Recent progress on hypernuclei

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Abstract. Some of last year’s progress made in hypernuclear physics is reviewed as follows: (i) resolving the $^5\Lambda$He overbinding problem in single-$\Lambda$ hypernuclei [1]; (ii) arguing that the onset of binding double-$\Lambda$ hypernuclei is most likely at $A=5$, with the neutral systems $^3\Lambda\Lambda n$ and $^4\Lambda\Lambda n$ unbound by a large margin [2]; and (iii) revising the calculated value of the loosely bound $^3\Lambda H$ lifetime to a level of $\sim 20\%$ shorter than the free $\Lambda$ lifetime [3], given recent claims from relativistic heavy ion experiments that $\tau(^3\Lambda H)$ is shorter than $\tau^\Lambda$ by as much as $\approx (30\pm8)\%$. Also discussed briefly in this context is the lifetime expected for the questionable $^3\Lambda n$ hypernucleus.

1. Introduction

Single- and double-$\Lambda$ hypernuclei provide a unique extension of nuclear physics into strange hadronic matter [4]. Experimental data on $\Lambda$ and $\Lambda\Lambda$ hypernuclei are unfortunately poorer both in quantity and quality than the data available on normal nuclei. Nevertheless, the few dozen $\Lambda$ separation energies $B_\Lambda$ of single-$\Lambda$ hypernuclei ($^A_\Lambda Z$) determined across the periodic table from $A=3$ to 208, and the three $\Lambda\Lambda$ hypernuclei ($^A_\Lambda\Lambda Z$) firmly established so far, provide a useful testground for the role of strangeness in dense hadronic matter, say in neutron star matter. Particularly meaningful tests of hyperon-nucleon and hyperon-hyperon strong-interaction models are possible in light $\Lambda$ and $\Lambda\Lambda$ hypernuclei, $A \leq 6$, which the three topics reviewed below are concerned with. Before focusing on each topic separately, I discuss briefly in the next section the pionless EFT ($\pi$EFT) framework which is used in two of these topics.

2. $\pi$EFT methodology

The leading-order (LO) $\pi$EFT interaction for aggregates of nucleons and $\Lambda$ hyperons consists of two-body and three-body $s$-wave contact interaction terms shown diagrammatically, together with the corresponding low-energy constants (LECs) listed alongside, in Fig. 1. Each LEC is labelled by the total Pauli-spin $S$ and isospin $I$ involved. Further contact terms, such as a three-$\Lambda\Lambda\Lambda$ term, appear only at subleading orders. These two-body and three-body contact interaction terms give rise to two-body and three-body potentials

$$V_2 = \sum_{IS} C_{\Lambda S}^{(I)} \sum_{i<j} \mathcal{P}_{IS}(ij)\delta_\lambda(\vec{r}_{ij}), \quad V_3 = \sum_{aIbS} D_{a\lambda b\lambda}^{(I)} \sum_{i<j<k} \mathcal{Q}_{IS}(ijk) \left( \sum_{\text{cyc}} \delta_\lambda(\vec{r}_{ij})\delta_\lambda(\vec{r}_{jk}) \right), \quad (1)$$

where $\mathcal{P}_{IS}$ project on $s$-wave $NN, \Lambda N, \Lambda\Lambda$ pairs with isospin $I$ and spin $S$ values associated with two-body LECs in Fig. 1. These LECs are fitted to low-energy two-body observables, e.g., to the corresponding $NN, \Lambda N, \Lambda\Lambda$ scattering lengths. Similarly, $\mathcal{Q}_{IS}$ project on $NN\Lambda, NNN$
Figure 1. Diagrammatic presentation of two-body (left) and three-body (right) contact terms, and their associated LEC input \( (C_1, \ldots, C_5) \) & \( (D_1, \ldots, D_5) \) to a LO \( \pi \)EFT calculation of light nuclei (upper), \( \Lambda \) hypernuclei (middle) and \( \Lambda \Lambda \) hypernuclei (lower), with values of spin \( S \) and isospin \( I \) corresponding to \( s \)-wave configurations. Figure adapted from Ref. [2].

\[ \delta_\lambda(\vec{r}) = \left( \frac{\lambda}{2\sqrt{\pi}} \right)^3 \exp\left( -\frac{\lambda^2}{4} \vec{r}^2 \right), \]  

whereby smearing a zero-range (in the limit \( \lambda \rightarrow \infty \)) Dirac \( \delta^{(3)}(\vec{r}) \) contact term over distances \( \sim \lambda^{-1} \). The cutoff parameter \( \lambda \) may be viewed as a scale parameter with respect to typical values of momenta \( Q \). To make observables cutoff independent, the LECs must be properly renormalized. Truncating \( \pi \)EFT at LO and using values of \( \lambda \) higher than the breakup scale of the theory which is of order \( 2m_\pi \) for the isoscalar \( \Lambda \) hyperon, observables acquire a residual dependence \( O(Q/\lambda) \) which diminishes with increasing \( \lambda \). Using such two-body \( V_2 \) and three-body \( V_3 \) regularized contact interaction terms, the \( A \)-body Schrödinger equation was solved by expanding the wave function \( \Psi \) in a correlated Gaussian basis, using the stochastic variational method (SVM).

Few-body \( \pi \)EFT calculations were first reported for nucleons in Refs. [5, 6] and recently extended to lattice nuclei [7–10]. Past hypernuclear applications are limited to \( \Lambda \) hypernuclei [11–13] and to \( A=4,6 \) \( \Lambda \)-core three-body calculations [14, 15]. The calculations reviewed below are the first systematic single- and double-\( \Lambda \) hypernuclear studies covering the full nuclear \( s \) shell.
3. Overbinding of $^5\Lambda$He

The overbinding of $^5\Lambda$He upon using fitted two-body $\Lambda N$ interactions, even when adding $\Lambda NN$ terms owing to $\Sigma$ hyperon excitation, was first recognized and stated clearly in a 1972 landmark paper by Dalitz et al. [16]. There, as well as in recent LO chiral effective field theory ($\chi$EFT) calculations [17], the $\Lambda$ separation energy $B_{\Lambda}(^5\Lambda$He) comes out as large as 6 MeV, well above the value $B_N^{\exp}(^5\Lambda$He)=$3.12\pm0.02$ MeV, as demonstrated in the first two main rows of Table 1. No truly ab-initio calculations of $B_{\Lambda}(^5\Lambda$He) using next-to-leading-order (NLO) $\chi$EFT interactions have ever been reported. Comprehensive NLO calculations for $A=3,4$ hypernuclei have recently been published [22], with results listed in the last two rows of the table. The cutoffs 500, 650 MeV chosen for the 2013 and 2019 versions, respectively, are motivated by looking for those cutoff values that correspond to $\Lambda$ well-depth values in the range $D_{\Lambda} = 28 - 30$ MeV. Given that such NLO versions fit the low-energy $\Delta p$ cross sections [23] better than the LO model [24] does, it is puzzling why the latest 2019 NLO version does so poorly for the $A=4$ hypernuclei, definitely worse than the LO calculation [18] does.

Table 1. Ground-state $\Lambda$ separation energies $B_{\Lambda}$ and excitation energies $E_x$ (in MeV) from several few-body calculations of $s$-shell $\Lambda$ hypernuclei, see text. Charge symmetry breaking is included in the $^4$He results from Ref. [18].

| Exp. | $B_{\Lambda}(^4\Lambda$H) | $B_{\Lambda}(^4\Lambda$He_{g.s}) | $E_x(^4\Lambda$He_{exc}) | $B_{\Lambda}(^5\Lambda$He) |
|------|-----------------|-----------------|-----------------|-----------------|
| $\chi$EFT(LO600) | 0.11(5) [19] | 2.44(3) [18] | 1.278 [18] | 5.82(2) [17] |
| $\chi$EFT(LO700) | – | 2.423 [18] | 1.941 [18] | 4.43(2) [17] |
| $\chi$EFT(NLO13500) | 0.135 [22] | 1.705 [22] | 0.915 [22] | – |
| $\chi$EFT(NLO19650) | 0.095 [22] | 1.530 [22] | 0.614 [22] | – |

Here I review a rather successful attempt to resolve the $^5\Lambda$He overbinding problem within the simpler EFT approach of pionless EFT ($g$EFT), limited at LO to nucleons and $\Lambda$-hyperons degrees of freedom, by means of precise SVM calculations of $s$-shell hypernuclei [1]. Note that the long-range $\Lambda N \to \Sigma N$ one-pion exchange (OPE) transition followed by an equally long-range $\Sigma N \to \Lambda N$ OPE transition is dominated by its central $S \to D \to S$ two-pion exchange component, which is partially absorbed in the $\Lambda N$ and $\Lambda NN$ LO contact LECs. Short-range $K$ and $K^*$ meson exchanges induce a mild $\Lambda N$ tensor force [25,26], the weakness of which is confirmed in shell-model observables of $p$-shell $\Lambda$ hypernuclear spectra [27]. Such momentum-dependent interaction terms which appear at subleading order in $g$EFT power counting, need to be introduced systematically in future applications to $p$-shell hypernuclei.

Apart from the two-body contact terms that are specified here by $NN$ and $\Lambda N$ spin-singlet and triplet scattering lengths, amounting to four low-energy constants (LECs), the theory uses additionally four three-body LECs: a pure $NNN$ LEC fitted to $B(^4\Lambda$H) and three $\Lambda NN$ LECs associated with the three possible $s$-wave $\Lambda NN$ systems, of which only $^4\Lambda$H($I=0, J^P=\frac{1}{2}^+$) is bound. Therefore, on top of fitting its binding energy, the binding energies of $^4\Lambda$H_{g.s.}($I=\frac{1}{2}, J^P=0^+$) and of $^4\Lambda$H_{exc.}($I=\frac{1}{2}, J^P=1^+$) are also fitted. The fitted LECs are used then, for a sequence of $\lambda$ cutoff values, to evaluate the binding energies of $^4$He and $^5\Lambda$He. Remarkably, $B(^4\Lambda$H) is reproduced well in the renormalization scale invariance limit $\lambda \to \infty$.

The $g$EFT approach was applied in SVM few-body calculations of $s$-shell hypernuclei, using several models of the $\Lambda N$ scattering lengths. The resulting $\Lambda$ separation energy values $B_{\Lambda}(^5\Lambda$He) are shown in Fig. 2 for two such models as a function of the cutoff $\lambda$. Common to all $\Lambda N$
models, the calculated $B_{\Lambda}^{\Lambda}(^5\text{He})$ values switch from about 2–3 MeV overbinding at $\lambda=1$ fm$^{-1}$ to less than 1 MeV underbinding between $\lambda=2$ and 3 fm$^{-1}$, and smoothly varying beyond, approaching a finite (renormalization scale invariance) limit at $\lambda \to \infty$. A reasonable choice of finite cutoff values in the present case is between $\lambda \approx 1.5$ fm$^{-1}$, which marks the $\pi$EFT breakup scale of $2m_\pi$, and 4 fm$^{-1}$, beginning at which the detailed dynamics of vector-meson exchanges may require attention. We note that for $\lambda \gtrsim 1.5$ fm$^{-1}$ all of the three $\Lambda NN$ state components are repulsive, as required to avoid Thomas collapse. Recent LO $\chi$EFT calculations [28] using induced $YNN$ repulsive contributions suggest that the $s$-shell overbinding problem extends to the $p$ shell. Interestingly, shell-model studies [27] reproduce satisfactorily $p$-shell ground-state $B_{\Lambda}$ values, essentially by using $B_{\Lambda}^{\text{exp}}(^5\text{He})$ for input, except for the relatively large difference of about 1.8 MeV between $B_{\Lambda}(^6\text{Li})$ and $B_{\Lambda}(^9\text{Be})$. In fact, it was noted long ago that strongly repulsive $\Lambda NN$ terms could settle it [29]. It would be interesting to test the $\Lambda NN$ interaction terms derived here in future shell-model studies and, perhaps, also in NS matter calculations such as by Lonardoni et al. [30] that are geared to resolve the ‘hyperon puzzle’ [31].

4. Onset of binding in $\Lambda\Lambda$ hypernuclei

The second topic reviewed here is the onset of binding $\Lambda\Lambda$ hypernuclei [2], using a methodology similar to that used for the previous topic. Reliable data on $\Lambda\Lambda$ hypernuclei are scarce: the Nagara event [32,33] is perhaps the only $\Lambda\Lambda$ hypernucleus determined unambiguously, identified as $^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He}$, with two more $\Lambda\Lambda$ hypernuclei, $^{\Lambda\Lambda}_{\Lambda\Lambda}\text{Be}$ and $^{13}_{\Lambda\Lambda}\text{B}$, that are also generally accepted [34]. The $^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He}$ datum $\Delta B_{\Lambda\Lambda}(^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He})=B_{\Lambda\Lambda}(^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He})-2B_{\Lambda}(^{3}_{\Lambda}\text{He})=0.67\pm0.17$ MeV [35] serves as a constraint, assuming also that the low-energy $\Lambda\Lambda$ interaction is weaker than the $\Lambda N$ interaction. The onset of $\Lambda\Lambda$ hypernuclear binding is then found at the isodoublet $^{\Lambda\Lambda}_{\Lambda\Lambda}\text{H}$–$^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He}$, with a $^{\Lambda\Lambda}_{\Lambda\Lambda}\text{H}$ bound state not definitively excluded.

A separation energy values $B_{\Lambda}(^{\Lambda\Lambda}_{\Lambda\Lambda}\text{H})$ from $\pi$EFT calculations [2] are shown in Fig. 5. Several representative values of the $\Lambda\Lambda$ scattering length were used, spanning a broad range of values suggested by analyses of $\Lambda\Lambda$ correlations observed recently in relativistic heavy-ion collisions and by analyzing the KEK-PS E522 [35] invariant mass spectrum in the reaction $^{12}\text{C}(K^-, K^+)^{\Lambda\Lambda}X$ near the $\Lambda\Lambda$ threshold; see Ref. [2] for detailed references. Here the choice of $a_{\Lambda\Lambda}$ determines the one $\Lambda\Lambda$ LEC required at LO, while the $\Lambda\Lambda N$ LEC was fitted to the $\Delta B_{\Lambda\Lambda}(^{\Lambda\Lambda}_{\Lambda\Lambda}\text{He})=0.67\pm0.17$ MeV datum [33]. Most calculations were made using the Alexander[B] $\Lambda N$ model with scattering lengths $a_{s,t}=-1.8,-1.6$ fm [23], but for cutoff $\lambda=4$ fm$^{-1}$ three
other $\Lambda N$ interaction models from Ref. [1] were also used, demonstrating that the $\Lambda N$ model dependence is rather weak when it comes to double-$\Lambda$ hypernuclei, provided $B_\Lambda$ values of single-$\Lambda$ hypernuclei for $A < 5$ are fitted to generate the necessary $\Lambda NN$ LECs. Calculated values of $B_\Lambda(\Lambda^5 \Lambda n)$ are also shown in the figure as a check. One observes that $\Lambda^5 \Lambda n$ comes out particle stable over a broad range of cutoff values used in the calculations. This is not the case for $\Lambda^4 \Lambda n$ which comes out unbound with respect to $\Lambda^3 \Lambda n$ for most of the permissible parameter space. Finally, ‘Tjon line’ correlations [36] found between $B_\Lambda(\Lambda^5 \Lambda n)$ and $B_\Lambda(\Lambda^6 \Lambda^\Lambda \Lambda n)$, when $\Delta B_{\Lambda \Lambda}(\Lambda^6 \Lambda^\Lambda \Lambda n)$ is varied within and also outside of its reported error-bar values, are demonstrated in the right panel of Fig. 3. Such correlations were noted already in old Faddeev calculations, e.g. Ref. [37].

Table 2. $\Lambda$ separation energies $B_\Lambda(\Lambda^A \Lambda^A n)$ for $A=3–6$, calculated using $a_{\Lambda \Lambda} = -0.8$ fm, cutoff $\lambda = 4$ fm$^{-1}$ and the Alexander[B] $\Lambda N$ interaction model [23], see text.

| Constraint (MeV) | $\Lambda^3 \Lambda n$ | $\Lambda^4 \Lambda n$ | $\Lambda^4 \Lambda^\Lambda \Lambda n$ | $\Lambda^5 \Lambda n$ | $\Lambda^6 \Lambda^\Lambda \Lambda n$ |
|-----------------|----------------------|----------------------|----------------------|----------------------|----------------------|
| $\Delta B_{\Lambda \Lambda}(\Lambda^6 \Lambda^\Lambda \Lambda n)$=0.67 | – | – | 1.21 | 3.28 |
| $B_\Lambda(\Lambda^4 \Lambda^\Lambda \Lambda n)$=0.05 | – | – | 0.05 | 2.28 | 4.76 |
| $B(\Lambda^4 \Lambda n)$=0.10 | – | 0.10 | 0.86 | 4.89 | 7.89 |
| $B(\Lambda^5 \Lambda n)$=0.10 | 0.10 | 15.15 | 18.40 | 22.13 | 25.66 |

To make $\Lambda^4 \Lambda n$ particle stable one may reduce the repulsive $\Lambda \Lambda N$ LEC from its value constrained by the $\Lambda^6 \Lambda^\Lambda \Lambda n$ datum, using representative values for $a_{\Lambda \Lambda}$ and the cutoff $\lambda$ for which $\Lambda^4 \Lambda n$ was found particle unstable. According to the first two rows in Table 2, this will overbind $\Lambda^6 \Lambda^\Lambda \Lambda n$ by $\approx 1.5$ MeV. Reducing further the $\Lambda \Lambda N$ LEC one binds the neutral systems, first $\Lambda^4 \Lambda n$ (third row) and then $\Lambda^5 \Lambda n$ (fourth row), at a price of overbinding further $\Lambda^6 \Lambda^\Lambda \Lambda n$. These results strongly suggest that the $A = 3, 4$ neutral $\Lambda \Lambda$ hypernuclei are unbound within a large margin.
5. $^3$H and $^3$He lifetime puzzles

The third topic discussed here is the hypertriton lifetime puzzle: why is the lifetime of the loosely bound $^3\Lambda$H shorter than the free-$\Lambda$ lifetime $\tau_\Lambda$ by as much as $\approx (30 \pm 8)\%$, as suggested in recent measurements using relativistic heavy ion collisions to produce light nuclei, anti-nuclei and hyperfragments \cite{38} such as $^3\Lambda$H? Also discussed briefly is the lifetime expected for the questionable $^3\Lambda$n \cite{39,40}. This 3-body hypernucleus has been found unanimously particle unstable in several few-body calculations cited below, including Ref. \cite{1} discussed in this report.

Measurements of the $^3\Lambda$H lifetime in emulsion or bubble-chamber experiments during the 1960s and early 1970s gave conflicting and puzzling results. Particularly troubling appeared a conference report by Block et al. claiming a lifetime of $\tau(3\Lambda H)=(95^{+19}_{-15})$ ps \cite{41}, to be compared with a free $\Lambda$ lifetime $\tau_\Lambda=(236^{+6}_{-6})$ ps measured in the same He chamber \cite{42}. Given the loose $\Lambda$ binding, $B_\Lambda(3\Lambda H)=0.13^{+0.05}_{-0.05}$ MeV, it was anticipated that $\tau(3\Lambda H)\approx \tau_\Lambda$, as argued by Rayet and Dalitz (RD) \cite{43} using a closure-approximation approach for $^3\Lambda$H.

Table 3. $^3\Lambda$H, $s\cdot(\frac{1}{2}^+)\cdot\gamma$ decay rate calculated in units of the free $\Lambda$ decay rate $\Gamma_\Lambda$ and listed in a year of publication order. The first row lists results for plane-wave pions, disregarding pion final state interaction (FSI) contributions which are listed in the second row. A calculated nonmesonic decay rate contribution of 0.017 from Ref. \cite{44} was added uniformly in obtaining the total decay rates listed in the third, last row.

| $\Gamma(3\Lambda H)$ model | 1966 | 1992 | 1998 | 2019 | 2019 |
|---------------------------|------|------|------|------|------|
| Without pion FSI          | 1.05 | 1.12 | 1.01 | 1.11 |
| Pion FSI contribution     | -0.013 | - | - | 0.11 |
| Total                     | 1.05 | 1.14 | 1.03 | 1.23 |

Table 3 lists $^3\Lambda$H, $s\cdot(\frac{1}{2}^+)\cdot\gamma$ decay rate values calculated by RD and in several subsequent solid calculations, all reaching similar results. Claims for large departures from the free $\Lambda$ value are, as a rule, incorrect or irreproducible. The RD methodology was also used by Congleton \cite{45} and by Gal and Garcilazo (GG) \cite{3}, with the latter one solving appropriate three-body Faddeev equations to produce a $^3\Lambda$H wavefunction. The Kamada et al. calculation \cite{46}, while also solving Faddeev equations for the $^3\Lambda$H wavefunction, accounted microscopically for the outgoing $3N$ phase space and FSI, thereby doing without a closure approximation. Pion FSI was considered only in two of these works, with differing results: (i) repulsion, weakly reducing $\Gamma(3\Lambda H)$ in RD; and (ii) attraction, moderately enhancing it in GG. The latter result is supported by the $\pi^-$-atom $1s$ level attractive shift observed in $^3$He \cite{47}. It is remarkable that $^3\Lambda$H decay is the only light hypernucleus decay where the low-energy pion $s$-wave FSI is expected to be attractive. The decays of $^4\Lambda$H, $^4$He and $^5$He involve pion-$^4$He FSI which is known from the $\pi^-$ atomic $^4$He $1s$ level shift to be repulsive \cite{48}.

Renewed interest in the $^3\Lambda$H lifetime problem arose by recent measurements of $\tau(3\Lambda H)$ in relativistic heavy ion experiments marked in Fig. 4 (STAR \cite{19}, HypHI \cite{50}, ALICE \cite{51}, STAR \cite{52} and ALICE \cite{53}) reporting values shorter by $(28\pm8)\%$ than $\tau_\Lambda=(263\pm2)$ ps \cite{38}. While enhancement of the free $\Lambda$ decay rate by up to $\approx 20\%$ is theoretically conceivable relying on the new GG calculation \cite{3} as recorded in the last column of Table 3, it appears inconceivable at present to reproduce a 30% or even larger enhancement suggested by some of the recent heavy-ion experiments. Note however that the most recent ALICE lifetime result \cite{53} is compatible within errors with the listed calculated values and also with $\tau_\Lambda$. 
Before closing this section I wish to make a few remarks on $^3\Lambda n$, conjectured by the HypHI GSI Collaboration [39] to be bound, while unbound in recent theoretical calculations [54–56]. In $^3\Lambda n$ decays induced by $\Lambda \to p + \pi^-$, where the $^3\Lambda n$ neutrons are spectators, the $^3\Lambda n$ weak decay rate is given in the closure approximation essentially by the $\Lambda \to p + \pi^-$ free-space weak-decay rate, whereas in $\Lambda \to n + \pi^0$ induced decays the production of a third low-momentum neutron is suppressed by the Pauli principle, and this $^3\Lambda n$ weak decay branch may be disregarded up to perhaps a few percents. Hence $\Gamma(^3\Lambda n)/\Gamma_\Lambda \approx 1.114 \times 0.641 = 0.714 \ [3]$, where the factor 1.114 follows from the difference between the recoil energies in the corresponding phase space factors, and the factor 0.641 is the free-space $\Lambda \to p + \pi^-$ fraction of the total $\Lambda \to N + \pi$ weak decay rate, giving rise to an estimated $^3\Lambda n$ lifetime of $\tau(^3\Lambda n) \approx 368$ ps, which should hold up to a few percent contribution from the $\pi^0$ decay branch. This lifetime is considerably longer than $181^{+30}_{-24} \pm 25$ ps or $190^{+37}_{-25} \pm 36$ ps deduced from the $n\pi^-$ and $\tau\pi^-$ alleged decay modes of $^3\Lambda n$ [39,40], providing a strong argument against the conjectured stability of $^3\Lambda n$.

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