation, and then those of the $aa\Lambda$ Faddeev calculation. The final section is devoted to a summary. Appendix A gives a brief comment on the rearrangement factors of three-body systems with two identical particles. The most general case with explicit spin-isospin degrees of freedom is discussed. In Appendix B, we derive a compact formula to calculate the $\Lambda\alpha$ Born kernel for arbitrary types of $\Lambda N$ interactions. Energies are in MeV and lengths in fm throughout, unless otherwise specified.

II. FORMULATION

A. Faddeev equation for systems with two identical clusters

In order to formulate the Faddeev equation for systems with two identical particles, we follow the notation of Refs. [43,44] as much as possible. The Jacobi-coordinate vectors are specified by the permutation $(\alpha\beta\gamma)$, which is a cyclic permutation of $(123)$. For example, the momentum vectors for the coordinate system $\gamma$ in the unit of $\hbar$ are defined by

$$p_\gamma = \frac{m_\alpha k_\alpha - m_\beta k_\beta}{m_\alpha + m_\beta},$$

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where \( k_\alpha, k_\beta, k_\gamma \) are single particle momenta of particles \( \alpha, \beta, \gamma \) with the masses \( m_\alpha, m_\beta, m_\gamma \) respectively, and \( M = m_\alpha + m_\beta + m_\gamma \) is the total mass. Three different sets of the Jacobi coordinates, \((p_1, q_1)\), \((p_2, q_2)\), and \((p_3, q_3)\), are related to each other in the standard relationship for the rearrangement. We choose the coordinate system \( \gamma = 3 \) as the standard set of Jacobi coordinates and assume that particles 1 and 2 are the two identical particles with a common mass \( m_1 = m_2 \). We incorporate the symmetry property for the exchange of particles 1 and 2 into the Faddeev formalism by assuming the total wave function \( \Psi(p, q) \) as

\[
\Psi(p_3, q_3) = \psi(p_3, q_3) \pm \varphi(-p_1, q_1) + \varphi(p_2, q_2)
\]

with \( \psi(p_3, q_3) = \pm \psi(p_3, q_3) \),

(2)

where the upper (lower) sign is applied for identical bosons (fermions). The requirement \( \Psi(-p, q) = \pm \Psi(p, q) \) is satisfied from this ansatz.

In the application to the \( \alpha \alpha \Lambda \) system, two \( \alpha \) clusters are numbered 1 and 2, and the \( \Lambda \) hyperon is numbered 3. Since the technique to handle the rearrangement of the Jacobi coordinates in the Faddeev formalism is well known [43], we only give the specification scheme of channels and the final Faddeev equation after partial-wave decomposition. We give expressions both in the Faddeev equation after partial-wave decomposition. We give expressions both in the Faddeev equation after partial-wave decomposition. We give expressions both in the Faddeev equation after partial-wave decomposition. We give expressions both in the Faddeev equation after partial-wave decomposition. We give expressions both in the Faddeev equation after partial-wave decomposition.

The partial-wave decomposed Faddeev equation for the two components \( \psi \) and \( \varphi \) in Eq. (2) is given by

\[
\psi_s(p, q) = \left[ E - \frac{\hbar^2}{4M \gamma} \left( p^2 + \frac{8 + \zeta}{4 \xi} q^2 \right) \right]^{-1} \int_0^{q_1^2} dq' \times \left( \sum_{\beta} \frac{1}{p_1} \hat{g}_{\beta s}(q, q'; x) \frac{1}{p_2} \varphi_s(p_2, q') \right),
\]

(5a)

\[
\varphi_s(p, q) = \left[ E - \frac{\hbar^2}{8M \gamma} \left( \frac{4 + \zeta}{\xi} p^2 + \frac{8 + \zeta}{4 + \xi} q^2 \right) \right]^{-1} \frac{1}{2} \int_0^{q_1^2} dq' \times \left( \sum_{\beta} \frac{1}{p_1} \hat{g}_{\beta s}(q, q'; x) \frac{1}{p_2} \psi_s(p_2, q') \right) + \left( \sum_{\beta} \frac{1}{p_1} \hat{g}_{\beta s}(q, q'; x) \frac{1}{p_2} \varphi_s(p_2, q') \right),
\]

(5b)

where \( \zeta = (M_\Lambda / M_\alpha) \) is the mass ratio of \( \Lambda \) to the nucleon and

\[
p_1 = p \left( q, \frac{1}{2} q', x \right), \quad p_2 = p \left( q, \frac{\zeta}{4 + \xi} q', x \right),
\]

\[
p_1' = p \left( q, \frac{1}{4 + \xi} q', x \right), \quad p_2' = p \left( q, \frac{1}{2} q', x \right),
\]

\[
\overline{p}_1 = p \left( q, \frac{4}{4 + \xi} q', x \right), \quad \overline{p}_2 = p \left( q, \frac{4}{4 + \xi} q', x \right),
\]

(5c)

with \( p(q, q'; x) = \sqrt{q^2 + q'^2 + 2qq'} \). The \( T \)-matrices, \( \hat{T}_\alpha \) and \( \hat{T}_{\beta} \), are discussed in Secs. II D and II C. The rearrangement factors for the \( \psi \) or \( \varphi \) is given by

\[
g_{\gamma \beta}(q, q'; x) = g_{\beta \gamma}(q', q; x) = \sum_{\lambda_1 \lambda_2} q^{\lambda_1 + \lambda_2} q'^{\lambda_1 + \lambda_2} \frac{1}{2} \left( \frac{\xi}{4 + \xi} \right)^{\lambda_1} \left( \frac{\zeta}{4 + \xi} \right)^{\lambda_2} \times \sum_k (2k + 1) g^{\lambda_1 \lambda_2 k} P_k(x),
\]

(6)

with \( P_k(x) \) is the Legendre polynomial of rank \( k \). The reduced rearrangement factor \( g^{\lambda_1 \lambda_2 k} \) is expressed as
with $j = \sqrt{2j+1}$, etc., and the spatial angular-momentum factor $G_{\lambda_1\lambda_2}(\ell_1,\ell_2)$ in Eq. (A9). For the $\varphi - \varphi$ type rearrangement, these factors are given by

$$g_{\beta\beta'}(q,q';x) = \sum_{\lambda_1+\lambda_2=\ell_1+\ell_2} \sum_{\lambda_1'\lambda_2'} q_{\lambda_1\lambda_2} q_{\lambda_1'\lambda_2'} \left( \frac{4}{4 + \zeta} \right)^{\lambda_1+\lambda_2} \sum_{k} (2k+1) g_{\lambda_1\lambda_2 k} P_k(x),$$

with

$$g_{\beta\beta'}(q,q';x) = \sum_{\lambda_1+\lambda_2=\ell_1+\ell_2} q_{\lambda_1\lambda_2} q_{\lambda_1'\lambda_2'} \left( \frac{4}{4 + \zeta} \right)^{\lambda_1+\lambda_2} \sum_{k} (2k+1) g_{\lambda_1\lambda_2 k} P_k(x).$$

B. Calculation of $e_\gamma$ and $e_\beta$

In this section, we derive some formulas to calculate expectation values of the two-cluster Hamiltonians, $h_\gamma + V_{RGM}(e_\gamma)$ and $h_\beta + V_\beta$, where $h_\gamma$ is the kinetic-energy operator of the $\gamma$ pair, etc. In the present application, $V_{RGM}(e_\gamma)$ is the $\alpha\alpha$ RGM kernel and $V_\beta$ is the $\Lambda\kappa$ kernel. We deal with the energy dependence of the $\alpha\alpha$ RGM kernel self-consistently by calculating

$$e_\gamma = \langle \Psi | h_\gamma + V_{RGM}(e_\gamma) | \Psi \rangle$$

for the normalized Faddeev solution $\Psi$. The potential term of the matrix element in Eq. (10) is most easily obtained from various matrix elements of the kinetic-energy operators. Suppose $\Psi$ is a sum of three Faddeev components, $\Psi = \psi_\alpha + \psi_\beta + \psi_\gamma$. Then the Faddeev equation $(E-H_0)\psi_\gamma = V_\gamma \psi_\gamma$ with $V_\gamma = V_{RGM}(e_\gamma)$ and $H_0 = h_\gamma + h_\gamma$ yields $\langle \Psi | V_\gamma | \Psi \rangle = \langle \psi_\gamma | E-H_0 | \psi_\gamma \rangle$. Thus Eq. (10) becomes

$$e_\gamma = E(\psi_\gamma | \Psi) - \langle \psi_\gamma | H_0 | \Psi \rangle + \langle \psi_\gamma | h_\gamma | \Psi \rangle.$$

We can write a similar equation also for the $\beta$ pair. We calculate $e_\beta$, although the self-consistent procedure is not necessary for the $\Lambda$ interaction. The kinetic energy term $\langle \Psi | h_\beta | \Psi \rangle$ is obtained from $\langle \Psi | h_\gamma | \Psi \rangle$ as follows. Using the momentum Jacobi coordinates in Eq. (1), we can easily show

$$\langle \psi_\beta | H_0 | \Psi \rangle = \sum_{\gamma} \int_{0}^{\infty} p^2 dq \left( \frac{p^2}{4M_N} + \frac{8 + \zeta}{4\xi} \right) \left[ \psi_\beta(p,q) \right]^2 \sum_{\gamma,\beta} \int_{0}^{\infty} q^2 dq \int_{-1}^{1} dx \left( \frac{\hbar^2}{4M_N} \right) \left( p^2 + \frac{8 + \zeta}{4\xi} \right) \frac{1}{p_1^{\lambda} g_{\beta\gamma}(q,q';x)} \frac{1}{p_2^{\gamma} \lambda_{\beta}(p_2,q')},$$

(16a)

Thus we find, for the $\alpha\alpha\Lambda$ system,

$$e_\beta = E(\psi_\beta | \Psi) + \frac{8 + \zeta}{2(4 + \zeta)} \langle \psi_\beta | H_0 | \Psi \rangle + \frac{4}{4 + \zeta} \langle \psi_\beta | H_0 | \Psi \rangle - \langle \psi_\beta | H_0 | \Psi \rangle.$$

(15)

We need to calculate the overlap matrix elements $\langle \psi_\beta | \Psi \rangle$, $\langle \psi_\beta | H_0 | \Psi \rangle$ and $\langle \psi_\beta | H_0 | \Psi \rangle$. For two identical particles with $m_\beta = m_\beta$, this relationship yields

$$h_\gamma + h_\beta = \frac{M}{m_\beta + m_\gamma} H_0 - \frac{2m_\beta}{m_\beta + m_\gamma} h_\gamma,$$

(13)

and

$$\langle \Psi | h_\beta | \Psi \rangle = \frac{M}{2(m_\beta + m_\gamma)} \langle \Psi | H_0 | \Psi \rangle - \frac{m_\beta}{m_\beta + m_\gamma} \langle \Psi | h_\gamma | \Psi \rangle.$$
\[
\langle \psi_{\beta}|H_0|\Psi \rangle = \sum_{\beta} \int_0^{\infty} p^2 dp q^2 dq \frac{\hbar^2}{8M_N} \left(4 + \frac{\xi}{\zeta} p^2 + \frac{8 + \xi}{4 + \zeta} \right) [\varphi_{\beta}(p, q)]^2 + \frac{1}{2} \sum_{\beta, \beta'} \int_0^{\infty} q^2 dq q'^2 dq' \int_{-1}^{1} dx \varphi_{\beta}(p_1, q) \frac{\hbar^2}{8M_N} \\
\times \left(4 + \frac{\xi}{\zeta} p_1^2 + \frac{8 + \xi}{4 + \zeta} q_1^2 \right) \frac{1}{p_1} \gamma_{\beta\beta'}(q, q'; x) \frac{1}{p_2} \varphi_{\beta'}(p_2, q') + \frac{1}{2} \sum_{\gamma, \beta} \int_0^{\infty} q^2 dq q'^2 dq' \int_{-1}^{1} dx \varphi_{\gamma}(p_1, q) \frac{\hbar^2}{4M_N} \\
\times \left(p_1^2 + \frac{8 + \xi}{4 + \zeta} q_1^2 \right) \frac{1}{p_1} \gamma_{\gamma\beta}(q, q'; x) \frac{1}{p_2} \varphi_{\beta}(p_2, q').
\]

The overlap integrals are obtained by setting \(H_0 \rightarrow 1\). Furthermore, \(\langle \Psi|h_{\gamma}|\Psi \rangle\) is given by

\[
\langle \Psi|h_{\gamma}|\Psi \rangle = \sum_{\gamma} \int_0^{\infty} p^2 dp q^2 dq \frac{\hbar^2}{4M_N} [\psi_{\gamma}(p, q)]^2 + \sum_{\gamma, \beta} \int_0^{\infty} q^2 dq q'^2 dq' \int_{-1}^{1} dx \psi_{\gamma}(p, q) \frac{\hbar^2}{2M_N} \frac{1}{p_1} \gamma_{\gamma\beta}(q, q'; x) \frac{1}{p_2} \varphi_{\beta}(p_2, q') \\
+ \sum_{\beta, \beta'} \int_0^{\infty} p^2 dp q^2 dq \varphi_{\beta}(p, q) \frac{\hbar^2}{8M_N} \left[p^2 + \left(\frac{8 + \xi}{4 + \zeta} \right) q^2 \right] \delta_{\beta, \beta'} + \frac{2(8 + \xi)}{4 + \zeta} p q f_{\beta\beta'}(p, q) \\
+ \sum_{\beta, \beta'} \int_0^{\infty} q^2 dq q'^2 dq' \int_{-1}^{1} dx \varphi_{\beta}(p_1, q) \frac{\hbar^2}{16M_N} (q^2 + q'^2 - 2qq'x) \frac{1}{p_1} \gamma_{\beta\beta'}(q, q'; x) \frac{1}{p_2} \varphi_{\beta'}(p_2, q').
\]

Here, \(f_{\beta\beta'}\) is given by

\[
f_{\beta\beta'} = \int d\hat{p}_2 d\hat{q}_2 (\hat{\beta}_2 \hat{\gamma}_2) (\hat{\beta} \hat{\gamma}) (\hat{\beta}_2 \hat{\gamma}_2) (\hat{\beta} \hat{\gamma}) = (-1)^{l_1+\ell_2} \ell_1 \ell_2 (\ell_1' 0 10) (\ell_2 010) (\ell_2' 010)
\]

\[
\times \left( \begin{array}{c}
-1 \ell_1 \\
\ell_2^l ' \ell_1' 1
\end{array} \right) \left( \begin{array}{c}
J L \\
\ell_1 I \ell_2 J
\end{array} \right) \left( \begin{array}{c}
J L \frac{1}{2} \\
\ell_1' I' \ell_2' \frac{1}{2}
\end{array} \right) \left( \begin{array}{c}
\ell_1 \ell_2 \frac{1}{L} \\
\ell_2^l \ell_1' 1
\end{array} \right) (LS-coupling).
\]

\[
C. \Lambda\alpha T\text{-matrix and effective } \Lambda N \text{ potentials}
\]

The \(\Lambda\alpha\) T-matrices are obtained by solving the Lippmann-Schwinger equation

\[
T_z(p, p'; E) = V_z(p, p') - \frac{4\pi}{(2\pi)^3} \frac{2\mu}{\hbar^2} \int_0^{\infty} k^2 dk V_z(p, k) \\
\times \frac{1}{\gamma^2 + k^2} T_z(k, p'; E),
\]

where \(\mu = (4\xi/4 + \zeta)M_N\) is the \(\Lambda\alpha\) reduced mass and \(E = -(h^2/2\mu)\gamma^2\) is a negative energy. The partial-wave components \(V_z(p, p')\) for the \(\Lambda\alpha\) Born kernel \(V(p, p')\) are defined through

\[
V(p, p') = 4\pi \sum_{\ell} V_z(p, p') \sum_{\ell} Y_{\ell m}(p)^* Y_{\ell m}(p'),
\]

and the \(\langle p|T_z(E)|p' \rangle\) in Eq. (5b) is related to \(T_z(p, p'; E)\) with an extra factor \(4\pi/(2\pi)^3\).

For the effective \(\Lambda N\) potential, we assume a Minnesota-type central force \([45]\)

\[
v_{\Lambda N} = \left[ v(1^E) \frac{1-P_\rho}{2} + v(3^E) \frac{1+P_\rho}{2} \right] \left[ \frac{u}{2} + \frac{2-u}{2} P_\rho \right],
\]

where \(v(1^E)\) and \(v(3^E)\) are simple two-range Gaussian potentials generated from the \(^1S_0\) and \(^3S_1\) phase shifts predicted by the quark-model \(\Lambda N\) interaction, fss2. We use the inversion method based on supersymmetric quantum mechanics, developed in Ref. [42], to derive phase-shift equivalent local potentials. These potentials are then fitted by two-range Gaussian functions. These are given by

\[
v(1^1S_0) = -128.0 \exp(-0.8908r^2) + 1015 \exp(-5.383r^2),
\]

\[
v(3^3S_1) = -56.31f \exp(-0.7517r^2) + 1072 \exp(-13.74r^2),
\]

where \(f = 1\) and \(r\) is the relative distance between \(\Lambda\) and \(N\). In the following, we call this effective \(\Lambda N\) potential the SB potential. Figure 1 shows that these potentials fit the low-energy behavior of the \(^1S_0\) and \(^3S_1\) \(\Lambda N\) phase shifts obtained by the full \(\Lambda N\) coupled-channel RGM calculations of fss2. In the \(^3S_1\) state, only the low-energy region is fitted, since the cusp region cannot be fitted in a single-channel
calculation. This potential overestimates the $^3S_1$ phase shift in the higher energy region. The procedure to calculate the \( \Lambda \alpha \) Born kernel for the simple \((0s)^4\) \( \alpha \)-cluster wave function is discussed in Appendix B. Here we only give the final result for the partial-wave components:

\[
V_i(q_f, q_i) = \sum_{i=1}^{4} \left[ X^i_d V_{e}^{i}(q_f, q_i; \kappa_i) + X^i_e V_{e}^{i}(q_f, q_i; \kappa_i) \right].
\] (23)

Here, \( X^i_d \) and \( X^i_e \) are spin-isospin factors defined in Eq. (B10) and tabulated in Table I for the present two-range Gaussian potentials. The explicit functional form of \( V_{e}^{i}(q_f, q_i; \kappa_i) \) and \( V_{e}^{i}(q_f, q_i; \kappa_i) \) are given in Eq. (B11).

In this paper, we also examine the \( \Lambda N \) effective forces [33] used by Hiyama et al. [7] for comparison. These potentials are generated from the \( G \)-matrix calculations of various OBEP potentials. They are parametrized as

<table>
<thead>
<tr>
<th>( i )</th>
<th>( X^i_d )</th>
<th>( X^i_e )</th>
<th>( \kappa_i )</th>
</tr>
</thead>
<tbody>
<tr>
<td>1, 2</td>
<td>((u/2)v_0(1S))</td>
<td>((1-u/2)v_0(1S))</td>
<td>((1-u/2)v_0(1S))</td>
</tr>
<tr>
<td>3, 4</td>
<td>((3u/2)v_0(3S))</td>
<td>((3(1-u/2)v_0(3S))</td>
<td>((1-u/2)v_0(1S))</td>
</tr>
</tbody>
</table>

**TABLE II.** Bound-state energies for the \( \Lambda \alpha \) system, \( E(\Lambda^\alpha(\alpha)) \) (in MeV), calculated by the original SB potential with \( f = 1 \). The h.o. width parameters, \( \nu = 0.275 \text{ fm}^{-2} \) and \( \nu = 0.257 \text{ fm}^{-2} \) are assumed for the \((0s)^4\) \( \alpha \)-cluster. The experimental value is \( E(\alpha) = -3.12 \pm 0.02 \text{ MeV} \).

<table>
<thead>
<tr>
<th>( \nu )</th>
<th>( \nu = 0.275 \text{ fm}^{-2} )</th>
<th>( \nu = 0.257 \text{ fm}^{-2} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>-4.975</td>
<td>-4.747</td>
</tr>
<tr>
<td>0.6</td>
<td>-4.946</td>
<td>-4.728</td>
</tr>
</tbody>
</table>

\[
u_{LNN} = \sum_{i=1}^{3} \left[ \left[ v_{0 \text{ even}}^{(i)} + v_{0 \text{ even}}^{(i)}(\sigma_1 \cdot \sigma_2) \right] \frac{1 + P_f}{2} 
\]

\[
+ \left[ v_{0 \text{ odd}}^{(i)} + v_{0 \text{ odd}}^{(i)}(\sigma_1 \cdot \sigma_2) \right] \frac{1 - P_f}{2} \right] e^{-i(\beta)^2}. \] (24)

Since the spin-spin term does not contribute to the spin-saturated \( \alpha \)-cluster, the spin-isospin factors in Eq. (23) (with \( \kappa_i \rightarrow 1/(\beta)^2 \), \( \beta = 1 - 3 \)) are given by

\[
X^i_d = 2(v_{0 \text{ even}}^{(i)} + v_{0 \text{ odd}}^{(i)}),
\]

\[
X^i_e = 2(v_{0 \text{ even}}^{(i)} - v_{0 \text{ odd}}^{(i)}). \] (25)
orthogonality condition. The $\Lambda\alpha$ phase shifts are also calculated, although there is no experimental information. The $S$-wave phase shift shows a monotonic decrease from 180° similar to Fig. 9 of Ref. [36]. In the energy region $E_{\text{cm}}(\Lambda\alpha)$=0–20 MeV, the phase shifts of the higher partial waves rapidly decrease, starting from 20°–30° for the $P$ wave. This implies that the $\Lambda\alpha$ potential is very much of the Wigner type, and our lack of knowledge of the $\Lambda\alpha$ interaction in higher partial waves may not become a serious problem in the Faddeev calculations.

D. $\alpha\alpha$ T-matrix and effective NN potentials

The $\alpha\alpha$ T-matrices used for the $3\alpha$ and $\alpha\alpha\Lambda$ Faddeev calculations are generated from the $\alpha\alpha$ RGM kernel which uses an effective NN potential similar to Eq. (21). In the notation used in Ref. [51], the $\alpha\alpha$ RGM kernel, $V_{\text{RGM}}(e)=V_D+V_D^{\text{Cl}}+G+G^{\text{Cl}}+\kappa K$ consists of the direct potential $V_D$, the direct Coulomb potential $V_D^{\text{Cl}}$, the sum of the exchange kinetic-energy and interaction kernels, $G=G^{K}+G^{V}$, the exchange Coulomb kernel $G^{\text{Cl}}$, and the exchange normalization kernel $K$. We have to eliminate redundant components from the energy-dependent partial wave $T$-matrices, $T_\lambda(p,p';E,e)$, which satisfy the Lippmann-Schwinger equation similar to Eq. (19). This is necessary only for the $S$-wave ($\lambda=0$) and $D$-wave ($\lambda=2$) components, for which there exist two and one h.o. Pauli redundant states, $u_{n\lambda}(p)$, respectively. Here, $u_{n\lambda}(p)$ are essentially the h.o. wave functions in the momentum representation with the total h.o. quanta $N=2n+\lambda=0$ and 2, satisfying $Ku_{n\lambda}=u_{n\lambda}$. They are explicitly given by

$$u_{n\lambda}(p)=(-1)^n\frac{(2\pi)^{3/2}}{\sqrt{4\pi}}R_{n\lambda}(p,\frac{1}{4\mu\nu})$$

with $\mu=2$, in terms of the standard three-dimensional h.o. radial wave function $R_{n\lambda}(r,\nu)$ in the coordinate representation. The RGM T-matrices defined in Ref. [14] are calculated by

$$\tilde{T}_\lambda(p,p';E,e)=T_\lambda(p,p';E,e)+\frac{\hbar^2}{4M_N}\frac{(\gamma^2+p'^2)(\gamma^2+p^2)}{(\gamma^2+\kappa^2)}$$

$$\times \left\{ \sum_{n=0}^{\infty} \frac{u_{n\lambda}(p)\overline{u}_{n\lambda}(p')}{u_{n\lambda}(p)u_{n\lambda}(p')} \right\}$$

for $\lambda = \begin{cases} 0 \\ 2 \end{cases}$, (27)

where $\gamma^2=-(4M_N E/\hbar^2)$ and $\kappa^2=(4M_N \kappa/\hbar^2)$. For the higher partial waves with $\lambda \geq 4$, we define $\overline{T}_\lambda(p,p';E,e)=\tilde{T}_\lambda(p,p';E,e)$. The RGM T-matrices in Eq. (27) satisfy the orthogonality condition

$$u_{n\lambda}(p) = \frac{4\pi}{(2\pi)^3} \frac{4M_N}{\hbar^2} \int_0^\infty p' \overline{P}_{\lambda}(p,p';E) \frac{u_{n\lambda}(p')}{\sqrt{p'^2+p'^2}}$$

for $n=0$, 1 ($\lambda=0$) and $n=0$ ($\lambda=2$). Owing to this relationship, we can prove the orthogonality of the total wave function Eq. (2) to the Pauli-forbidden states $u_{n\lambda}(p)$.

For the effective NN force, we mainly use the three-range Minnesota (MN) force [45] with the exchange-mixture parameter, $u=0.946$, 87, and the h.o. width parameter, $\nu=0.257$ fm$^{-2}$, for the $(0s)^4\alpha$-clusters. We also use the two-range Volkov No.1 (VN1) and No.2 (VN2) forces [52], in order to compare our $3\alpha$ results with the microscopic RGM [38,53] and GCM [39] calculations. The Majorana parameters $m$ of the Volkov forces and the h.o. width parameters are $m=0.575$ and $\nu=0.2515$ fm$^{-2}$ for VN1, and $m=0.59$ and $\nu=0.275$ fm$^{-2}$ for VN2. The $\alpha\alpha$ RGM calculations using these effective NN forces and the complete Coulomb kernel reasonably reproduce the empirical $\alpha\alpha$ phase shifts of the $S$, $D$, and $G$-waves, as well as the $S$-wave resonance near the $\alpha\alpha$ threshold. However, the best fit to the experiment is obtained by the three-range MN force. For the VN2 force, the $s$-wave resonance appears as a bound state with the binding energy $B_{\alpha\alpha}=245$ keV. Although the VN1 force reproduces this resonance, the overall fit to the $\alpha\alpha$ phase shifts is less impressive compared to the MN force. In the RGM calculation, the precise determination of the resonance energy is not easy even in the two-$\alpha$ system, because of the presence of the Coulomb force. In the present Lippmann-Schwinger formalism in the momentum representation, the method by Vincent and Phatak [54] is used for solving the scattering problem with full Coulomb force at the nucleon level. We find that the $0^+$ resonance energy is 0.18 and 0.14 MeV for VN1 force and the MN force, respectively. This should be compared with the experimental value 0.092 MeV.

For the Coulomb force in the $3\alpha$ and $\alpha\alpha\Lambda$ Faddeev calculations, we use the cut-off Coulomb force at the nucleon level

$$v_{ij}^\text{Cl}(r)=\frac{1}{2} \frac{\tau(i) \tau(i) e^2}{2} \theta(R_C-r),$$

with the cut-off radius $R_C$, although an exact treatment of the point Coulomb force exists for bound-state nuclear three-body problems with two charged particles [55]. Here $\theta(x)$ is the Heaviside step function. For the most compact $3\alpha$ ground state, this approximation with $R_C=10$ fm is good enough to obtain 1–2 keV accuracy. The exchange Coulomb kernel for Eq. (29) is calculated analytically. The partial-wave decomposition of the $\alpha\alpha$ RGM kernel is carried out numerically using the Gauss-Legendre 20-point quadrature formula, when the Coulomb force is not included. When the cut-off Coulomb force with a $R_C=14$ fm is employed, it is increased to the 30-point quadrature formula to obtain an accuracy within 1 keV for the exchange Coulomb kernel. The direct Coulomb term is separately integrated with a sufficient number of numerical integration points.
III. RESULTS

To solve the Faddeev equation, we discretize the continuous momentum variable \( p \) (or \( q \)) for the Jacobi coordinate vectors, using the Gauss-Legendre \( n_1 \)-point (\( n_2 \)-point) quadrature formula, for each of the three intervals of 0–1 fm\(^{-1} \), 1–3 fm\(^{-1} \) and 3–6 fm\(^{-1} \). The small contribution from the intermediate integral over \( p \) beyond \( p_0=6 \) fm\(^{-1} \) in the \( \alpha \alpha \) T-matrix calculation is also taken into account by using the Gauss-Legendre \( n_3 \)-point quadrature formula through the mapping \( p=p_0+\tan(\pi (1+x)/4) \). We need \( n_1 \geq 10 \) and \( n_3 = 5 \), so that 35 points are at least necessary to follow up the inner oscillations of the two-\( \alpha \) bound-state wave function and the necessary T-matrices for solving the Faddeev equation. These \( n_3 \) points for \( p > 6 \) fm\(^{-1} \) are, however, not included for solving the Faddeev equation, since it causes a numerical instability for the interpolation. The momentum range \( q = 6 \) fm\(^{-1} \) is also discretized by the \( n_3 \) point formula just as in the \( p \) discretization. We take \( n_1 = n_2 = n_3 = 15 \) for the \( 3 \alpha \) system and 10-10-5 for the \( \alpha \alpha \Delta \) system, respectively, unless otherwise specified. The modified spline interpolation technique developed in Ref. [56] is employed to generate the rearrangement matrices. For the large-scale diagonalization of non-symmetric matrices, the Arnoldi-Lanczos algorithm developed in the ARPACK subroutine package [57] is very useful.

A. \( 3 \alpha \) Faddeev calculation

In order to make sure that our Faddeev equation is solved correctly, we first carried out the standard \( 3 \alpha \)-particle Faddeev calculation by using the angular-momentum dependent Ali-Bodmer potential of \( d \) type (\( \text{ABd} \)). We find that the \( 3 \alpha \) energy, \( E_{3\alpha} = -6.423 \) MeV without Coulomb force, is consistent with previous calculations [9]. Here, we used \( \hbar^2/M_a = 10.4465 \) MeV fm\(^2 \) and \( e^2 = 1.44 \) MeV fm for comparison. When the cut-off Coulomb force is included, our value \(-1.527 \) MeV is 4 keV lower than the \(-1.523 \) MeV given in Table I of Ref. [9]. This difference is due to a slightly different treatment of the Coulomb force between the two calculations. The small \( 3 \alpha \) binding energy implies that the Ali-Bodmer phenomenological \( \alpha \alpha \) potential cannot describe the ground state of \(^{12}\text{C} \) with a compact shell-model like structure.

On the other hand, the present \( 3 \alpha \) model interacting via the \( \alpha \alpha \) RGM kernel gives enough binding and a large overlap with the compact shell-model-like component. Table III lists the results of such Faddeev calculations for the ground state of the \( 3 \alpha \) system with and without the Coulomb force. The \( \alpha \alpha \) RGM kernels are generated from the VN1, VN2, and MN forces. When the Coulomb effect is included, the cut-off Coulomb force with \( R_C = 10 \) fm is employed. In the last column in Table III, \( c_{(0)} \) implies the overlap amplitude of the \( 3 \alpha \) bound-state function with the \( SU_3 \) (04) shell-model configuration. We find that all three effective \( NN \) forces yield binding energies comparable with the experimental value \( |E_{3\alpha}^{\text{exp}}| = 7.275 \) MeV, although the result of of the \( MN \) force is a little too large. The dominant component of these \( 3 \alpha \) ground states is the \( SU_3 \) (04) shell-model configuration.

TABLE III. Results of \( 3 \alpha \) Faddeev calculations, using the \( \alpha \alpha \) RGM kernel, with and without the Coulomb effect. The parenthesisized numbers indicate the results when the cut-off Coulomb force with \( R_C=10 \) fm are included at the nucleon level. Partial waves up to \( \lambda_{\text{max}} \) are included in \( \alpha \alpha \) and (2\( \alpha \)-\( \alpha \)) channels. The heading \( E_{3\alpha} \) is the expectation value of the two-\( \alpha \) Hamiltonian with respect to the \( 3 \alpha \) bound state solution, \( E_{3\alpha} \), the \( 3 \alpha \) bound-state energy, and \( c_{(0)} \) the overlap between the \( 3 \alpha \) bound-state wave function and the \( SU_3 \) (04) shell-model configuration. For the MN force, the result of the variational calculation using the translationally invariant h.o. basis (h.o. var.) is also given for comparison, where h.o. quanta up to \( N=60 \) are included.

<table>
<thead>
<tr>
<th>Force</th>
<th>( \lambda_{\text{max}} )</th>
<th>( E_{2\alpha} )</th>
<th>( E_{3\alpha} )</th>
<th>( c_{(0)} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>VN1</td>
<td>4</td>
<td>9.657 (10.887)</td>
<td>-10.751 (-5.206)</td>
<td>0.900 (0.879)</td>
</tr>
<tr>
<td></td>
<td>8</td>
<td>9.530 (10.778)</td>
<td>-10.927 (-5.365)</td>
<td>0.896 (0.875)</td>
</tr>
<tr>
<td>VN2</td>
<td>4</td>
<td>8.853 (9.608)</td>
<td>-11.202 (-5.781)</td>
<td>0.826 (0.795)</td>
</tr>
<tr>
<td></td>
<td>8</td>
<td>8.449 (9.505)</td>
<td>-11.415 (-5.967)</td>
<td>0.821 (0.790)</td>
</tr>
<tr>
<td>MN</td>
<td>4</td>
<td>12.032 (13.603)</td>
<td>-15.616 (-9.433)</td>
<td>0.979 (0.973)</td>
</tr>
<tr>
<td></td>
<td>8</td>
<td>11.904 (13.481)</td>
<td>-15.779 (-9.591)</td>
<td>0.978 (0.971)</td>
</tr>
<tr>
<td>h.o. var.</td>
<td>11.903 (13.480)</td>
<td>-15.781 (-9.594)</td>
<td>0.978 (0.971)</td>
<td></td>
</tr>
</tbody>
</table>

In Table IV we compare the \( 3 \alpha \) ground-state energies \( E_{3\alpha} \) predicted in the present three-cluster Faddeev formalism, with those obtained by fully microscopic calculations, \( E_{3\alpha}^{\text{full}} \). We find that the present three-cluster equation gives \( 3 \alpha \) energies which are only 1.5–1.8 MeV higher than those of the fully microscopic \( 3 \alpha \) RGM or GCM calculations. This implies that the three-cluster exchange effect, which is neglected in our three-cluster formalism, but is present in the fully microscopic three-cluster RGM kernel, is attractive in nature, and is not as large as the repulsive three-body force claimed necessary in the semi-microscopic \( 3 \alpha \) models [2,4]. This is mainly because the \( 3 \alpha \) model space used by these authors does not exclude the \( 3 \alpha \) Pauli-forbidden components.

TABLE IV. Comparison of the \( 3 \alpha \) ground-state energies, predicted by the present model \( (E_{3\alpha}) \) and by fully microscopic calculations \( (E_{3\alpha}^{\text{full}}) \). The experimental value is \( E_{3\alpha}^{\text{exp}} = -7.275 \) MeV. The present model is the Faddeev calculation using the \( \alpha \alpha \) RGM kernel, including the cut-off Coulomb force with \( R_C=10 \) fm. The heading \( E_{\alpha} \) implies the internal energy of the \( (0s)^4 \) \( \alpha \)-cluster with the h.o. width parameter \( \nu \), \( E_{\text{tot}} \) the total energy from the RGM ([53] for MN and [38] for VN2) or GCM ([39] for VN1) calculations, and \( E_{3\alpha}^{\text{full}} = E_{\text{tot}} - 3E_{\alpha}^{\text{exp}} \).

<table>
<thead>
<tr>
<th>Force</th>
<th>( \nu ) (fm(^{-2} ))</th>
<th>( E_{\alpha}^{\text{full}} )</th>
<th>( E_{\text{tot}} )</th>
<th>( E_{3\alpha}^{\text{full}} )</th>
<th>( E_{3\alpha} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>VN1</td>
<td>0.257</td>
<td>-27.0</td>
<td>-87.9</td>
<td>-6.9</td>
<td>-5.37</td>
</tr>
<tr>
<td>VN2</td>
<td>0.275</td>
<td>-27.3</td>
<td>-89.4</td>
<td>-7.5</td>
<td>-5.97</td>
</tr>
<tr>
<td>MN</td>
<td>0.257</td>
<td>-23.9</td>
<td>-83.0</td>
<td>-11.4</td>
<td>-9.59</td>
</tr>
</tbody>
</table>
TABLE V. Kinetic- and potential-energy contributions to the three-α energy \( E_{3\alpha} \) calculated from \( \langle H_0 \rangle = 2(3e_{2\alpha} - E_{3\alpha}) \) and \( \langle V \rangle = 3(E_{3\alpha} - 2e_{2\alpha}) \). The shell-model (04) component, \( c_{(04)} \), is large if \( \langle H_0 \rangle \) is large.

<table>
<thead>
<tr>
<th>Force</th>
<th>( e_{2\alpha} )</th>
<th>( E_{3\alpha} )</th>
<th>( \langle H_0 \rangle )</th>
<th>( \langle V \rangle )</th>
<th>( c_{(04)} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>VN1</td>
<td>10.778</td>
<td>-5.366</td>
<td>75.402</td>
<td>-80.768</td>
<td>0.875</td>
</tr>
<tr>
<td>VN2</td>
<td>9.503</td>
<td>-5.969</td>
<td>68.958</td>
<td>-74.927</td>
<td>0.790</td>
</tr>
<tr>
<td>MN</td>
<td>13.481</td>
<td>-9.592</td>
<td>100.068</td>
<td>-109.660</td>
<td>0.971</td>
</tr>
</tbody>
</table>

accurately, unlike the one used in the present Faddeev formalism.

In Tables III and IV, we also find that the three-range MN force gives a somewhat large overbinding of 2–4 MeV, if the 3α energy \( E_{3\alpha} \) is measured from the 3α threshold. The decomposition of the 3α energy to the kinetic-energy and potential-energy contributions in Table V implies that this overbinding is due to the large cancellation between these two contributions. In this respect, it is interesting to note that the \( \alpha \) clusters with \( n \sim 0.257 \text{ fm}^{-2} \) (which gives the correct rms radius \( r_{\alpha} = (3/4) \rho = 1.48 \text{ fm} \) [58] for the simple (0s) \( \alpha \)-cluster) give less binding in the framework of the orthogonality condition model (OCM) [8]. If the h.o. constant parameter \( \nu \) is small, a proper treatment of the \( \alpha \alpha \) exchange kernel seems to be essential in order to obtain a large binding energy of the 3α ground state. This is reasonable since the large overlap of two \( \alpha \)-clusters implies the importance of nucleon exchange effects.

**B. \( \alpha\alpha\Lambda \) Faddeev calculation**

For a detailed description of the \( \alpha\alpha\Lambda \) bound states in the Faddeev calculation, it is important to make sure that the result is converged with respect to the following three conditions:
1. convergence with respect to the momentum discretization points,
2. convergence with respect to the extension of partial waves included,
3. convergence with respect to the cut-off radius \( R_C \), when the cut-off Coulomb force is included.

Among them, the Coulomb effect is the most difficult, since the \( T \)-matrix of the full Coulomb force is divergent at the diagonal part and the strong oscillation in the momentum representation in the cut-off Coulomb case does not lead to the correct answer, unless the numerical angular-momentum projection of the \( \alpha\alpha \) Coulomb kernel (especially the direct Coulomb term) is accurately performed. As to the partial waves, we can easily enumerate all possible angular-momentum states of \( \alpha^3 \text{Be} \) for the \( L^\pi = 0^+ \) ground state with \( J = 1/2 \) and the \( L^\pi = 2^+ \) excited state with \( J = 5/2 \) and 3/2 in the \( LS \) coupling scheme. If no \( \Lambda \alpha \) spin-orbit force is introduced, the \( J = 5/2 \) and 3/2 excited states are degenerate and the \( LS \)-coupling scheme is more efficient than the \( jj \)-coupling scheme to reduce the number of channels coupled in the calculation. In the following, the angular-momentum truncation is specified by \( \lambda_{\max} \ell_{\max} \) values for the \( \alpha\alpha \) and \( \Lambda\alpha \) pairs. For example, \( D-P \) in the ground-state calculation implies a 4-channel calculation and \( G-G \) in the \( L^\pi = 2^+ \) calculation a 19-channel calculation. The largest model space adopted is \( I-I \), which is an 11-channel calculation for \( L^\pi = 0^+ \) and a 28-channel calculation for \( L^\pi = 2^+ \). Note that the variational calculation in Ref. [7] uses a rather restricted model space, i.e., a 3-channel calculation with \( \lambda_{\max} = 2 \) and \( \ell_{\max} = 0 \), although the meaning of angular-momentum truncation is a little different from ours. For the momentum discretization points, we find that the energy change due to the increase of \( n_1 \), \( n_2 \), \( n_3 \) is very much \( R_C \) dependent. It is usually positive if we go from \( n_1 \), \( n_2 \), \( n_3 \) = 5-5 -5 to \( n_1 \), \( n_2 \), \( n_3 \) = 10-10-5 when the Coulomb force is not included, but it turns out negative when \( R_C = 10 \) and 14 fm. This implies that the Faddeev calculation without Coulomb force usually overestimates the binding energy, if the number of momentum discretization points is not large enough. Since the cut-off Coulomb kernels are oscillating, too small a number of momentum discretization points such as in \( n_1 \), \( n_2 \), \( n_3 \) = 5-5-5 case is dangerous when \( R_C \) is very large like \( R_C = 10 \) and 14 fm. The orthogonality to the Pauli-forbidden states also deteriorates when the number of momentum discretization points is too small. The squared norm of the Pauli-forbidden components contaminating the total wave function is typically \( 10^{-5} - 10^{-6} \) when \( n_1 \), \( n_2 \), \( n_3 \) = 5-5-5, but is improved to less than \( 10^{-13} \) for \( n_1 \), \( n_2 \), \( n_3 \) = 10-10-5. In this paper, we mainly show the results of \( n_1 \), \( n_2 \), \( n_3 \) = 10-10-5, since the energy gain by further extension to \( n_1 \), \( n_2 \), \( n_3 \) = 15-15-5 is usually less than 1 keV, when the cut-off Coulomb force with \( R_C = 10 - 14 \) fm is included.

The energy gain of the ground state, \( \Delta E \), and that of the self-consistent \( e_{2\alpha} \) value by the increase of the maximum angular-momentum values, \( \lambda_{\max} \ell_{\max} \), are shown in Table VI in the cases when we use the VN2 or MN forces for the \( \alpha\alpha \) interaction and the SB force for the \( \Lambda\alpha \) interaction. In these calculations the cut-off Coulomb force with \( R_C = 6 \) fm is employed. If the \( S \)-wave calculation is extended to include the \( D \)-wave, the energy gain is about 1 MeV for VN2+SB and 1.2 MeV for MN+SB. The energy gain mainly comes from the partial-wave component with \( \ell_1 = \ell_2 = 1 \) of the \( \alpha^3 \text{He} \) channel. The effect of the partial wave \( \ell_1 = \ell_2 = 2 \) is rather small; i.e., about 50 (VN2)–60 (MN) keV. Needless to say, the exact energy gain largely depends on the character of the \( \Lambda \) odd force. The ground-state energy is further improved by 7 (VN2)–5 (MN) keV and 0.03 (VN2)–0.0 (MN) keV, according to the extension to the \( G \)- and \( I \)-wave, respectively. On the other hand, \( e_{2\alpha} \) is improved by 165 (VN2)–288 (MN) keV, 6 keV and 0.5–0.6 keV, according to the extension to the \( D \)-, \( G \)- and \( I \)-waves, respectively. Thus, in conclusion, partial waves up to the \( D \)-wave are sufficient within 10 keV accuracy. If we wish to have a 1 keV accuracy, we need to take into account at least up to the \( G \)-wave. This implies that the partial-wave truncation in the Faddeev formalism is very efficient and the result converges very rapidly, according to the increase of the partial waves taken into account.

Table VII shows the \( R_C \) dependence of the two-α energy \( E(\alpha^3 \text{Be}) \), the self-consistently determined \( e_{2\alpha} \), the three-cluster ground-state energy \( E(\alpha^3 \text{Be}) \), the \( \Lambda \) separation energy defined by \( B_\Lambda(\alpha^3 \text{Be}) = E(\alpha^3 \text{Be}) - M_\Lambda - E(\alpha^3 \text{Be}) \), and the expecta-
tion value of the $\Delta \alpha$ Hamiltonian, $\epsilon_{\Delta \alpha}$, when the momentum discretization points with $n_1-n_2-n_3=10-10-5$ and the partial waves up to $I-I$ are used in the MN plus SB model. The energy increase (and the accumulated one) due to the increase of $R_C$ is also shown with the plus sign in the second (and the third) row. We find that the ground-state energy $E(\Lambda\alpha\beta)$ increases by 1.621 MeV when we move from $R_C=0$ to $R_C=6$ fm, which is larger than 1.127 MeV calculated for the free two-\alpha bound state. This seems to be natural, since the two-\alpha subsystem is more compact in the $\Lambda\alpha\beta$ system. The energy increase in the self-consistently determined $E_{2\alpha}$ values is 1.435 MeV, which is about 200 keV smaller than the energy increase in $E(\Lambda\alpha\beta)$, but is still larger than in the free two-\alpha bound state by about 300 keV. This observa-

TABLE VII. Cut-off radius ($R_C$) dependence of the Coulomb energies in the two-\alpha bound state energy $E(\Lambda\alpha\beta)$, the two-\alpha expectation value $E_{2\alpha}$, the three-body bound state energy $E(\Lambda\alpha\beta)$, the $\Lambda$ separation energy $B_{\Lambda}(\Lambda\alpha\beta)$, and the $\Lambda\alpha$ expectation value $E_{\Lambda\alpha}$. Calculations are carried out by using $n_1-n_2-n_3=10-10-5$ and the partial waves up to $I-I$. The three-range MN force and the SB force are used with $\nu=0.257$ fm$^{-2}$ for the h.o. width parameter of the $\alpha$-clusters. The energy increase (and the accumulated one) due to the increase of $R_C$ is also shown with the plus sign in the second (and third) row. The experimental $\Lambda$ separation energy is $B_{\Lambda\alpha}(\Lambda\alpha\beta)$=6.71\pm0.04 MeV. The suffix "ext" stands for extrapolation.

<table>
<thead>
<tr>
<th>$R_C$ (fm)</th>
<th>0</th>
<th>6</th>
<th>10</th>
<th>14</th>
<th>\infty</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E(\Lambda\alpha\beta)$</td>
<td>-1.260</td>
<td>-0.133</td>
<td>&gt;0</td>
<td>+1.127</td>
<td>(+1.260)</td>
</tr>
<tr>
<td>$E_{2\alpha}$</td>
<td>-0.384</td>
<td>1.051</td>
<td>1.180</td>
<td>1.181</td>
<td>(1.181)_{ext}</td>
</tr>
<tr>
<td>&amp;</td>
<td>+1.435</td>
<td>+0.129</td>
<td>+0.001</td>
<td>-</td>
<td>(-1.564)</td>
</tr>
<tr>
<td>&amp;</td>
<td>(+1.565)</td>
<td>(+1.565)</td>
<td>(+1.565)_{ext}</td>
<td>-</td>
<td>(-1.564)</td>
</tr>
<tr>
<td>$E(\Lambda\alpha\beta)$</td>
<td>-8.543</td>
<td>-6.922</td>
<td>-6.837</td>
<td>-6.837</td>
<td>(-6.837)_{ext}</td>
</tr>
<tr>
<td>&amp;</td>
<td>+1.621</td>
<td>+0.085</td>
<td>+0.000</td>
<td>-</td>
<td>(-1.706)</td>
</tr>
<tr>
<td>&amp;</td>
<td>(+1.706)</td>
<td>(+1.706)</td>
<td>(+1.706)_{ext}</td>
<td>-</td>
<td>(-1.706)</td>
</tr>
<tr>
<td>$B_{\Lambda\alpha}(\Lambda\alpha\beta)$</td>
<td>7.283</td>
<td>6.789</td>
<td>&gt;6.837</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>$E_{\Lambda\alpha}$</td>
<td>1.390</td>
<td>1.228</td>
<td>1.260</td>
<td>1.261</td>
<td>(+1.261)_{ext}</td>
</tr>
</tbody>
</table>

The $\Delta \alpha$ Hamiltonian, $\epsilon_{\Delta \alpha}$, when the momentum discretization points with $n_1-n_2-n_3=10-10-5$ and the partial waves up to $I-I$ are used in the MN plus SB model. The energy increase (and the accumulated one) due to the increase of $R_C$ is also shown with the plus sign in the second (and the third) row. We find that the ground-state energy $E(\Lambda\alpha\beta)$ increases by 1.621 MeV when we move from $R_C=0$ to $R_C=6$ fm, which is larger than 1.127 MeV calculated for the free two-\alpha bound state. This seems to be natural, since the two-\alpha subsystem is more compact in the $\Lambda\alpha\beta$ system. The energy increase in the self-consistently determined $E_{2\alpha}$ values is 1.435 MeV, which is about 200 keV smaller than the energy increase in $E(\Lambda\alpha\beta)$, but is still larger than in the free two-\alpha bound state by about 300 keV. This observa-

TABLE VI. Energy gain for the ground state ($\Delta E$) and that of the self-consistent $E_{2\alpha}$ value ($\Delta E_{2\alpha}$) in keV, for the extension of the maximum angular-momentum values, $\lambda_{\text{max}}$-$\ell_{\text{max}}$. The cut-off Coulomb force with $R_C=6$ fm is included.

<table>
<thead>
<tr>
<th>Force</th>
<th>VN2+SB</th>
<th>MN+SB</th>
</tr>
</thead>
<tbody>
<tr>
<td>$n_1-n_2-n_3$</td>
<td>5-5-5</td>
<td>10-10-5</td>
</tr>
<tr>
<td>$S-S\rightarrow D-P$</td>
<td>-954</td>
<td>-954</td>
</tr>
<tr>
<td>$D-P\rightarrow D-D$</td>
<td>-50</td>
<td>-50</td>
</tr>
<tr>
<td>$D-D\rightarrow G-G$</td>
<td>-7</td>
<td>-7</td>
</tr>
<tr>
<td>$G-G\rightarrow I-I$</td>
<td>-0.03</td>
<td>-0.03</td>
</tr>
</tbody>
</table>

Here we have shown the kinetic-energy and potential-energy contributions separately in each energy, and $c(40)$ is the overlap amplitude of the $\Lambda\alpha\beta$ ground-state wave function with the shell-model (40) wave function. [Note that the sum of the $E_{2\alpha}$ potential energy and twice of the $E_{\Lambda\alpha}$ potential energy is the potential energy of $E(\Lambda\alpha\beta)$, but this is not true for the kinetic-energy terms.] We have also carried out the similar analysis in the VN2+SB model. The converged result of the VN2+SB forces, including the cut-off Coulomb force with $R_C=14$ fm, is given by

$$E(\Lambda\alpha\beta) = 27.35 - 34.18 = -6.837 \text{ MeV},$$

$$E_{2\alpha} = 19.46 - 18.27 = 1.181 \text{ MeV},$$

$$E_{\Lambda\alpha} = 9.215 - 7.954 = 1.261 \text{ MeV},$$

$$c_{(40)} = 0.695.$$  

(30)
If we compare this result with Eq. (30) for the MN force, we find that the energy gain by the more attractive VN2 force is only 42 keV. This result is rather surprising, if we consider that the VN2 force gives a two-α bound state with energy $E_{2\alpha}=-245$ keV. The $\Lambda\alpha$ interaction by the SB force is also more attractive than in the MN force case due to the different choice of the h.o. width parameter $\nu$. In other words, the ground state energy of $^8\Lambda$Be is not much affected by the poor $\alpha\alpha$ and $\Lambda\alpha$ interactions, as long as we find a well-converged value by taking enough partial waves and a large number of momentum discretization points. On the other hand, the $E_{2\alpha}$ and $E_{\Lambda\alpha}$ values for the MN force are larger than those for the VN2 force by almost 500 keV. This may be related to the difference of $\nu$ values in the two calculations. The smaller $\nu$ value, 0.257 fm$^{-2}$, in the MN force calculation means more extended $\alpha$-clusters than in the VN2 calculation ($\nu$ =0.275 fm$^{-2}$), which implies in turn that the relative wave functions in the $2\alpha$ and $\Lambda\alpha$ subsystems should be more compact in the MN case. This can be confirmed by comparing the kinetic-energy contributions in $E^\alpha_{\Lambda}(\Lambda)$, $E_{2\alpha}$ and $E_{\Lambda\alpha}$, in Eqs. (30) and (31). For example, the kinetic-energy contribution in $E_{2\alpha}$ is 13.64 MeV in the VN2 case, while in the MN case it has a much larger value 19.46 MeV. The compactness of the $\alpha\Lambda\alpha$ relative wave function in the MN case is also reflected in the fact that $c_{(40)}$ is larger in the MN case, even though the binding energy is smaller. Comparing the result in Eq. (30) with the experimental value $E^\text{exp}_{\Lambda}(\Lambda)$ = -6.62±0.04 MeV, we can conclude that the MN+SB combination overrides the $^8\Lambda$Be ground-state energy by 220 keV. This is partly because our SB potential is of the pure Serber type ($\nu=1$). If we choose $\nu=0.82$ for the SB force, the combination with the present MN force and $\nu=0.257$ fm$^{-2}$ yields $E^\alpha_{\Lambda}(\Lambda)$=−6.621 MeV. In this case, the $^8\Lambda$Be bound-state energy is −3.105 MeV.

We list the results of various $\Lambda\alpha$ effective forces used by Hiyama et al. in Table VIII, when they are used in combination with the MN force for the $\alpha\alpha$ RGM kernel. The calculations are carried out with $n_1-n_2-n_3=10-10-5$, $R_C=10$ fm, and the partial waves up to the $G$-wave, to obtain the converged results with the accuracy of 1 – 2 keV.

Table IX lists $\alpha\Lambda\alpha$ Faddeev calculations for the $2^+$ excited state, including the cut-off Coulomb force with $R_C=10$ fm. The momentum discretization points with $n_1-n_2-n_3=10-10-5$ are employed. When the partial waves are restricted to D-S or S-D, the $2^+$-state energy is located above the $\alpha$-$^3\Lambda$He threshold with the threshold energy −3.12 MeV. The listing therefore starts from the $7$-channel calculation with D-P. We find that the result is almost converged with $I-I$ and $R_C=14$ fm, within the accuracy of 1 keV. The final result for the $2^+$ excited state in the MN+SB model is

$$E^{(\alpha\Lambda\alpha)(\Lambda)}_{\Lambda}(\Lambda) = 21.21 - 28.09 = -6.879 \text{ MeV},$$
$$E^{(\alpha\alpha\alpha)(\Lambda)}_{\alpha\alpha\alpha} = 13.64 - 12.99 = 0.649 \text{ MeV},$$
$$E^{(\alpha\alpha\alpha\alpha)(\Lambda)}_{\alpha\alpha\alpha\alpha} = 8.264 - 7.548 = 0.715 \text{ MeV},$$
$$c_{(40)} = 0.569. \quad (31)$$

If we compare Eq. (30) for the MN force, we find that the energy gain by the more attractive VN2 force is only 42 keV. This result is rather surprising, if we consider that the VN2 force gives a two-$\alpha$ bound state with energy $E_{2\alpha}=-245$ keV. The $\Lambda\alpha$ interaction by the SB force is also more attractive than in the MN force case due to the different choice of the h.o. width parameter $\nu$. In other words, the ground state energy of $^8\Lambda$Be is not much affected by the poor $\alpha\alpha$ and $\Lambda\alpha$ interactions, as long as we find a well-converged value by taking enough partial waves and a large number of momentum discretization points. On the other hand, the $E_{2\alpha}$ and $E_{\Lambda\alpha}$ values for the MN force are larger than those for the VN2 force by almost 500 keV. This may be related to the difference of $\nu$ values in the two calculations. The smaller $\nu$ value, 0.257 fm$^{-2}$, in the MN force calculation means more extended $\alpha$-clusters than in the VN2 calculation ($\nu$ =0.275 fm$^{-2}$), which implies in turn that the relative wave functions in the $2\alpha$ and $\Lambda\alpha$ subsystems should be more compact in the MN case. This can be confirmed by comparing the kinetic-energy contributions in $E^\alpha_{\Lambda}(\Lambda)$, $E_{2\alpha}$ and $E_{\Lambda\alpha}$, in Eqs. (30) and (31). For example, the kinetic-energy contribution in $E_{2\alpha}$ is 13.64 MeV in the VN2 case, while in the MN case it has a much larger value 19.46 MeV. The compactness of the $\alpha\alpha\Lambda$ relative wave function in the MN case is also reflected in the fact that $c_{(40)}$ is larger in the MN case, even though the binding energy is smaller. Comparing the result in Eq. (30) with the experimental value $E^\text{exp}_{\Lambda}(\Lambda)$ = -6.62±0.04 MeV, we can conclude that the MN+SB combination overrides the $^8\Lambda$Be ground-state energy by 220 keV. This is partly because our SB potential is of the pure Serber type ($\nu=1$). If we choose $\nu=0.82$ for the SB force, the combination with the present MN force and $\nu=0.257$ fm$^{-2}$ yields $E^\alpha_{\Lambda}(\Lambda)$=−6.621 MeV. In this case, the $^8\Lambda$Be bound-state energy is −3.105 MeV.

We list the results of various $\Lambda\alpha$ effective forces used by Hiyama et al. in Table VIII, when they are used in combination with the MN force for the $\alpha\alpha$ RGM kernel. The calculations are carried out with $n_1-n_2-n_3=10-10-5$, $R_C=10$ fm, and the partial waves up to the $G$-wave, to obtain the converged results with the accuracy of 1 – 2 keV.

Table IX lists $\alpha\Lambda\alpha$ Faddeev calculations for the $2^+$ excited state, including the cut-off Coulomb force with $R_C=10$ fm. The momentum discretization points with $n_1-n_2-n_3=10-10-5$ are employed. When the partial waves are restricted to D-S or S-D, the $2^+$-state energy is located above the $\alpha$-$^3\Lambda$He threshold with the threshold energy −3.12 MeV. The listing therefore starts from the $7$-channel calculation with D-P. We find that the result is almost converged with $I-I$ and $R_C=14$ fm, within the accuracy of 1 keV. The final result for the $2^+$ excited state in the MN+SB model is

$$E_{2\alpha} = 21.55 - 17.54 = 4.013 \text{ MeV},$$
$$E_{\Lambda\alpha} = 9.481 - 7.930 = 1.551 \text{ MeV},$$
$$c_{(40)} = 0.645. \quad (32)$$

If we compare Eqs. (30) and (32), we find that the 3 MeV excitation energy of the $2^+$ state mainly comes from an increase of the two-$\alpha$ kinetic energy (2 MeV) and from the two-$\alpha$ potential energy (1 MeV). This clearly shows the rotational nature of the ground $0^+$ and excited $2^+$ states, com-
TABLE IX. Same as Table VIII, but for the $L^s=2^+$ excited state with $R_C=14$ fm.

<table>
<thead>
<tr>
<th>Force</th>
<th>$\lambda_{\text{max}}$</th>
<th>$\ell_{\text{max}}$</th>
<th>$E_\Lambda(\Lambda^\Lambda\text{Be})$</th>
<th>$E_{2\alpha}$</th>
<th>$E_{\Lambda\Lambda}$</th>
<th>$c_{(4)}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>SB</td>
<td>$D-P$</td>
<td>$-3.797$</td>
<td>$3.987$</td>
<td>$1.528$</td>
<td>$0.643$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$D-D$</td>
<td>$-3.874$</td>
<td>$4.014$</td>
<td>$1.536$</td>
<td>$0.645$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$G-G$</td>
<td>$-3.926$</td>
<td>$4.013$</td>
<td>$1.550$</td>
<td>$0.645$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
<td>$-3.926$</td>
<td>$4.013$</td>
<td>$1.551$</td>
<td>$0.645$</td>
<td></td>
</tr>
<tr>
<td>NS</td>
<td>$D-P$</td>
<td>$-3.700$</td>
<td>$3.920$</td>
<td>$1.518$</td>
<td>$0.639$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$D-D$</td>
<td>$-3.772$</td>
<td>$3.946$</td>
<td>$1.525$</td>
<td>$0.641$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$G-G$</td>
<td>$-3.831$</td>
<td>$3.942$</td>
<td>$1.544$</td>
<td>$0.641$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
<td>$-3.831$</td>
<td>$3.943$</td>
<td>$1.544$</td>
<td>$0.641$</td>
<td></td>
</tr>
<tr>
<td>ND</td>
<td>$D-P$</td>
<td>$-4.377$</td>
<td>$4.027$</td>
<td>$1.130$</td>
<td>$0.648$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$D-D$</td>
<td>$-4.518$</td>
<td>$4.071$</td>
<td>$1.134$</td>
<td>$0.651$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$G-G$</td>
<td>$-4.553$</td>
<td>$4.066$</td>
<td>$1.137$</td>
<td>$0.651$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
<td>$-4.553$</td>
<td>$4.067$</td>
<td>$1.138$</td>
<td>$0.651$</td>
<td></td>
</tr>
<tr>
<td>NF</td>
<td>$D-P$</td>
<td>$-3.853$</td>
<td>$3.825$</td>
<td>$1.223$</td>
<td>$0.637$</td>
<td></td>
</tr>
<tr>
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<td>$D-D$</td>
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<tr>
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<td>$G-G$</td>
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<td>$1.236$</td>
<td>$0.639$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
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<td>$3.849$</td>
<td>$1.236$</td>
<td>$0.639$</td>
<td></td>
</tr>
<tr>
<td>JA</td>
<td>$D-P$</td>
<td>$-3.645$</td>
<td>$3.805$</td>
<td>$1.380$</td>
<td>$0.635$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$D-D$</td>
<td>$-3.710$</td>
<td>$3.827$</td>
<td>$1.385$</td>
<td>$0.637$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$G-G$</td>
<td>$-3.762$</td>
<td>$3.825$</td>
<td>$1.401$</td>
<td>$0.637$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
<td>$-3.762$</td>
<td>$3.826$</td>
<td>$1.402$</td>
<td>$0.637$</td>
<td></td>
</tr>
<tr>
<td>JB</td>
<td>$D-P$</td>
<td>$-3.460$</td>
<td>$3.775$</td>
<td>$1.507$</td>
<td>$0.632$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$D-D$</td>
<td>$-3.510$</td>
<td>$3.792$</td>
<td>$1.511$</td>
<td>$0.633$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$G-G$</td>
<td>$-3.568$</td>
<td>$3.793$</td>
<td>$1.535$</td>
<td>$0.634$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$I-I$</td>
<td>$-3.568$</td>
<td>$3.794$</td>
<td>$1.535$</td>
<td>$0.634$</td>
<td></td>
</tr>
</tbody>
</table>

The exchange term of the $\Lambda\alpha$ folding potential may also contribute to this difference.

If we arrange the effective $\Lambda N$ forces in Table X in the order of more attractive nature, we find

$$
\text{ND}(-7.483) > \text{NF}(-6.906) > \text{SB}(-6.837) > \text{NS}(-6.742) > \text{JA}(-6.677) > \text{JB}(-6.474). \tag{33}
$$

The experimental value $-6.62\pm 0.04$ MeV is located between JA and JB. However, this does not mean that the Jülich potentials JA and JB are the most correct $\Lambda N$ interactions. It is well known that the spin-spin central terms of these Jülich potentials are completely wrong and that they fail to reproduce the observed energy spectrum of the $^3\text{H}$ and $^4\text{He}$ systems [60]. As for the $2^+$ excitation energy, all the results in Table X are between 2.91 and 2.93 MeV. They are too small by 110–130 keV with respect to the average value 3.04 MeV of the two resonances recently observed by $\gamma$-ray spectroscopy [30,31]. Since the experimental error bars are at most $\pm 40$ keV even in the $(K,\pi)$ reaction [26], this is a meaningful disagreement. It would be interesting to examine the $\ell s$ splitting of the $5/2^+-3/2^+$ states, by introducing a small $\Lambda N$ spin-orbit force predicted by our quark-model interaction.

In order to show that the present $\alpha\alpha$ RGM kernel gives a better result than simple $\alpha\alpha$ potentials, we show in Table XI some results of $\alpha\alpha\Lambda$ Faddeev calculations using the Ali-Bodmer potential, ABd [12], and the Buck, Friedrich, and Wheatley potential, BFW [11]. In these cases, there needs to be no self-consistent procedure to determine $\varepsilon_{2\alpha}$. We only use the SB potential for the $\Lambda\alpha$ interaction, since results with other effective $\Lambda N$ forces are easily evaluated from the above discussion in the case of the $\alpha\alpha$ RGM kernel. In these $\alpha$-particle models, we customarily use $\hbar^2/M_\alpha = 10.4465$ MeV fm$^2$ and $e^2 = 1.44$ MeV fm. The momentum discretization points with $n_1-n_2-n_3=15-10-5$ are employed.
TABLE XI. $\alpha\alpha\Lambda$ Faddeev calculations for the $L^z=0^+$ ground state by the Ali-Bodmer (ABd) [12] and Buck, Friedrich, and Wheatley (BFW) [11] $\alpha\alpha$ potentials. The SB $\Lambda\Lambda$ force is used for the $\Lambda\alpha$ interaction. The cut-off Coulomb force is included at the nucleon level with $R_C=10$ fm. The h.o. width parameters of the $\alpha$-clusters are assumed to be $\nu=0.271$ 27 fm$^{-2}$ (ABd) and $\nu=0.257$ fm$^{-2}$ (BFW). The parameters $\hbar^2/M_a=10.4465$ MeV fm$^2$ and $\varepsilon^2=1.44$ MeV fm are used. Partial waves up to $\ell_{\max}$ in the $\alpha\alpha$ channel and those up to $\ell_{\max}$ in the $\Lambda\alpha$ channel. The momentum discretization points with $n_1,n_2,n_3=15,10,5$ are employed. The $\Lambda\alpha$ bound-state energy $E(\Lambda\alpha)$ is given in the first column.

<table>
<thead>
<tr>
<th></th>
<th>$\ell_{\max}$</th>
<th>$E(\Lambda\alpha)$</th>
<th>$E(\alpha\alpha \Lambda)$</th>
<th>$c(40)$</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>$\varepsilon_{2\alpha}$</td>
<td>$\varepsilon_{\Lambda\alpha}$</td>
<td></td>
</tr>
<tr>
<td>ABd+SB</td>
<td>-3.183</td>
<td>-6.409</td>
<td>0.970</td>
<td>-0.503</td>
</tr>
<tr>
<td></td>
<td>-3.066</td>
<td>-5.544</td>
<td>0.861</td>
<td>1.776</td>
</tr>
<tr>
<td>BFW+SB</td>
<td>S-S</td>
<td>-7.091</td>
<td>1.013</td>
<td>-0.532</td>
</tr>
<tr>
<td></td>
<td>D-P</td>
<td>-7.147</td>
<td>1.013</td>
<td>-0.526</td>
</tr>
<tr>
<td></td>
<td>-3.183</td>
<td>-7.153</td>
<td>1.018</td>
<td>-0.518</td>
</tr>
<tr>
<td></td>
<td>-3.066</td>
<td>-7.038</td>
<td>1.155</td>
<td>1.979</td>
</tr>
<tr>
<td></td>
<td>D-D</td>
<td>-7.153</td>
<td>1.018</td>
<td>-0.517</td>
</tr>
<tr>
<td></td>
<td>G-G</td>
<td>-7.153</td>
<td>1.018</td>
<td>-0.517</td>
</tr>
<tr>
<td></td>
<td>I-I</td>
<td>-7.043</td>
<td>1.161</td>
<td>1.980</td>
</tr>
</tbody>
</table>

For the $\alpha\alpha$ Coulomb potential, the folding potential of the cut-off Coulomb force with the (0s)$^3$ shell-model wave function is used with $R_C=10$ fm. The h.o. width parameter of the (0s)$^4$ $\alpha$-cluster for this Gaussian folding is $\nu=0.271$ 27 fm$^{-2}$ in the ABd case and $\nu=0.257$ fm$^{-2}$ in the BFW case. In the ABd case, this $\nu$ value corresponds to the Coulomb-force parameter $\beta=(\sqrt{3}/(2\times1.44))=0.6014$ fm$^{-1}$ and the $\alpha$ rms radius, $r_\alpha=(3/4\sqrt{\nu})=1.44$ fm. Since this $\nu$ value is also used for the $\alpha$-cluster folding for the $\Lambda\alpha$ potential, the $\Lambda\alpha$ bound-state energy $E(\Lambda\alpha)$ is a little shifted from the fitted experimental value $-3.12$ MeV. [The different $\hbar^2/M_N$ value also affects this difference.] Since the energy change is only about 0.06 MeV, we do not readjust the potential parameters of the $\Lambda\Lambda$ force. In the BFW case, the $\nu$ value, 0.257 fm$^{-2}$, corresponds to $\beta=(4\nu/\sqrt{3})=0.5853$ 38 fm$^{-1}$ and the rms radius of the $\alpha$-cluster, $r_\alpha=(\sqrt{3}/(2\beta)=1.48$ fm. In this case the difference of the $\Lambda\alpha$ bound-state energy, 0.054 MeV, from 0.32 MeV is solely from the different $\hbar^2/M_N$ value. The bound-state solutions of the BFW potential are used for the pairwise Pauli-forbidden states. The elimination of the Pauli-forbidden components from the three-body total wave function is always inspected by calculating their squared norm, which is of the order of $10^{-13}$.

We find that the $\alpha\alpha\Lambda$ ground-state energy by the ABd potential is lower than the result of the MN force in Eq. (30) by 0.3 MeV. Note that even in this case the energy gain from the higher partial waves than the $S$ wave is appreciable, i.e., 0.7 MeV. This implies that the $S$-wave assumption adopted by Filikhin and Gal [34] is not valid. They used a little different version of the Ali-Bodmer potential (type (a) with 125 MeV modified by 120 MeV) and obtained $E(\alpha\alpha \Lambda)=\approx-6.55$ MeV in the $S$-wave approximation. We expect an energy gain of about 0.7 MeV from the higher partial waves and their result is overbound, in comparison with the experimental value, $-6.62\pm0.04$ MeV. In Table XI, we find that the BFW potential gives a better result than the Ali-Bodmer force, but the energy is still lower than in the MN force case by 0.2 MeV. In this case we find that the effect of partial waves higher than the $S$ wave is quite appreciable, i.e., $-1.5$ MeV. This is of course due to the inner oscillation of the relative wave function between the two $\alpha$-clusters in the $\alpha\alpha\Lambda$ ground state. The shell-model like (40) components are about 0.7 in amplitude, which is appreciably larger than $c(40)\approx 0.5$ in the Ali-Bodmer case.

IV. SUMMARY

The three-cluster Faddeev formalism using two-cluster resonating-group method (RGM) kernels opens a way to solve few-baryon systems interacting via quark-model baryon-baryon interactions without spoiling essential features of the RGM kernel, i.e., the non-locality, the energy dependence proportional to the exchange normalization kernel, and the existence of pairwise Pauli-forbidden states in some specific channels. In this paper, we have applied this formalism to three-cluster systems involving $\alpha$-clusters, i.e., the $3\alpha$ and $\alpha\alpha\Lambda$ systems. These systems involve all of the above three features for the microscopic interactions between composite particles. In particular, the $\alpha\alpha$ interaction is a prototype of composite-particle interactions, in which the fully microscopic RGM calculation is easy and very successful. It, however, involves a somewhat complex kernel structure composed of three non-trivial Pauli-forbidden states, and the energy-dependence of the interaction is rather strong in the Pauli-allowed model space. In the present Faddeev formulation, the Pauli-forbidden components between pairwise clusters are completely eliminated from the total wave function of the three clusters. This can be achieved by introducing a special type of RGM $\widetilde{T}$-matrix calculated from the two-cluster RGM kernel, which satisfies the $T$-matrix version of the orthogonality conditions to the relative motion between two clusters. The on-shell and half off-shell properties of the $\widetilde{T}$-matrix are just the same as those of the ordinary $T$-matrix. This RGM $\widetilde{T}$-matrix involves a relative energy of two clusters as a parameter, which is determined self-consistently by calculating the expectation value of the two-cluster Hamiltonian with respect to the total wave function resulting from the Faddeev equation. The Faddeev equation using $\widetilde{T}$-matrices is equivalent to the pairwise orthogonality condition model (OCM) of three-cluster systems, interacting via two-cluster RGM kernels. A nice point of this formalism is that the underlying nucleon-nucleon (NN) and hyperon-nucleon (YN) interactions are more directly related to the structure of three-cluster systems than in the models assuming simple two-cluster potentials.
We have first applied the present formalism to the ground state of the $3\alpha$ system by using three different types of effective $NN$ forces, the two-range Volkov forces, No.1 (V1N) and No. 2 (V2N), and the three-range Minnesota (MN) force. The three-range MN force reproduces the $S$-, $D$- and $G$-wave $\alpha\alpha$ phase shifts quite well in the simple $(0s)^4$-model of the $\alpha$ clusters. The comparison with the $3\alpha$ RGM calculation has shown that the present three-cluster formalism using only the $\alpha\alpha$ RGM kernel gives a good approximation to the microscopic $3\alpha$ model. The difference of the ground-state energies predicted by these two models is less than 2 MeV. The effect of the antisymmetrization among three $\alpha$-clusters, which is neglected in our formulation, is attractive and is not so large, as long as the Pauli-allowed model space of the $3\alpha$ system is properly treated. It is also shown that the three-range MN force gives a lower ground-state energy than the two-range V1N and V2N forces, resulting in a somewhat large overbinding of 2–4 MeV, if the $3\alpha$ ground-state energy is measured from the $3\alpha$ threshold.

The application to the $^9\text{Be}$ system has proved that our three-cluster formalism is soundly extended to the systems with two identical clusters, in addition to the systems of three identical clusters like the $3\alpha$ system and the triton system. Here we have introduced a new effective $\Lambda N$ force, called the SB force, which is made from the quark-model predictions of the $\Lambda N$ phase shifts by using an inversion method based on supersymmetric quantum mechanics [42]. The SB force consists of two simple two-range Gaussian potentials which reproduce the low-energy behavior of the $^3S_0$ and $^3S_1$ $\Lambda N$ phase shifts predicted by $\Lambda N-\Sigma N$ coupled-channel RGM calculations using the model fss2 [46]. Since any central and single-channel effective $\Lambda N$ force leads to the well-known overbinding problem of $^5\text{He}$ by about 2 MeV [50], the attractive part of the $^3S_1$ $\Lambda N$ potential is reduced by about 10% to reproduce the empirical $\Lambda$-separation energy. $B^{\exp}_{\Lambda\text{Be}}(\Lambda\text{He}) = 3.12\pm 0.02$ MeV. The odd-state $\Lambda N$ force is assumed to be zero (pure Serber type). In addition to this SB force, we have also used the effective $\Lambda N$ forces in Ref. [7] for comparison. The $\Lambda\alpha$ interactions are generated from these $\Lambda N$ effective forces by the folding procedure with respect to the $(0s)^4$ h.o. wave function of the $\alpha$ clusters.

In the $\alpha\alpha\Lambda$ Faddeev calculation, sufficient partial waves up to $\lambda_{\text{Max}} = \ell_{1\text{Max}} = 6$ are included both in the $\alpha\alpha$ and $\Lambda\alpha$ pairs since the relative wave functions between two $\alpha$-clusters are oscillating at least in the relative $S$- and $D$-waves. The detailed analysis shows that the partial waves up to the $D$-wave are sufficient if we do not mind a 10 keV inaccuracy. If we wish to obtain a 1 keV accuracy, we need to take into account at least up to the $G$-wave. This implies that the partial wave truncation is very efficient even in the present Faddeev formalism. The energy gain due to partial waves higher than the $S$-wave is about 1 MeV for the V1N force and 1.2 MeV for the MN force, when these $\alpha\alpha$ interactions are used in combination with the SB force for the $\Lambda\alpha$ interaction. The Coulomb effect between the two $\alpha$-clusters is included by a cut-off Coulomb force at the nucleon level. The cut-off radius, $R_c = 10–14$ fm seems to be sufficient for a 1–2 keV accuracy. In the present formalism, the structure change of two $\alpha$-clusters inside $^9\text{Be}$ is clearly identified by calculating the kinetic-energy contribution in the two-$\alpha$ expectation value $\varepsilon_{2\alpha}$. The comparison of the Coulomb contributions in the $\alpha\alpha$ bound state, $\varepsilon_{2\alpha}$, and the $^9\text{Be}$ ground state with respect to the change of $R_c$ is very useful to measure the compactness of the two-$\alpha$ configurations in various environments. It is confirmed that the $0^+$ ground state and the $2^+$ exited state of $^9\text{Be}$ are well described by the contracted two-$\alpha$ cluster structure with a weakly coupled $\Lambda$-particle in the dominant $S$-wave component. In the present calculation using only central forces, the three-range MN force and the SB potential with the pure-Serber character can reproduce the ground-state and excitation energies of $^9\text{Be}$ within an accuracy of 100–200 keV. The results in Ref. [7] based on the OCM framework are also confirmed within 100 keV accuracy. On the other hand, the simple $\alpha$-particle model using the Ali-Bodmer $\alpha\alpha$ potential, Abd [12], and the OCM using the deep Buck, Friedrich, and Wheatley $\alpha\alpha$ potential, BFW [11], with bound-state Pauli-forbidden states give an overbinding of the $^9\text{Be}$ ground state by 530 and 420 keV, respectively, when the SB force is used for the $\Lambda\alpha$ interaction. Although these energies are rather similar, the effect of partial waves higher than the $S$-wave is very different, i.e., 0.7 MeV in the Ali-Bodmer case and 1.5 MeV in the BFW case. It is natural that the $\alpha\alpha$ interactions which yield an oscillatory behavior of the $\alpha\alpha$ relative wave functions, like our RGM kernel and the BFW potential, need more partial waves with a larger energy gain.

There are still many problems left for future studies. First of all, the readjustment of the $^3S$ attractive part of the SB $\Lambda N$ potential is unsatisfactory from the viewpoint of using the fundamental baryon-baryon interactions. The Brueckner rearrangement effect in $^4\text{He}$ is fairly large even for the rather stable $\alpha$-cluster [50]. In this sense, there is still no consistent description of the $s$-shell and $p$-shell hypernuclei even at the level of using effective baryon-baryon interactions. A microscopic description of the $\Lambda\alpha$ interaction may need more detailed analysis based on the $G$-matrix theory, for which the folding formula given in Appendix B is very useful. In order to describe the $^9\text{Be}$ excited states realistically, we need to introduce the $\Lambda\alpha$ spin-orbit force and solve the Faddeev equation in the $jj$-coupling scheme. The recent $\gamma$-ray spectroscopy experiment [30,31] indicates a very small spin-orbit splitting for the possible $5/2^+$ and $3/2^+$ resonances. It is interesting to examine the $LS$ components of the quark-model $\Lambda N$ interaction, in which the antisymmetric $LS$ interaction ($LS^{-\nu}$) is by about a factor two larger than in the Nijmegen models. We expect a large cancellation between the ordinary $LS$ interaction and this $LS^{-\nu}$ interaction. An interesting application of the present Faddeev formalism and the $\Lambda\alpha$ $T$-matrix derived in this study is to the recent Nagara event [61] for $^9\text{He}$. For the $\Lambda\Lambda$ interaction, we can use the coupled-channel $\Lambda\Lambda-\Xi NN-\Sigma \Sigma$ $T$-matrix of the quark-model interaction, fss2. A preliminary result [41] shows that fss2 is at present the only model which can reproduce an appropriate strength of the $\Lambda\Lambda$ interaction, $\Delta B^{\exp}_{\Lambda\Lambda} = 1.01\pm 0.20$ MeV, deduced from the Nagara event. In a separate paper [22], we have also reported another application of the present three-cluster Faddeev formalism to the hypertriton system, in which the quark-model $NN$ and $YN$ interactions are explicitly
used in the $\Delta NN$ and $\Sigma NN$ coupled-channel Faddeev formalism. In this system, a complete Pauli-forbidden state at the quark level exists in the $\Delta N \Sigma N$ subsystem.

**ACKNOWLEDGMENTS**

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**APPENDIX A: REARRANGEMENT FACTORS OF THREE-CLUSTER SYSTEMS WITH TWO IDENTICAL PARTICLES**

In this Appendix, we give a brief comment on the definition of the rearrangement factors in the Dirac notation for general three-body systems with two identical particles or clusters. The incorporation of spin-isospin degrees of freedom is essential for further applications to the hypertriton system [22] and the $\Lambda\Lambda\alpha$ system [41]. When one uses the Dirac notation, it is important to fix a coordinate system of the representation. We choose the standard system of the Jacobi coordinates with $\gamma=3$, and introduce the Jacobi coordinates in the momentum space, $\mathbf{p}=\mathbf{p}_3$ and $\mathbf{q}=\mathbf{q}_3$. The other Jacobi coordinates $\mathbf{p}_1$, $\mathbf{q}_1$, etc., are similarly defined. For an arbitrary function $\psi(\mathbf{p}, \mathbf{q}; 123)$ in $\gamma=3$, the effect of the cyclic permutation $P_{(123)}$ of the symmetric group $S_3$ is

$$P_{(123)} \psi(\mathbf{p}_3, \mathbf{q}_3; 123) = P_{(123)}^2 \psi(\mathbf{p}_1, \mathbf{q}_1; 231) = \psi(\mathbf{p}_3, \mathbf{q}_3; 123),$$

(A1)

where $123$ in $\psi(\mathbf{p}, \mathbf{q}; 123)$ stands for the spin-isospin variables. For the transposition $P_{(12)}$, Eq. (1) yields

$$P_{(12)} \psi(\mathbf{p}_3, \mathbf{q}_3; 123) = \psi(-\mathbf{p}_3, \mathbf{q}_3; 213),$$

$$P_{(12)} \psi(\mathbf{p}_1, \mathbf{q}_1; 231) = \psi(-\mathbf{p}_2, \mathbf{q}_2; 132),$$

$$P_{(12)} \psi(\mathbf{p}_2, \mathbf{q}_2; 312) = \psi(-\mathbf{p}_1, \mathbf{q}_1; 321).$$

(A2)

Note that the momentum suffix $\alpha$ in $\mathbf{p}_\alpha$, $\mathbf{q}_\alpha$, and the sign of $\mathbf{p}_\alpha$, etc., are uniquely specified by the sequence of 123. For example, $\langle \mathbf{p}_2, \mathbf{q}_2 | \beta \rangle$ in Eq. (4) actually implies $\langle \mathbf{p}_2, \mathbf{q}_2; 312 | \beta \rangle$. In the following, we always use an abbreviated notation, $\psi=\psi(\mathbf{p}_3, \mathbf{q}_3; 123)$, in the standard coordinate system $\gamma=3$. The total wave function, $\Psi(\mathbf{q}_3, \mathbf{q}_3; 123)$, in Eq. (2) is then compactly expressed as

$$\Psi = \psi \pm P_{(12)} P_{(123)}^2 \varphi + P_{(123)}^2 \varphi.$$

(A3)

If we write the Faddeev equation in terms of $\psi$ and $\varphi$, it reads

$$\psi = G_0 T (1 \pm P_{(12)}) P_{(123)}^2 \varphi,$$

$$\varphi = G_0 T [P_{(123)} \psi \pm P_{(123)} \varphi],$$

(A4)

with $T=T_{12}$ and $T=P_{(123)} T_{13} P_{(123)}^{-1}$, where $T_{12}$ and $T_{13}$ are the two-body $T$-matrices in the three-body space.

The definition of the rearrangement factors in the Dirac notation is based on the assumption

$$\langle \mathbf{p}_3, \mathbf{q}_3; 123 | P_{(123)}^2 \varphi \rangle = P_{(123)}^2 \varphi(\mathbf{p}, \mathbf{q}; 312) = \varphi(\mathbf{p}, \mathbf{q}; 231),$$

$$= \int d\mathbf{p}' d\mathbf{q}' \delta(\mathbf{p}' - \mathbf{p}_2) \delta(\mathbf{q}' - \mathbf{q}_2) \times P_{(123)}^2 \varphi(\mathbf{p}', \mathbf{q}'; 123),$$

(A5)

where the function $\varphi(\mathbf{p}, \mathbf{q}; 123)$ is $\varphi$ in Eq. (A3) and $P_{(123)}^2 \varphi$ operates only on the spin-isospin variables of $\varphi(\mathbf{p}', \mathbf{q}'; 123)$. With this $\varphi$ in the $\beta=2$ channel in mind, the standard procedure of the partial wave decomposition gives the following definition for the first-type rearrangement factor $g_{\gamma \rho}(\mathbf{q}, \mathbf{q}'; x)$:

$$\langle \mathbf{p}, \mathbf{q} | P_{(123)}^2 | \mathbf{p}', \mathbf{q}' ; \beta \rangle_{\gamma \rho} = \frac{1}{2} \int_{-1}^{1} dx \frac{\delta(p-p_1) \delta(p' - p_2)}{p^{\alpha+2} p'^{\alpha+2}} g_{\gamma \rho}(\mathbf{q}, \mathbf{q}' ; x) = \sum_{123} \int d\mathbf{p} d\mathbf{q} d\mathbf{p}' d\mathbf{q}' \langle \gamma | \hat{\rho} | \mathbf{q}; 123 \rangle \times \delta\left(p + \mathbf{q} + \frac{m_2}{m_2 + m_1} \mathbf{q}ight) \delta\left(p' - \mathbf{q} - \frac{m_3}{m_3 + m_1} \mathbf{q}'\right) P_{(123)}^2 \varphi(\mathbf{p}', \mathbf{q}'; 123) \beta.$$

(A6)

Here, $p_1$ and $p_2$ are given in Eq. (5c) with a general mass factor, $\zeta=(4m_3/m_1)$. With this mass modification, Eq. (6) is valid with a more complete reduced rearrangement factor.
The angular-momentum factors
\[ G_{Ls}(\lambda, \ell, \ell_2) = \left\{ \frac{(2\lambda + 1)! (2\ell + 1)!}{(2\ell_2)!} \right\}^{1/2} \lambda \ell \sum_{\ell_1 \ell_2} \langle \lambda \ell_1 | \ell_2 \rangle \langle \lambda \ell_2 | \ell_1 \rangle \langle \lambda \ell \rangle \langle \lambda \ell_1 | \ell_2 \rangle \langle \lambda \ell_2 | \ell_1 \rangle \langle \lambda \ell \rangle \] (A7)

Here the square bracket implies the unitary form of the 9j coefficients and the quantum numbers are specified by
\[ \left\{ \begin{array}{c} \gamma = [(\lambda \ell)LS]J_J; \langle TT \rangle, \\ \beta = [(\ell_1 \ell_2)LS]J_J; \langle TT \rangle \end{array} \right\} (LS-coupling), \]
\[ \left\{ \begin{array}{c} \gamma = [(\lambda_1 \lambda_2)I(\ell_1 \ell_2)J_J; \langle TT \rangle, \\ \beta = [(\ell_1' \ell_2')I(\ell_1 \ell_2)J_J; \langle TT \rangle \end{array} \right\} \] (jj-coupling). (A8)

The angular-momentum factors \( G_{Ls}(\lambda, \ell, \ell_2) \) with \( \lambda_1 = 0 - \lambda, \lambda_2 = 0 - \ell \) are given by
\[ G_{Ls}(\lambda, \ell, \ell_2) = G_{Ls}(\lambda_2, \ell, \ell_2) = \left\{ \frac{(2\lambda + 1)! (2\ell + 1)!}{(2\ell_2)!} \right\}^{1/2} \lambda \ell \sum_{\ell_1 \ell_2} \langle \lambda \ell_1 | \ell_2 \rangle \langle \lambda \ell_2 | \ell_1 \rangle \langle \lambda \ell \rangle \langle \lambda \ell_1 | \ell_2 \rangle \langle \lambda \ell_2 | \ell_1 \rangle \langle \lambda \ell \rangle \] (A9)

with respect to the interchange between \( p, q, \beta \) and \( p', q', \beta' \), since \( m_1 = m_2 \).

**APPENDIX B: A USEFUL FORMULA FOR THE \( \Lambda \alpha \) BORN KERNEL**

The general procedure to calculate Born kernels of the s-shell clusters, developed in Ref. [62], can also be used to calculate the \( \Lambda \alpha \) Born kernel
\[ V(q, q') = \langle e^{iqr} | V | e^{iq' r} \rangle = \langle e^{iqr} \xi_{\Lambda} \phi_\alpha | \sum_{j=2}^5 u_{1j} | e^{iq' r} \xi_{\Lambda} \phi_\alpha \rangle, \]
(B1)

where \( \phi_\alpha \) is the internal wave function of the \( \alpha \) cluster, \( \xi_{\Lambda} \) is the spin wave function of the \( \Lambda \) particle and \( u_{1j} \) is an effective \( \Lambda \Lambda \) interaction. The essential part of this method lies in the correct treatment of the c.m. motion which is handled by the procedure given in Ref. [63]. This method makes it possible to deal with the most general form of the \( \Lambda \Lambda \) interaction with non-static effects like the \( G \)-matrix \( \Lambda \Lambda \) interaction. In this method, \( V(q, q') \) in Eq. (B1) is calculated from an integral form of the GCM kernel through
\[ V(q, q') = \langle \delta(X_c) e^{iqr} \xi_{\Lambda} \phi_\alpha | \sum_{j=2}^5 u_{1j} | e^{iq' r} \xi_{\Lambda} \phi_\alpha \rangle = \left( \frac{\gamma}{2 \pi} \right)^{3/2} e^{(1/4)(q'^2 + q^2)} \int da db e^{-iqr + q'r} G(a, b), \]
(B2)
Here, $\gamma_G = (4 + \xi)\nu$, $\nu = 4\xi\nu/(4 + \xi)$ with $\xi = M_A/M_N$, and $\psi_\alpha(R)$ and $\psi_\alpha(R)$ are the h.o. shell model wave functions of $\Lambda$ and $\alpha$, centered at $R$, with the width parameters $\xi\nu$ and $\nu$, respectively. First we calculate spatial integrals for the spatial part $u\cdot v$, and $\nabla u\cdot \nabla v$. These four integrals with $j = 2 - 5$ are all equal because of the antisymmetry of the $\alpha$ cluster. We need to calculate spatial integrals for $u = u(x_1 - x_2)$ and $v = v(x_1 - x_2)P_r$, which we call the direct term and the exchange term, respectively. It is important to note that the space exchange operator $P_r$, operates only on the single-particle coordinates $x_1$ and $x_2$, and does not exchange the $\Lambda$ and $N$ masses. The procedure to interchange these masses $M_A$ and $M_N$ simultaneously like in Ref. [7] leads to an erroneous expression [see Eq. (A.1) of [7]], which is apparently wrong since the RGM kernel $\langle \delta(r-a)\xi_\lambda\Phi_\alpha|\delta(r-b)\xi_\lambda\Phi_\alpha \rangle$ should not involve the mass dependence. The correct expression is the one in which one sets $M_N = M_A$ in their Eq. (A.1) [see Eq. (B8) below]. The most general form of the two-body $\Lambda N$ matrix elements for the translationally invarient $u$ is parametrized as

$$\langle p_1p_2|u(p_1',p_2')\rangle = \frac{1}{(2\pi)^3} \delta(p - p')u(k',q';P).$$

with $p = (p_1 - q_2)/(\xi + 1)$, $P = p_1 + p_2$ (also $P'$ for $p_1', p_2'$), and $k' = p - p'$, $q' = (p + p')/2$. For the matrix element Eq. (B4), the spatial part of the GCM kernel in Eq. (B3) is calculated to be

$$G_{\text{space}}^{\alpha}(a,b) = \frac{1}{(2\pi)^3} \left(\frac{4 + \xi}{3\xi}\right)^{3/2} \int dP dk' dq'|u(k',q';P) \times \exp\left[-\frac{1}{6\nu\xi + 1} \frac{(q' - 1)^2 + k'^2}{2\nu}\right] + i(a - b) \cdot \left(q' + \frac{\xi + 1}{\xi + 1}P\right) + i(\xi + 1)k'^2].$$

(B5)

If we use Eq. (B5) in Eq. (B2), we can perform the integrals over $a$ and $b$ and obtain two delta functions. Thus we can perform the integrals over $k'$ and $q'$ and obtain a compact formula

$$G_{\text{space}}^{\alpha}(a,b) = \left(\frac{\gamma_G}{2\pi}\right)^{3/2} \int dR \langle \psi_\alpha(R) | \psi_\alpha(0) \rangle \times \sum_{j=2}^5 \phi_j(R) \psi_\alpha(R).$$

(B3)

where $k = q_1 - q_2$, and $q = (q_1 + q_2)/2$. For a simple local Gaussian interaction, we find

$$V(r.q.;P) = \frac{1}{3\nu + \frac{8}{3}\kappa} \exp\left[-\frac{1}{4}(\frac{3}{8\nu} + \frac{1}{2}\kappa)k^2\right] \text{ for } u(r) = \exp(-kr^2).$$

Then the $P$ integral is carried out easily and we obtain

$$V_{\alpha}(q_1 q_2) = \left(\frac{\gamma_G}{2\pi}\right)^{3/2} \int dP e^{-(3/2\nu + \frac{1}{2}\kappa)k^2} \times u\left(k - \frac{\xi + 1}{\xi + 1}P, P + \frac{3}{4}q\right).$$

(B6)

If we further incorporate the spin-isospin factors, the full $V(q_1, q_2)$ is given by

$$V(q_1, q_2) = X_dX_eV_{d}(q_1, q_2) + X_eV_{e}(q_1, q_2).$$

(B9)

with the spin-isospin factors defined by

$$\left\{ \begin{array}{l} X_d = (\xi_\lambda\chi_\alpha) \sum_{j=2}^5 \phi_j \xi_\lambda(\chi_\alpha) \end{array} \right.$$
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