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Faddeev calculation of the hypertriton using the $SU_6$ quark-model nucleon-nucleon and hyperon-nucleon interactions

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Quark-model nucleon-nucleon and hyperon-nucleon interactions by the Kyoto-Niigata group are applied to the hypertriton calculation in a three-cluster Faddeev formalism using the two-cluster resonating-group method kernels. The most recent model, fss2, gives a reasonable result similar to the Nijmegen soft-core model NSC89, except for an appreciable contribution of higher partial waves.

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I. INTRODUCTION

The QCD-inspired spin-flavor $SU_6$ quark model for the baryon-baryon interaction, proposed by the Kyoto-Niigata group, is a unified model for the complete baryon octet ($B_8 = N, \Lambda, \Sigma$ and $\Xi$), which has achieved very accurate descriptions of the nucleon-nucleon ($NN$) and hyperon-nucleon ($YN$) interactions [1–3]. In particular, the $NN$ interaction of the most recent model fss2 [2] is accurate enough to compare with the modern realistic meson-exchange models. These quark-model interactions can be used for realistic calculations of few-baryon and few-cluster systems, once an appropriate three-body equation is formulated for the pairwise interactions described by the resonating-group method (RGM) kernels. The desired three-cluster equation should be able to deal with the nonlocality and the energy dependence intrinsically involved in the exchange-RGM kernel. Furthermore, the quark-model description of the $YN$ and hyperon-hyperon ($YY$) interactions in the full coupled-channel formalism sometimes involves a Pauli forbidden state at the quark level, which excludes the most compact $(0s)^6$ spatial configuration, resulting in the strongly repulsive nature of the interactions in some particular channels. We have recently formulated a three-cluster equation which uses two-cluster RGM kernels explicitly [4]. This equation exactly eliminates three-cluster redundant components by requiring the orthogonality of the total wave function to the pairwise two-cluster Pauli-forbidden states. The explicit energy dependence inherent in the exchange RGM kernel is self-consistently treated. This equation is entirely equivalent to the Faddeev equation which uses a singularity-free $T$ matrix (which we call the RGM $T$ matrix) generated from the two-cluster RGM kernel.

We first applied this formalism to three di-neutron and three-alpha systems, and obtained complete equivalence between the Faddeev calculations and variational calculations using the translationally invariant harmonic-oscillator (h.o.) basis [4,5]. This formalism was also applied to the Faddeev calculation of the three-nucleon bound state [6], employing the off-shell $T$ matrices which are derived from the nonlocal and energy-dependent RGM kernels for our quark-model $NN$ interactions, fss2 and FSS. The model fss2 yields the 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 [7]. The effect of the charge dependence of the two-body $NN$ interaction is estimated to be $0.19$ MeV for the triton binding energy [8]. 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, $0.5 \sim 1$ MeV [9], predicted by many Faddeev calculations employing modern realistic meson-theoretical $NN$ interactions. The charge rms radii for $^3$H and $^3$He are also correctly reproduced. The non-local description of the short-range repulsion in the quark model is essential to reproduce the large binding energy and the correct size of the three-nucleon bound state without reducing the D-state probability of the deuteron.

Here we apply our quark-model $NN$ and $YN$ interactions to the hypertriton ($^3$H) with the small separation energy of the $\Lambda$ particle, $B_\Lambda^{exp} = 130 \pm 50$ keV. Since the $\Lambda$ particle is far apart from the two-nucleon subsystem, the on-shell properties of the $\Lambda N$ and $\Sigma N$ interactions are expected to be well reflected in this system. In particular, this system is very useful to determine the relative strength of $^3S_0$ and $^3S_1$ interactions in our framework. We will be able to fine tune the quark model interaction to the hypertriton binding energy. This enables firmer quark-model predictions for the $^3S_0$ and

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\[ \sum_{i=1,2} |\mu_{3i}\rangle\langle \mu_{3i}| \Psi_\lambda = \lambda |\Psi_\lambda \rangle \] (4)

in the antisymmetric model space, \( \Psi_\lambda(123) = - \Psi_\lambda(213) \), which we denote as \( \Psi_\lambda \in [11] \) in the following. The projection operator on the Pauli-allowed space, \( P \), is defined by taking the model space spanned by the eigenvectors with \( \lambda = 0 \)

\[ P = \sum_{\lambda=0} |\Psi_\lambda\rangle\langle \Psi_\lambda| . \] (5)

The three-cluster Faddeev equation, which is completely equivalent to Eq. (1), is derived by using some nice properties satisfied by \( P \), which is briefly discussed in the Appendix. The total wave function of the hypertriton system is expressed as a superposition of two independent Faddeev components \( \psi \) and \( \phi \): \( P\Psi = \psi + (1 - P)\phi \) with \( \psi \in [11] \). The coupled-channel Faddeev equation reads

\[ \psi = G_0(E) T_{12}(E,\epsilon_{12})(1 - P_{12})\phi, \] (6)

where \( P_{12} \) is the exchange operator of particles 1 and 2, \( G_0(E) = 1/(E - H_0) \) is the three-body free Green function for the negative energy \( E < 0 \), \( T_{12}(E,\epsilon_{12}) \) is the NN \( T \) matrix in the three-body space, and \( T_{31}(E,\epsilon_{31}) \) is the redundancy-free \( \Lambda N \Sigma N \bar{T} \) matrix in the coupled-channel formalism. These \( T \) matrices are generated from the NN and \( Y N \) RGM kernels \( V_{RGM}(\epsilon_{NN}) \) and \( V_{RGM}(\epsilon_{YN}) \), respectively, according to the prescription essentially given in Ref. [4]. The energy dependence involved in these kernels is treated self-consistently by calculating the matrix elements of the two-cluster Hamiltonian, which will be discussed in some detail.

For the \( NN \) sector, \( T_{12}(E,\epsilon_{12}) \) in the \( \Lambda N \) and \( \Sigma N \) spaces are given by the two-body \( T \) matrix \( t(\omega, \epsilon) \) through a simple replacement of the starting-energy argument

\[ T_{12}(E,\epsilon_{12}) = \begin{pmatrix} t_{NN} & 0 \\ 0 & \frac{\xi}{t_{NN}} \end{pmatrix}, \]

\[ t_{NN} = t_{NN}(E - \frac{\hbar^2}{M_N} \frac{\xi_\Lambda + 2}{4 \xi_\Lambda} \epsilon_{NN}^2), \]

\[ \frac{\xi}{t_{NN}} = t_{NN}(E - \Delta M_{\Sigma} \frac{\hbar^2}{M_N} \frac{\xi_\Sigma + 2}{4 \xi_\Sigma} \epsilon_{NN}^2), \] (7)

where \( M_N \) is the nucleon mass, \( \xi_\Lambda = M_\Lambda/M_N \) is the mass ratio for \( \beta = \Lambda \) or \( \Sigma \), and \( q \) is the momentum of the residual \( \Lambda \) or \( \Sigma \). The \( NN \) relative energy in the three-body space, \( \epsilon_{NN} \), is determined from

\[ \epsilon_{NN} = \langle \Psi | P | \Psi_{NN} + V_{RGM}(\epsilon_{NN}) | \Psi \rangle, \] (8)

which is actually the sum of the \( \Lambda N \) and \( \Sigma N \) components.

For the \( Y N \) sector, the situation is more complicated since the \( \Lambda N \Sigma N \) coupling involves a complete Pauli forbidden state \( |\mu_{33}\rangle \) in the three-body space and the difference of the momentum dependent starting energies in the \( \Lambda N \) and \( \Sigma N \).
channels is not constant because of the two types of reduced masses between $YN$ and $N$. In fact, the equation satisfied by $\tilde{T}_3(E, \epsilon_{31})$ can never be reduced to the two-cluster coupled-channel equation in the free space. We request

$$T_{31}(E, \epsilon_{31}) = V_{31}^{\text{RGM}}(\epsilon_{31}) + V_{31}^{\text{RGM}}(\epsilon_{31}) G_0(E) T_{31}(E, \epsilon_{31})$$

$$\tilde{T}_3(E, \epsilon_{31}) = T_{31}(E, \epsilon_{31}) - (E - H_0) |u_{31}|$$

$$\times [E - \epsilon_{31} - h_2 - \Delta m_3]^{-1} |u_{31}| (E - H_0). \quad (9)$$

However, this equation is not actually satisfied, since the derivation is based on the unrealistic assumptions

1. $[E - \epsilon_{31} - h_2 - \Delta m_3]$ is channel independent,
2. $[\epsilon_{31} - h_2 - V_{31}^{\text{RGM}}(\epsilon_{31})]|u_{31}|$ is satisfied.

The second condition is only approximately satisfied since, in the strict RGM framework, it is not possible to use the empirical internal energies of clusters and reduced masses for the relative kinetic-energy operators. In fact, the correct expression for (2) is

$$\langle u_{31} \rangle [1 + G_0(\omega) \tilde{T}_{31}(E, \epsilon_{31})] = 0, \quad (11)$$

which is essential to yield the orthogonality of the total wave function through the second equation, Eq. (6): $\langle u_{31} | \psi + (1 - P_3 \chi) \rangle = 0 \quad (i=1,2)$. 

Fortunately, these problems are completely solved by simply adding a small correction term to the RGM kernel, which is a procedure developed in Ref. [15] for making it possible to use the empirical values of the internal energies and reduced masses in the RGM formalism. In Ref. [15], we have slightly modified the original two-cluster RGM equation $[\epsilon_{31} - h_2 - V_{31}^{\text{RGM}}(\epsilon_{31})]\chi = 0$ and considered the following RGM equation in the OCM (orthogonality condition model) form:

$$\Lambda_0[\omega - h_0^{\text{exp}} - V_{\text{RGM}}(\epsilon_0)] \chi = 0. \quad (12)$$

From here on, we omit the subscript 31 or $YN$ as much as possible, in order to simplify the notations. For example, $V_{\text{RGM}}(\epsilon_0) = V_{31}^{\text{RGM}}(\epsilon_{31})$ and $\Lambda_0 = \Lambda_{31} = 1 - |u\rangle\langle u|$ with $|u\rangle = |u_{31}\rangle$ in Eq. (12). Furthermore, $\epsilon = \epsilon_{AN} - E_{\text{int}}$ in Eq. (12) uses the empirical internal energy, $E_{\text{int}}^{\text{exp}} = \Delta m_3$, and the relative kinetic-energy operator, $h_0^{\text{exp}} = h_2 - 2 \mu^{\text{exp}} \nabla^2$, uses the empirical reduced mass, $\mu^{\text{exp}} = \mu_{AN}^{\text{exp}}$ or $\mu_{SN}^{\text{exp}}$. On the other hand, we need to use $\epsilon_0$ with the calculated internal energies in the RGM kernel. It is shown in Ref. [15] that Eq. (12) is converted to Eq. (3) by simply adding $\Delta G$ to $V_{\text{RGM}}(\epsilon_0)$; i.e., $V_{31}^{\text{RGM}}(\epsilon_{31}) = V_{\text{RGM}}(\epsilon_0) + \Delta G$. The explicit expression of $\Delta G$ is given in the paper [or Case 0 case in Eq. (16) below]. We use the same idea to eliminate the channel dependence of $[E - \epsilon_{31} - h_2 - \Delta m_3]$ in (1). Let us start from the off-shell extension of the $\Lambda N - \Sigma N$ coupled-channel RGM equation in the following OCM form:

$$\Lambda_0[\omega - h_0^{\text{exp}} - V_{\text{RGM}}(\epsilon_0)] \chi |u\rangle (\omega - \epsilon)(\omega - \epsilon) \chi = 0. \quad (13)$$

The solution of this equation satisfies the orthogonality, $\langle u | \chi \rangle = 0$ for $\omega \neq \epsilon$. The various energies in Eq. (13) are usually channel dependent; namely, if we use the label $\beta = \Lambda$ or $\Sigma$ to specify the $\Lambda N$ or $\Sigma N$ channel, the diagonal matrix elements of these energies are given by

$$\omega = E - \frac{\hbar^2}{2M_N} \left( \frac{\zeta + 2}{\zeta + 1} \right) q^2 - E_{\text{int}}^{\text{exp}},$$

$$e = e_{\text{AN}} - \frac{\hbar^2}{2M_N} \left( \frac{\zeta + 2}{\zeta + 1} \right) q^2 - E_{\text{int}}^{\text{exp}},$$

$$e_0 = e_{\text{AN}} - E_{\text{int}}.$$ 

where $E$ is the negative three-body energy, $q$ this time is the momentum of the residual nucleon. We can prove that Eq. (13) is equivalent to the following Schrödinger-type RGM equation:

$$\omega \chi = \left[ E_{\text{int}}^{\text{exp}} + V_{\text{RGM}}^{\text{mod}}(\epsilon) \right] \chi, \quad (15)$$

where a newly defined RGM kernel $V_{\text{RGM}}^{\text{mod}}(\epsilon)$ is given by

$$V_{\text{RGM}}^{\text{mod}}(\epsilon) = V_{\text{RGM}}(\epsilon_0) + \Delta G,$$

$$\Delta G = \Lambda (\Delta E_{\text{int}} + \Delta \epsilon + \Delta h_0) - (\Delta E_{\text{int}} + \Delta \epsilon + \Delta h_0).$$

$$\Delta E_{\text{int}} = E_{\text{int}}^{\text{exp}} - E_{\text{int}}, \quad \Delta h_0 = h_0^{\text{exp}} - h_0,$$

$$\Delta \epsilon = \frac{\hbar^2}{2M_N} \left( \frac{1}{\zeta + 1} - \frac{1}{\zeta + 1} \right) q^2.$$ 

The RGM $T$ matrix $\tilde{T}_{31}(E, \epsilon_{31})$ is therefore formulated for this modified RGM kernel $V_{\text{RGM}}^{\text{mod}}(\epsilon)$. By repeating the same process as to derive Eq. (9) with respect to $V_{\text{RGM}}^{\text{mod}}(\epsilon)$, we can find that Eq. (9) [and also Eq. (11)] is just valid if we replace (see Appendix for details)

$$V_{31}^{\text{RGM}}(\epsilon_{31}) \rightarrow V_{\text{RGM}}^{\text{mod}}(\epsilon),$$

$$[E - \epsilon_{31} - h_2 - \Delta m_3]^{-1} \rightarrow [E - e_{\text{AN}} - \left( h_2 \right)_\Lambda]^{-1}. \quad (17)$$

In order to determine $e_{\text{AN}}$ or $\epsilon_0$, we approximate $\omega$ in Eq. (15) as

$$\omega = \begin{pmatrix} E - \frac{\hbar^2}{2M_N} \left( \frac{\zeta + 2}{\zeta + 1} \right) q^2 & 0 \\ 0 & E - \frac{\hbar^2}{2M_N} \left( \frac{\zeta + 2}{\zeta + 1} \right) q^2 - \Delta M_{\Sigma N} \end{pmatrix} \begin{pmatrix} e_{\text{AN}}^0 \\ e_{\Sigma N}^0 \end{pmatrix}$$

with keeping the relationship $\epsilon_{\Sigma N} = e_{\text{AN}} - \Delta M_{\Sigma N}$ in the free space, and first calculate a $\Lambda N - \Sigma N$ averaged value.
\[ e_{YN} = \langle P| \hat{H}_0^{\text{exp}} + V_{\text{mod}}^{\text{RGM}}(\theta)|P\rangle. \] (19)

Then, from \( e_{YN} = e_{AN} \langle \Psi | \Psi \rangle^\Delta + e_{\Sigma N} \langle \Psi | \Psi \rangle^\Sigma \) and \( \Delta M_{\Sigma A} = e_{AN} - e_{\Sigma N} \), we find
\[ e_{\Lambda N} = e_{YN} + \Delta M_{\Sigma A} P_{\Sigma}, \]
\[ e_{\Sigma N} = e_{YN} - \Delta M_{\Sigma A} (1 - P_{\Sigma}), \] (20)
where \( P_{\Sigma} = \langle \Psi | \Psi \rangle^\Sigma \) is the probability of the \( \Sigma NN \) component admixed in the hypertriton wave function.

### III. RESULT

The angular-momentum states of the \( NN \)-channel are specified by \[ \{(\lambda s)I, (\ell s)I\}|I,(\ell \pm \frac{1}{2})r_{\lambda}\rangle \rangle \}, \]
where \( (\lambda s)I \) stands for the two-nucleon state, \( \gamma \) and \( \ell \) is the isospin value \( t=0 \) for \( \Lambda \) and \( t=1 \) for \( \Sigma \). Due to the antisymmetric property of the two nucleons, \(-1)^{l+s+1}=1\), we find the \( ^3E \) and \( ^1O \) states contribute to the \( \Lambda NN \) channel and the \( ^1E \) and \( ^3O \) states contribute to the \( \Sigma NN \) channel. For the \( YN-N \) channel, the angular-momentum states are specified by \[ (\ell \ell J)I_1, (\ell \ell J)I_2\rangle |I, (\ell \pm \frac{1}{2})r_{\lambda}\rangle \rangle \rangle \). Since the isospin of the hypertriton is zero, only the isospin \( T=1/2 \) sector of the \( \Lambda-N \)-\( \Sigma N \) interaction contributes to the hypertriton calculation. All the partial waves of the orbital angular momentum are possible for each of the \( \Lambda NN \) and \( \Sigma NN \) channels, which makes the number of channels for a particular partial-wave truncation just three times larger than in the triton Faddeev calculation. The hyperon species of the \( YNN \) channels are uniquely specified by the isospin value \( t=0 \) or \( 1 \). For the orbital part, the parity conservation requires \(-1)^{l+s+1} = (-1)^{l+s+1} = 1\). The channel truncation is specified by the maximum value of the total angular momenta of the pairwise baryons, \( l \) and \( s \), which we denote \( J \). As an example, all the channels of the standard 15-channel calculation with \( S \) and \( D \) waves are listed in Table I.

For the numerical calculation, we discretize the continuous momentum variable \( p \) (or \( q \)) for the Jacobi coordinate vectors, using the Gauss-Legendre \( n_1 \) (or \( n_2 \)) point quadrature formula, for each of the three intervals of \( 0-1 \) \( \text{fm}^{-1} \), \( 1-3 \) \( \text{fm}^{-1} \) and \( 3-6 \) \( \text{fm}^{-1} \). The small contribution from the intermediate integral over \( p \) beyond \( p_0=6 \) \( \text{fm}^{-1} \) in the two-body \( T \)-matrix calculation is also taken into account by using the Gauss-Legendre \( n_1 \)-point quadrature formula through the mapping \( p=p_0+\tan\left(\pi(1+x)/4\right)\). These \( n_1 \) points for \( p \) are not included for solving the Faddeev equation Eq. (6), since it causes a numerical inaccuracy for the interpolation.] The momentum region \( q=6 \text{ fm}^{-1} \) is also discretized by the \( n_1 \)-point formula just as with the \( p \) discretization case. We take \( n_1=n_2=n_3=10-10.5 \) as is used for the triton Faddeev calculation in Refs. [6,7], for which well converged results are obtained. The partial-wave decomposition of the two-cluster RGM kernel is carried out numerically using the Gauss-Legendre 20-point quadrature formula. The modified spline interpolation technique developed in Ref. [16] is employed for simplifying the treatment of the rearrangement of the Jacobi momentum coordinates. The Faddeev formalism with two identical particles or clusters is discussed in Ref. [17],

### TABLE I

The channel quantum numbers of the hypertriton included in 15-channel Faddeev calculation with \( S \) and \( D \) waves. For the \( NN \)-\( Y \) channels, \( l \) and \( j \) are coupled to the total angular momentum \( \frac{1}{2}^+ \) and the isospin coupling is \( \left[ \frac{1}{2}^+ \right] \). For the \( YN \)-\( YN \) channels, these are \[ [l_1^l_{1/2}] \] and \[ [l_1^d_{1/2}, l_1^s_{1/2}] \]. The isospin value \( t \) specifies the hyperon species; i.e., \( Y=\Lambda \) for \( t=0 \) and \( Y=\Sigma \) for \( t=1 \).

| \( NN-Y \) | \( |l_{1/2}^j \rangle |l_{1/2}^j \rangle \) | \( t \) |
|---|---|---|
| 1 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 2 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 3 | \( ^3D_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 4 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 5 | \( ^3S_0 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |

| \( YN-N \) | \( |l_{1/2}^j \rangle |l_{1/2}^j \rangle \) | \( t \) |
|---|---|---|
| 1 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 2 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 3 | \( ^3D_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 4 | \( ^3S_1 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |
| 5 | \( ^3S_0 \) | \( \left( \frac{3}{2} \right)^{1/2} \) |

### ACKNOWLEDGMENTS

Together with some formulas for calculating the matrix elements of the two-cluster Hamiltonian. For the diagonalization of the large nonsymmetric matrix appearing in solving Faddeev equations, the Arnoldi-Lanczos algorithm developed in the ARPACK subroutine package [18] is very useful.

Table II shows the results of the Faddeev calculations using fss2 and our previous model FSS. In the 15-channel calculation including the \( S \) and \( D \) waves of the \( NN \) and \( YN \) interactions, we have already obtained \( B_\Lambda = -0.80 \text{ keV} \) for fss2. The convergence with the extension to the higher partial waves is very rapid, and the total angular momentum of the baryon pairs with \( J=4 \) is good enough for 1 keV accuracy. As for the converged \( B_\Lambda \) values with 150-channel \( \Lambda NN \) and \( \Sigma NN \) configurations, we obtain \( B_\Lambda = 289 \text{ keV} \) and the \( \Sigma NN \) component \( P_\Sigma = 0.80\% \) for the fss2 prediction, and \( B_\Lambda = 878 \text{ keV} \) and \( P_\Sigma = 1.36\% \) for FSS. Table III shows the correlation between the \( \Lambda \) separation energies \( B_\Lambda \) and the \( ^1S_0 \) and \( ^3S_1 \) effective range parameters of FSS, fss2 and NSC89 \( \Lambda N \) interactions. Although all of these \( \Lambda N \) interactions reproduce the low-energy \( \Lambda N \) total cross section data below \( p_\Lambda \sim 300 \text{ MeV/c} \) within the experimental error bars, our quark-model interactions seem to be more attractive than the Nijmegen soft-core potential NSC89 [13]. The model FSS gives a large overbinding since the \( ^1S_0 \) \( \Lambda N \) interaction is strongly attractive. The phase-shift difference of the \( ^1S_0 \) and \( ^3S_1 \) states at \( p_\Lambda \sim 200 \text{ MeV}/c \) is \( \delta(^1S_0) - \delta(^3S_1) \sim 29^\circ \) for FSS, while \( \delta(^1S_0) - \delta(^3S_1) \sim 7^\circ \) for fss2. Since the present fss2 result is still slightly overbound, this difference should be somewhat smaller in order to reproduce the correct experimental value \( B_\Lambda = 130 \pm 50 \text{ keV} \). From the two results given by fss2 and FSS, we can extrapolate that, in our quark model, the desired difference is \( 0^\circ \sim 2^\circ \).

In order to make sure that this extrapolation gives a good estimation, we modify the \( \kappa \)-meson mass of the model fss2...
from the original value, \( m_{ee^2} = 936 \text{ MeV} \) [2], to 1000 MeV, and repeat the whole calculation. It is known that this modification makes the \( ^1S_0 \) \( \Lambda N \) interaction less attractive and the \( ^3S_1 \) more attractive. We obtain \( B_\Lambda = 145 \text{ keV} \) with \( P_\Sigma = 0.53\% \). The effective range parameters of this modified fss2 interaction are \( a_s = -2.15 \text{ fm} \), \( r_s = 3.05 \text{ fm} \), and \( a_t = -1.80 \text{ fm} \), \( r_t = 2.87 \text{ fm} \). The phase-shift difference is only 1.3° and the total cross section of the \( \Lambda N \) scattering increases at most 10 mb at \( P_\Lambda = 100 \text{ MeV} / c \) from 286 mb to 296 mb, which is still within the experimental error bars.

It should be kept in mind that the effective range parameters or the \( S \)-wave phase-shift values determined in this way are very much model dependent, since the \( B_\Lambda \) value is not solely determined by these quantities. It depends on how higher partial waves contribute and also on the details of the \( \Lambda N - \Sigma N \) coupling of a particular model. A nice extrapolation shown here is based on the similarity of the models fss2 and FSS, which have a common framework for the quark sector and the effective meson-exchange potentials.

Table II also shows that the expectation value of the \( NN \) Hamiltonian, \( \epsilon_{\text{NN}} \), determined self-consistently is rather close to the deuteron energy \( -\epsilon_d \), especially in fss2. This feature is even marked if we decompose these energies to the kinetic-energy and potential-energy contributions. Table IV shows this decomposition with respect to fss2, FSS and NSC89. (For this comparison, we use the definition of the kinetic-energy part of the deuteron by \( h_{d} = \langle \chi_d | V_{NN} | \chi_d \rangle / \langle \chi_d | \chi_d \rangle \), where \( \chi_d \) is the RGM relative wave function between the neutron and the proton.) In fss2, the kinetic energy of the NN subsystem is 1.88 MeV larger than that of the deuteron, which implies that the NN subsystem shrinks by the effect of the outer \( \Lambda \) particle, in comparison with the deuteron in the free space. In NSC89, this difference is even smaller; i.e., 1.18 MeV. These results are consistent with the fact that the hypertriton in NSC89 is more loosely bound (\( B_\Lambda = 143 \text{ keV} \)) than in fss2 (289 keV), and the \( \Lambda \) particle is very far apart from the NN cluster. The \( \Sigma NN \) probability in NSC89 is \( P_\Sigma = 0.5\% \) [10,12], Table IV also lists the kinetic-energy and potential-energy decompositions for the averaged \( YN \) expectation value \( \epsilon_{YN} \) and the total energy \( E \). The kinetic energies of \( \epsilon_{YN} \) are much smaller than those of \( \epsilon_{NN} \), which indicates that the relative wave functions between the hyperon and the nucleon are widely spread in the configuration space. The comparison of the total-energy decomposition shows that the wave functions of fss2 and NSC89 may be very similar. A clear difference between fss2 and NSC89 appears in the roles
of higher partial waves. The energy increase due to the higher partial waves than the $S$ and $D$ waves is 91 keV in fss2 and 168 keV in FSS, respectively. On the other hand, the results in Ref. [10] imply that this is only 20–30 keV in the case of NSC89. This difference can originate from both of the $NN$ and $YN$ interactions. Since the characteristics of the meson-theoretical $YN$ interactions in higher partial waves are $a$ priori unknown, more detailed analysis of the fss2 results might shed light on the adequacy of the quark-model baryon-baryon interactions.

IV. SUMMARY

In this study, we have carried out the Faddeev calculations of the hypertriton, using the recent quark-model $NN$ and $YN$ interactions, FSS [1] and fss2 [2,3]. These are realistic interactions which describe all the available $NN$ and $YN$ data, by incorporating the effective meson-exchange potentials at the quark level. Since these quark-model baryon-baryon interactions are formulated in the RGM framework, they are nonlocal, energy dependent, and sometimes involve the Pauli-forbidden component at the quark level. The hypertriton is an appropriate place to investigate the roles of the compact $SU_4\ (11)_5$ component, which is completely Pauli forbidden in the $^1S_0\ \Lambda N-\Sigma N$ channel coupling with the isospin $T=1/2$. In order to deal with this off-shell effect of the quark-model interaction precisely, we have formulated a type of the Faddeev equation which explicitly employs the two-cluster RGM kernels [4,5]. The energy dependence of the RGM kernel is self-consistently treated, by calculating the expectation values of the two-cluster Hamiltonian with respect to the obtained solutions of the Faddeev equation [17]. We have especially paid attention to how to extend the microscopic description of the $\Lambda N-\Sigma N$ coupling in the hypertriton system without spoiling the effect of the Pauli principle at the quark level. The present study is the second application of this formalism to the few-baryon systems interacting via the quark-model baryon-baryon interactions, following our previous one to the triton system [6,7]. The hypertriton is well suited to investigate the on-shell properties of the $\Lambda N$ and $\Sigma N$ interactions, since the hyperon is very far apart from the two-nucleon cluster.

We have found that our quark-model interaction fss2 gives a reasonable result for the hypertriton properties, which is rather similar to the result of the Nijmegen soft-core model NSC89 [13]. The $\Lambda$ separation energy given by fss2 is $B_\Lambda =289$ keV, which is a little too large in comparison with the experimental value $B_\Lambda^{exp}=130\pm50$ keV. The admixture of the $\Sigma NN$ component is $P_\Sigma=0.80\%$. Modifying the $\kappa$-meson mass of fss2 from the original value, $m_\kappa c^2=936$ MeV, to 1000 MeV leads to the almost correct $\Lambda$-separation energy 145 keV with $P_\Sigma=0.53\%$. Unlike the NSC89 result, the effects of higher partial waves up to the $G$ wave are rather important in the quark-model $NN$ and $YN$ interactions. If we use the dominant $S$-wave character of the $\Lambda N$ interaction in the hypertriton system, the $^1S_0\ \Lambda N$ interaction of the model fss2 is slightly too attractive. It is a future problem to investigate whether or not a reduction in the $^1S_0$ attraction like the modification to $m_\kappa c^2=1000$ MeV produces a favorable feature for the level spacing of the $0^+$ and $1^+$ states of $\Lambda H$ and $\Lambda$He systems. The fairly large charge symmetry breaking in these systems is also an important issue to understand the hyperon-nucleon interactions in detail.

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APPENDIX: DERIVATION OF THE THREE-CLUSTER FADDEEV EQUATION FOR THE $\Lambda NN-\Sigma NN$ SYSTEM

In this Appendix, we discuss some essential points to derive the Faddeev equation Eq. (6) from the three-cluster equation Eq. (1). In the derivation, we extensively use the following properties of the projection operator $P$ defined in Eq. (5):

(i) $\Lambda_i P = P \Lambda_i \ (i = 1, 2)$,

\[ \text{where } \Lambda_i = 1 - |u_{3i}|^2 |u_{3i}|, \]

(ii) when $\Psi \in [11]$, \ $\forall \ (u_{3i}|\Psi) = 0 \Rightarrow P \Psi = \Psi$,

(iii) when $\Psi \in [11]$, \ $P \Psi = 0 \Rightarrow \exists \ |\Psi|$,

\[ \Psi = |u_{3f2}|^2 - |u_{3f1}|. \]  

(A1)

Using the property (i), we can replace $V_{3i}^{RGM}(e_{3i})$ in Eq. (1) by $v_{3i}(e_{3i}) = \Lambda_i V_{3i}^{RGM}(e_{3i}) \Lambda_i$, or

\[ V_{3i}(E, e_{3i}) = (E - H_0) - \Lambda_i (E - H_0) \Lambda_i + v_{3i}(e_{3i}). \]  

(A2)

We further use the property (iii) for the whole equation and introduce the ansatz for the Faddeev components, $P \Psi = \psi + (1 - P_{12}) \phi$, to derive a pair of equations

\[ [E - H_0 - V_{12}^{RGM}(e_{12})] \psi = V_{12}^{RGM}(e_{12})(1 - P_{12}) \phi, \]

\[ [E - H_0 - V_{31}(E, e_{31})] \phi = V_{31}(E, e_{31})(\psi - P_{12} \phi) + |u_{3f2}|. \]  

(A3)

In the second equation, we note that

\[ E - H_0 - V_{31}(E, e_{31}) = \Lambda_{31}\ [E - H_0 - V_{31}^{RGM}(e_{31})]\Lambda_{31}, \]

(A4)

and introduce the projected two-body Green function in the three-body space, $G_{\Lambda_{31}}(E, e_{31})$, which satisfies

\[ G_{\Lambda_{31}}(E, e_{31}) \Lambda_{31}(E - H_0 - v_{31}(e_{31}))\Lambda_{31} = \Lambda_{31}. \]  

(A5)

This can be easily constructed through...
by using the two-body Green function $G_{v_31}(E, e_{31}) = [E - H_0 - V_{RGM}(e_{31}) + i0]^{-1}$ in the three-body space. The essential equation we need for deriving the full Green function $G_{31}(E, e_{31}) = [E - H_{31} - V_{RGM}(e_{31}) + i0]^{-1}$ is the decomposition

$$E - H_0 - V_{RGM}(e_{31}) = (E - e_0 - h_2 - \Delta m_3)$$

$$- \Lambda_{31}(E - e_{31} - h_2 - \Delta m_3) \Lambda_{31}$$

$$+ \Lambda_{31}[E - H_0 - V_{RGM}(e_{31})] \Lambda_{31}$$

$$= |u_{31}(E - e_{31} - h_2 - \Delta m_3)|^2$$

$$+ \Lambda_{31}[E - H_0 - V_{RGM}(e_{31})] \Lambda_{31},$$

(A7)

but the last equality is actually not satisfied since $(E - e_{31} - h_2 - \Delta m_3)$ is channel dependent. This difficulty is avoided by using $V_{RGM}^{mod}(\epsilon)$ in Eq. (16), in place of $V_{RGM}(e_{31})$. In fact, we find that

$$E - e_0 - h_2 - \Delta m_3 - (h_{31} - h_0) + (\Delta E_{int} + \Delta \epsilon + \Delta h_0)$$

$$= E - e_{31} - (h_2)$$

(A8)

is channel independent. Here $e_0$ and $(h_{31} - h_0)$ terms appear since $|u_{31}(E - e_{31} - h_2 - \Delta m_3)|^2$ actually satisfies Eq. (10) and not Eq. (3). This makes it possible to derive our basic relationship

$$G_0(E)T_{31}(E, \epsilon) = G_{31}(E, \epsilon)V_{RGM}^{mod}(\epsilon)$$

$$= G_{31}(E, \epsilon) - |u_{31}(E - e_{31})|$$

$$\cdot \frac{1}{(E - e_{31} - (h_2) - \Delta m_3)}$$

$$\times \frac{1}{(E - e_{31} - (h_2)}$$

$$\times (u_{31}(E - H_0),$$

(A9)

where all the kernels are defined by using $V_{RGM}^{mod}(\epsilon)$. From Eq. (A9) we can easily prove the second Faddeev equation, Eq. (6), and the orthogonality condition, Eq. (11).