Effect of a realistic three-body force on the spectra of medium-mass hypernuclei

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Abstract

We adopt the Hartree–Fock method in the proton–neutron–Λ formalism and the nucleon–Λ Tamm–Dancoff approximation to study the energy spectra of medium-mass hypernuclei. The formalism is developed for a potential derived from effective field theories which includes explicitly the 3-body NNN forces plus the YN LO potential. The energy spectra of selected medium-mass hypernuclei are presented and their properties discussed. The present calculation is the first step of a project devoted to ab initio studies of hypernuclei in medium and heavy mass regions. This may provide a guide for a better understanding of the YN interactions at momentum scales not accessible in few-body hypernuclei.

Keywords: many-body models, mean-field approximation, three-body force, hypernuclei

(Some figures may appear in colour only in the online journal)

1. Introduction

Hypernuclei have been studied within many theoretical approaches. Let us mention some of them. Faddeev or Yakubovsky equations based on realistic baryon–baryon interactions have been used to study the $A = 3$ and $A = 4$ systems [1]. Ab initio calculations of spectra of light hypernuclei up to the $p$-shell have been performed within no-core–shell model [2, 3] and fermionic molecular dynamics [4]. Medium-mass and heavy hypernuclei have been studied mostly within phenomenological models. The list includes calculations performed within the antisymmetrized molecular dynamics [5], the shell model [6], relativistic [7] and Skyrme [8, 9] mean field theories.

There are very few theoretical approaches which use realistic interaction to describe the hypernuclei above the $p$-shell. Among them, let us mention the auxiliary field diffusion Monte Carlo method [10, 11] which adopts local baryonic forces and the hypernuclear mean-field model using realistic nucleon–nucleon (NN) and nucleon–Λ ($N\Lambda$) interactions [12, 13].

The first requirement for a mean field approach is that the calculation must provide a satisfactory description of the density distributions, root mean square (rms) charge radii, and single-particle energies of the nucleons in the nuclear cores. In fact, the single-particle energies of $\Lambda$ in hypernuclei are mainly affected by the nuclear core rms radii and the density distributions. This is a consequence of the Bertlmann–Martin inequalities which link the splitting of the $\Lambda$ 0s and 0p single-particle levels with the nuclear rms radii [14]. On the other hand, HF is not able to reproduce the correct nuclear binding energies. Ground state correlation are to be included for this purpose [15].

The HF calculations using only realistic 2-nucleon NN interactions underestimate the nuclear radii [16] and overestimate the relative gaps among the nucleon single-particle levels [17]. Their correct description was partly achieved by including a phenomenological repulsive density dependent (DD) term simulating the NNN interaction [16, 17]. In our previous work [13], in fact, such a DD force was added to the optimized chiral NN potential NNLO$_\text{opt}$ [18] to generate the hypernuclear mean field.

In order to avoid free phenomenological parameters it is desirable to use a 3-body NNN interaction deduced from effective field theories. Such a term is included in the chiral NN + NNN NNLO$_\text{opt}$ potential derived recently [19]. This
interaction was optimized so as to reproduce the low-energy NN scattering data as well as binding energies and radii of selected nuclei up to oxygen and carbon isotopes [19]. It was used successfully in a recent HF + random phase approximation calculation of nuclear multipole resonances [20] and in a calculation based on the self-consistent Green’s function approach (SCGF) to study the potential bubble nucleus $^{34}$Si [21].

In this paper we adopt the above potential within the HF method in the proton–neutron–$\Lambda$ formalism to generate the mean field for hypernuclei with an even–even nuclear cores.

Additionally, we develop the nucleon–$\Lambda$ Tamm–Dancoff approximation (NA TDA) to describe hypernuclei consisting of one $\Lambda$ coupled to the even–odd or odd–even nuclear cores. The 3-nucleon NNN force is included in the derivation of the NA TDA only at the 2-body level of its normal-ordered form. The residual 3-body term gives zero contribution and is, therefore, ignored.

In this paper, we first implement the chiral NN + NNN NNLOsat potential as force acting among nucleons, and we study the influence of its 3-nucleon NN term on the rms charge radii, and nucleon single-particle energies of the nuclei $^{16}$O, $^{40}$Ca.

We investigate the effect of the NNN force on the single-particle energies of $\Lambda$ in $^{16}$O and $^{40}$Ca and on the energy spectra of the $^{16}$O, $^{40}$K. The interaction of $\Lambda$ with nucleons is described by the NA–$\Lambda$ channel of the chiral YN LO potential [22]. We do not take into account the influence of the $\Lambda$–$\Sigma$ mixing. This issue is discussed in the final section of the text.

The methods presented here are the first step of our project. Our long-term goal is to go beyond the mean-field approximation by including more complex many-body configurations.

This goal will be achieved by a straightforward extension to hypernuclei of the equation of motion phonon method (EMPm) developed for nuclear structure studies [23] and adopted extensively for light and heavy even–even [24–26] and medium-mass odd–even nuclei [27–29].

2. Theoretical formalism

The adopted Hamiltonian has the structure

$$H = T + V_{NN} + V_{NNN} + V_{NA} - T_{CM},$$

where $T$ is the kinetic energy operators of nucleons and $\Lambda$, $V_{NN}$, $V_{NNN}$, and $V_{NA}$ stand for the 2-body NN, NA, and the 3-body NNN interactions. The term $T_{CM}$ is the center-of-mass kinetic operator

$$T_{CM} = \frac{1}{2!(A - 1)m_N + m_\Lambda} \left( \sum_{a=1}^A \vec{p}_a^2 + 2 \sum_{a<c} \vec{p}_a \cdot \vec{p}_c \right),$$

where $A$ is the baryon number, $\vec{p}_a$ the momentum operator of the $a$th particle, $m_N \approx 938$ MeV and $m_\Lambda \approx 1116$ MeV are, respectively, the masses of nucleon and $\Lambda$.

In the second quantization formalism, the Hamiltonian (1) becomes

$$H = \sum_{ij} t_{ij}^{\Lambda} a_i^\dagger a_j + \sum_{ij} t_{ij}^N b_i^\dagger b_j + \sum_{ij} t_{ij}^{\nu} c_i^\dagger c_j + \frac{1}{4} \sum_{ijkl} V_{ijkl}^{NN} a_i^{\dagger} a_j^{\dagger} a_k b_l + \frac{1}{4} \sum_{ijkl} V_{ijkl}^{NN} b_j^{\dagger} b_k^{\dagger} b_i b_l$$

$$+ \sum_{ijkl} V_{ijkl}^{NNN} a_i^{\dagger} b_j^{\dagger} b_k^{\dagger} a_l b_l$$

$$+ \frac{1}{36} \sum_{ijklmn} V_{ijklmn}^{NNN} a_i^{\dagger} a_j^{\dagger} a_k^{\dagger} b_m b_n + \frac{1}{36} \sum_{ijklmn} V_{ijklmn}^{NNN} b_j^{\dagger} b_k^{\dagger} b_l^{\dagger} a_m a_l$$

$$+ \frac{1}{4} \sum_{ijkl} V_{ijkl}^{\nu} a_i^{\dagger} b_j^{\dagger} b_k^{\dagger} a_l b_l$$

$$+ \frac{1}{4} \sum_{ijkl} V_{ijkl}^{\nu} b_j^{\dagger} b_k^{\dagger} b_l^{\dagger} b_i b_i,$$

where $a_i^{\dagger}$, $b_j^{\dagger}$ and $c_i^{\dagger}$ are, respectively, the creation (annihilation) operators for protons, neutrons, and $\Lambda$. The hypernuclear mean field is generated by solving the self-consistent HF equations

$$t_{ij}^{\nu} = \sum_{kl} V_{ijkl}^{NN} P_{lj}^{\nu} + \sum_{kl} V_{ijkl}^{NN} P_{lj}^{\nu} + \sum_{kl} V_{ijkl}^{\nu} P_{lj}^{\nu}$$

$$+ \frac{1}{2} \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\nu} + \frac{1}{2} \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\nu}$$

$$+ \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\nu} = \delta_i^{\nu} \delta_j^{\nu},$$

$$t_{ij}^{N} = \sum_{kl} V_{ijkl}^{NN} P_{lj}^{N} + \sum_{kl} V_{ijkl}^{NN} P_{lj}^{N} + \sum_{kl} V_{ijkl}^{\nu} P_{lj}^{\nu}$$

$$+ \frac{1}{2} \sum_{klmn} V_{klmn}^{NN} P_{lk}^{N} + \frac{1}{2} \sum_{klmn} V_{klmn}^{NN} P_{lk}^{N}$$

$$+ \sum_{klmn} V_{klmn}^{NN} P_{lk}^{N} = \epsilon_i^{N} \delta_j^{N},$$

$$t_{ij}^{\Lambda} = \sum_{kl} V_{ijkl}^{NN} P_{lj}^{\Lambda} + \sum_{kl} V_{ijkl}^{NN} P_{lj}^{\Lambda} + \sum_{kl} V_{ijkl}^{\nu} P_{lj}^{\nu}$$

$$+ \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\Lambda} + \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\Lambda}$$

$$+ \sum_{klmn} V_{klmn}^{NN} P_{lk}^{\Lambda} = \epsilon_i^{\Lambda} \delta_j^{\Lambda},$$

where $\rho_{mn}^{\nu} = \langle HF | a_m^{\dagger} a_n | HF \rangle$, $\rho_{mn}^{N} = \langle HF | b_m^{\dagger} b_n | HF \rangle$ and $\rho_{mn}^{\Lambda} = \langle HF | c_m^{\dagger} c_n | HF \rangle$ are the proton, neutron and $\Lambda$ densities, respectively, and $\epsilon_i^{N}$, $\epsilon_i^{\Lambda}$, $\epsilon_i^{\nu}$ the single-particle energies.

The solution of the HF equations yields the single-particle energies of $\Lambda$ bound in the even–even nuclear core.

The states of the $\Lambda$ bound in the odd–even or even–odd core are determined in the NA TDA. We first solve the HF equations (4), (5) for the nuclear core and, subsequently, solve equation (6) to obtain the self-consistent single-particle energies $\epsilon_i^{\nu}$ and basis states for $\Lambda$. We, then, introduce the proton-$\Lambda$ and neutron-$\Lambda$ phonon operators

$$R_i^{\nu,p\Lambda} = \sum_{ph} \rho_{ph}^{\nu,p\Lambda} c_i^{\dagger} a_h,$$

$$R_i^{\nu,n\Lambda} = \sum_{ph} \rho_{ph}^{\nu,n\Lambda} c_i^{\dagger} b_h,$$
where \( a_{nljm} = a_{nljm} \) and \( b_{nljm} = b_{nljm} \). These operators represent a general superposition of the \( \Lambda \)-particle (p) nucleon-hole (h) configurations. The coefficients \( r_{ph}^{p,\Lambda} \) are fixed by solving the following eigenvalue equations

\[
\sum_{ph} ((E_p^{p,\Lambda} - E_h^{p,\Lambda})^2 \delta_{pp'} \delta_{hh'}) r_{ph}^{p,\Lambda} = (E_p^{p,\Lambda} - E_h^{p,\Lambda}) r_{ph}^{p,\Lambda},
\]

and

\[
\sum_{ph} ((E_p^{n,\Lambda} - E_h^{n,\Lambda})^2 \delta_{pp'} \delta_{hh'}) r_{ph}^{n,\Lambda} = (E_p^{n,\Lambda} - E_h^{n,\Lambda}) r_{ph}^{n,\Lambda}.
\]

### 3. Calculation details and results

In order to solve the HF eigenvalue problem we represent the Hamiltonian \((\mathbf{3})\) in the harmonic oscillator (HO) basis. The nucleon and \( \Lambda \) HO states depend on the oscillator lengths \( b_N \) and \( b_L \) which are fixed by the frequency \( \omega_{\text{HO}} \) through the relation

\[
b_{N(\Lambda)} = \frac{\hbar}{\sqrt{m_{N(\Lambda)} \omega_{\text{HO}}}}.
\]

The HF solutions do not depend on the parameter \( \hbar \omega_{\text{HO}} \) if the basis is large enough. We have put \( \hbar \omega_{\text{HO}} = 16 \text{ MeV} \). As we shall see, we reach a good convergence by including all two-body and three-body matrix elements under the restrictions \( \{ij\}: 2n_i + l_i + 2n_j + l_j \leq N_{\text{max}} \) and \( \{ijk\}: 2n_i + l_i + 2n_j + l_j + 2n_k + l_k \leq N_{\text{max}} \), with \( N_{\text{max}} = 12 \).

#### 3.1. HF investigation of \( ^{16}\text{O} \) and \( ^{40}\text{Ca} \) cores

The radial density distributions, charge radii, nucleon single-particle energies, and binding energies of \( ^{16}\text{O} \), \( ^{40}\text{Ca} \) are calculated by solving the HF equations \((4), (5)\) for nucleons only (\( \rho^\Lambda = 0 \)).

As shown in figure 1 the radial density distribution in \( ^{40}\text{Ca} \) converges with \( N_{\text{max}} \) more slowly than in \( ^{16}\text{O} \). In both nuclei, however, a saturation value is reached. A good convergence is reached also for the neutron single-particle energies \( e_i^n \) (figure 2).

In order to study the effect of the 3-nucleon NNN component of the NNLOsat potential on the rms charge radii and neutron single-particle energies we perform calculations with and without this term.

The mean square charged radius is given by

\[
\langle r_{ch}^2 \rangle = \left( 1 - \frac{1}{4A} \right) \frac{1}{Z} \langle r_p^2 \rangle + R_p^2 + \frac{N}{Z} R_n^2 + \frac{3\hbar^2}{4m_p^2 c^2}.
\]

where \( \langle r_p^2 \rangle = \int dr r^4 \rho_p^2(r) \) is the mean-square proton point radius, \( R_p = 0.8775(51) \text{ fm} \), \( R_n = -0.1149(27) \text{ fm}^2 \), and \( \frac{3\hbar^2}{4m_p^2 c^2} = 0.033 \text{ fm}^2 \) [19].

As shown in figure 1, the charge radii \( r_{ch} \) calculated with the 2-nucleon NN interaction only are too small. They are enhanced, in much better agreement with the experiments, once the NNN force is included. It is worth to notice that the phenomenological DD term adopted in our previous study produces a similar effect [13].
The 3-body NNN force has also the very important effect of reducing the separation between the neutron single-particle levels, as clearly illustrated in figure 3 for $^{16}$O and $^{40}$Ca. An analogous effect was produced by the phenomenological DD term \([13]\).

The effect of the 3-body NNN force on nuclear radii and nucleon single-particle energies is caused by its repulsive character. This term, in fact, reduces strongly the binding energies of $^{16}$O and $^{40}$Ca which are underestimated by a factor $\sim 3$ and $\sim 3.7$ respectively (table 2), indicating that ground state correlations are needed.

The strong impact on the physical observables induced by more complex configurations emerges from the calculations using the NNLO$_{sat}$ potential within a coupled cluster \([19]\) and a SCGF approach \([21]\).

The effect of the correlations has been also studied within the EMPM using the NNLO$_{opt}$ potential \([15]\).

This calculation yielded a contribution to the binding energy coming from two-phonon correlations comparable to the HF contribution. The impact on the charge radius was very modest instead.

On the ground of this EMPM calculation and the results obtained in \([19, 21]\) with NNLO$_{sat}$, we expect that the ground state correlations estimated within the EMPM should counterbalance the contribution coming from the strongly repulsive NNN part of the NNLO$_{sat}$ potential thereby filling to a large extent the gap with experiments.

### 3.2. Spectra of $^{16}$O, $^{17}$O, $^{40}$K and $^{41}$Ca hypernuclei

The hypernuclear energy spectra are affected by the strong dependence of the YN LO potential on the regulator cutoff parameter $\lambda$. We put $\lambda = 550$ MeV. This value yields a more bound $\Lambda$ particle with the energies closer to the empirical values. In general, bigger cutoff $\lambda$ causes an overall upward shift of the energy levels (figure 4). However, the separations among single-particle levels remain relatively stable.

Figures 5 and 6 point out the good convergence of the energy spectra with respect to the parameter $N_{max}$ of the studied hypernuclei.

As shown in figure 7, the energy gaps among the $\Lambda$ single-particle levels in $^{16}$O and $^{40}$K are significantly reduced once the NNN interaction is included. The relative distances among $s$-, $p$-, and $sd$-major shells are in fair agreement with the empirical levels. The whole spectra, however, are shifted...
upward with respect to the empirical one due to the underestimation of the binding energies. Clearly, more complex configurations need to be included.

The NNN interaction affects strongly also the relative energy spectra in $^{16}_{\Lambda}$O and $^{40}_{\Lambda}$Ca. As illustrated in figure 8, it reduces strongly the gaps among the levels in $^{16}_{\Lambda}$O in very close agreement with the experiments [34] and an analogous reduction of the separation energies is obtained for $^{40}_{\Lambda}$K (figure 9) where no experimental data are available.
Figure 9. The relative energies of $^{40}\Lambda K$ calculated with NN, and NN + NNN forces. The red (blue) lines represent the states with the negative (positive) parity.

4. Summary and outlook

The results presented here emphasize the impact of the 3-nucleon NNN component of the chiral NNLOsat potential on nuclei and hypernuclei.

In nuclei, in particular $^{16}\text{O}$ and $^{40}\text{Ca}$, the NNN force improves the description of the radial densities by flattening their distributions, enhances the nuclear radii in agreement with the empirical ones, and yields nucleon single-particle energies in close correspondence with the empirical levels. However, its strong repulsive character brings the theoretical binding energies far from the experimental data.

In hypernuclei, we obtain the $\Lambda$ single-particle levels of $^{17}_{\Lambda}\text{O}$, $^{41}_{\Lambda}\text{Ca}$ with realistic energy gaps and in close correspondence with the observed levels by an appropriate rigid translation of the whole spectra. This shift is very sensitive to the regulator cutoff $\lambda$ of the NY LO potential. This potential grows linearly with $\lambda$ and, therefore, affects little the relative energies of the hypernuclear spectra studied here. The relative energies of $^{16}_{\Lambda}\text{O}$ and $^{40}_{\Lambda}\text{K}$ also are in better agreement with the available experimental data once the NNN force is included.

To fill the gap with the experimental nuclear binding energies and to reproduce the absolute energies of hypernuclei, for the adopted value of the regulator cutoff $\lambda$, it is necessary to include more complex excitations.

We plan to study the effect of the ground state correlations and the coupling of the $\Lambda$-particle to many particle-hole excitations of the core resorting to the EMPM already mentioned in the introduction.

This method constructs and solves a set of equations of motion to generate an orthonormal multiphonon basis built of TDA phonons. The solution of the eigenvalue equation in the multiphonon basis so constructed yields highly correlated states, including the ground state. It can be implemented for any realistic Hamiltonian, does not rely on any approximation and takes the Pauli principle into full account.

We extend the method to hypernuclei with even–even and odd–even nuclear cores. In the two cases we couple, respectively, the $\Lambda$ states and the NA TDA phonons to the multiphonon excitations of the nuclear cores.

We intend to apply this theoretical framework to the hypernuclei $^{4}\Lambda K$, and $^{8}\Lambda K$, whose production is being planned at JLab [35]. Our calculations can provide the hypernuclear wave functions needed for the theoretical analysis of this production.

Another important step is the inclusion of the $\Lambda-\Sigma$ mixing in the $YN$ interaction. This will be accounted for by following the procedure of [36] which incorporates the NA–N$\Sigma$ part of the chiral LO $YN$ interaction into the NA–NA channel. The procedure is based on the SRG transformation [37] in the Wegner’s formulation [38], which leads to the suppression of the $\Lambda-\Sigma$ conversion terms in the Hamiltonian.

The SRG transformed interaction must include the 2-body $YN$ as well as the 3-body $YN$. The SRG transformed $YN$ force alone overestimates by several MeV the binding of $\Lambda$ in the $p$-shell hypernucliei [36] with respect to the experimental values. This over-binding is reduced to a large extent by the strongly repulsive SRG induced $YN$ term. We expect an analogous behavior also in heavier hypernucliei. In general, it would be also worth to implement other realistic $YN$ interactions like the chiral next-to-leading order $YN$ interaction [39].

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