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Author(s): Benjamin F. Gibson, T-16

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Hypernuclear Physics, a Brief Past and Bright Future

B.F. Gibson ^{a*}

^aTheoretical Division, Los Alamos National Laboratory
Los Alamos, New Mexico 87545, USA

From its cosmic ray origins in 1952, interest in hypernuclear physics has focused upon the differences between conventional nuclei and their hypernuclear counterparts. Explore here are some of the unique aspects of few-baryon systems with non-zero strangeness, including features of the hyperon-nucleon interaction which distinguish it from the nucleon-nucleon force. Examine is the question of whether our models can extrapolate beyond the realm of conventional nuclear physics in which they were constructed or whether they are merely exquisite tools for interpolation within the world of zero strangeness.

1. Introduction

A primary reason for investigating the structure and reactions of baryon systems is to achieve an understanding the fundamental baryon-baryon force in the realm of non-perturbative QCD. Few-baryon systems play an essential role, because one can calculate complete solutions to test a particular baryon-baryon interaction ansatz. Hypernuclei, exotic nuclei containing one or more hyperons ($Y = \Lambda, \Sigma$, or Ξ) are crucial to this investigation, because they permit one to probe models based upon our experience in the nonstrange sector; they lie outside of the conventional world where our models were developed. That is, we can test whether our sophisticated models of the nucleon-nucleon (NN) interaction extrapolate successfully beyond the zero strangeness region in which the parameters were determined, or whether the models merely interpolate.

The presence of the strangeness degree of freedom (flavor) adds a new dimension to our evolving picture of nuclear physics. We shall see that the physics of hypernuclei is both novel and puzzling, stretching our intuition and analysis capability beyond that developed during the more than half century that we have explored conventional nuclear physics. The hypernuclear sector of hadronic physics is not just a simple extension of zero-strangeness phenomena.

2. Our Strangeness-Zero Experience

In an attempt to relate nonperturbative QCD to physical observables, a number of theorists have turned to chiral perturbation theory and effective field theory. However, in the case of the NN interaction we have already available rather sophisticated meson-exchange and one-boson-exchange potentials which embody important characteristics of

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these approaches: a one-pion-exchange tail and a quantitative fit to the low-energy scattering parameters and deuteron properties. Moreover, thorough partial wave analysis of the NN scattering data have determined the number of free parameters required to represent a particular partial wave. Therefore, the potentials developed by the Argonne group and the Nijmegen group and the Bonn group provide superb models which satisfy the constraints of chiral perturbation theory or finite range effective field theory. These potential models form the basis for our successful “traditional approach” to calculating physical observables for few-nucleon systems with strangeness zero ($S = 0$).

The traditional approach to nuclear physics can most easily be defined in terms of the model assumptions:

- Nuclei consist only of nucleons – other degrees of freedom are suppressed.
- Nucleons move slowly within the nucleus – nonrelativistic dynamics prevails.
- Nucleons interact primarily via pairwise (two-body) forces.

This is an enormous simplification of the physics, but it does account amazingly well for much experimental data. Nonetheless, our ability to calculate has achieved the precision required to see differences between the traditional model predictions and experiment. Therefore, much of the research during the past decade has been focussed upon extensions: meson exchange currents (MEC), three-body forces (3BF), NN - $N\Delta$ coupling, relativistic dynamics, quark-gluon substructure, Such effects can be included to fine tune the model, but their contributions are often small and difficult to distinguish from one another.

It has now been reasonably established [1,2] that the low energy observables “scale” with the trinucleon binding energy. A summary of results for charge radii, wave function probabilities, magnetic moments, Coulomb energy, asymptotic normalization constants, and Nd scattering lengths can be found in Ref. [3]. Because of this scaling property, the triton discussion can be limited to results for the binding energy. Benchmark results exist for a variety of realistic potentials, where realistic implies

- strong spin-isospin dependence ($V_{nn} \neq V_{np}$),
- strong tensor force (OPEP is essential, providing up to 3/4 of the potential energy in 3H and 4He),
- strong short range repulsion (the probability of NN overlap at such separations should be small),

in addition to a reasonable fit to the NN scattering data. Charge-dependent potential models (with $V_{nn} \neq V_{np}^s \neq V_{np}^t$) have been used to estimate the triton binding energy to be $B(^3H) = 7.6$ MeV. That is, a *local* potential model, which fits the NN observables as well as a proper phase shift analysis implies is possible, leads to underbinding of the triton by about 0.85 MeV, and a corresponding failure to describe the low energy physical observables. This missing energy is less than 2% of the 50 MeV of potential energy in the system, and probably provides a more quantitative description of the triton than we had any right to expect *a priori*.

Such underbinding of the triton by local potential models led theorists to ask [4] about the role of three-nucleon forces. Adding a two-pion-exchange three-body force (3BF) to the Hamiltonian, adjusted to reproduce the triton binding energy, indeed scales the other physical observables into agreement with experiment. That is, a nonrelativistic Hamiltonian composed of a local NN potential plus a suitable 3BF can yield approximately the correct value for $B(^3\text{H})$. Moreover, it leads to the correct binding of ^4He (see, *e.g.*, the GFMC results of Carlson [5]), enhancing the binding by some 3 MeV, as predicted by the strong correlation among the ^3H and ^4He binding energies established by Tjon [6]. Similar results were later confirmed by Gloeckle and coworkers. Although the 3BF approach is but one means of achieving the desired increase in binding, it seems not unreasonable when such a model 3BF is found to contribute some 18 MeV to the binding of ^7Li and the nucleus is underbound by only some 2 MeV [7].

Finally, it should be noted that even though relativistic corrections have been estimated *a posteriori* to be small, the short range regularization of the potentials in vogue almost certainly ensures that such estimates are small. Moreover, by fitting the nonrelativistic potential to NN scattering data, one is unable to delineate what relativistic effects are already subsumed in the model.

3. Strangeness of Hypernuclei

As stated above the question of prime importance is whether our models, developed to describe conventional nuclei and nuclear reactions, extrapolate beyond the $S = 0$ realm? Or do they merely provide exquisite interpolation schemes? We shall see below that strange particle physics is not a simple extension of known $S = 0$ phenomena. $S \neq 0$ physics observations have been found to be new; the physics is different.

3.1. A Brief Chronology

Hypernuclei were first discovered in 1952 in a balloon-flown emulsion stack by Danysz and Pniewski [8]. Such cosmic ray observations of hypernuclei were followed by pion and proton beam production in emulsions and later in ^4He bubble chambers. The weak decay of the Λ into a $\pi^- + p$ was used to identify the Λ hypernucleus as well as to determine binding energies, spins, and lifetimes up to $A = 15$ [9–12]. By 1957 eight species had been observed, and in 1962 charge symmetry breaking in the Λ separation energy difference for the $A = 4$ mirror pair was confirmed. The first $\Lambda\Lambda$ hypernucleus was reported in 1963 [13], produced by Ξ^- capture and identified by the sequential π^- weak decay of both Λ s. In the late 1960s much of the Λp scattering data base was obtained using hydrogen bubble chambers [14–17].

The systematic investigation of hypernuclei occurred with the arrival of separated K^- beams, which were suited to counter techniques [18]. The (K^-, π^-) reaction produced information about excited states, particularly in the p shell. Particle-stable excited states in the $A = 4$ isodoublet were discovered [11,19] through observation of de-excitation γ s. In the 1980s interest turned to weak decays of Λ hypernuclei [20,21] and the possible existence of Σ hypernuclei [22]. In the 1990s interest focused upon “H” dibaryon [23] searches, the identification of single-particle states in heavy Λ hypernuclei [24], the identification of new dynamical symmetries [25], the observation of Ξ hypernuclei [26], and the search for strange matter.

3.2. Novel Aspects of $S = -1$ Physics

Despite the limited resources and manpower devoted to hypernuclear physics a number of features that differentiate $S = -1$ physics from conventional $S = 0$ physics have been discovered. The experimental observations show that hypernuclear physics is novel. In particular, one observes:

- anomalous binding energies
- a vanishing spin-orbit force
- significant three-body-force effects
- striking charge symmetry breaking
- new dynamical symmetries
- ground-state spin inversion
- puzzling nonmesonic weak decays
- anomalous π^+ mesonic decay

The $S = -1$ sector of hadronic physics is not just a simple extension of $S = 0$ phenomena.

3.3. Binding Energy Systematics

The available data for few-body Λ hypernuclei come primarily from emulsion experiments [9–11,19] – binding energies and weak decay properties. We restrict our discussion to binding energies, because the $S = 0$ sector has taught us that binding energies determine the low energy observables. In the study of hypernuclei, it is customary to quote the Λ -separation energy:

$$B_\Lambda(\Lambda A) = B(\Lambda A) - B(A-1).$$

This has the advantage of removing first order Coulomb effects in mirror pairs plus allowing one to focus upon the physics arising from the hyperon interaction.

In the $S = 0$ sector we observe that the ratio of neutron separation energies for neighboring s-shell nuclei is approximately 3:

- $B_n(^3H)/B_n(^2H) \simeq 6/2 = 3$,
- $B_n(^4He)/B_n(^3H) \simeq 20/6 \simeq 3$.

If the physics of few-body systems were similar, then we might anticipate a factor of 3 in the ratio of Λ separation energies for neighboring Λ hypernuclei. Using $B_\Lambda(^4H) \simeq 2$ MeV as our base, we would then predict:

- $B_\Lambda(^5He) \simeq 3 \times B_\Lambda(^4H) \simeq 6$ MeV,
- $B_\Lambda(^3H) \simeq \frac{1}{3} \times B_\Lambda(^4H) \simeq \frac{2}{3}$ MeV.

Simple, central force calculations using $\bar{V}_{N\Lambda}$ fitted to $B_\Lambda(^4\text{H})$ plus low-energy Λp scattering data confirm [27–31] this simple analysis.

However, such is not the case in nature. The real world is more complex. The separation energies for the s-shell systems are quoted in Table 1 along with the measured γ -ray de-excitation energies [19] for the two species with particle-stable excited states. The $A=6$ entry [32] is $B_{\Lambda\Lambda} = B(^6\text{He}) - B(^4\text{He})$. Thus, experimentally we observe that $B_\Lambda(^5\text{He}) \simeq$

Table 1
 Λ -separation energies and excitation energies in MeV.

hypernucleus	B_Λ	E_γ
$^3_\Lambda\text{H}$	0.13 ± 0.05	
$^4_\Lambda\text{H}$	2.04 ± 0.04	1.04 ± 0.04
$^4_\Lambda\text{He}$	2.39 ± 0.03	1.15 ± 0.04
$^5_\Lambda\text{He}$	3.10 ± 0.02	
$^6_{\Lambda\Lambda}\text{He}$	10.9	

3.1 MeV, only some 50% of the above anticipated value, and $B_\Lambda(^3_\Lambda\text{H}) \simeq 0.13$ MeV, only some 20% of the above anticipated value. Our $S = 0$ model experience does not naively extrapolate to $S = -1$. Even at first glance, one sees quickly that the physics of the s-shell hypernuclei is puzzling.

4. $S = -1$ Status Report

4.1. The Hyperon-Nucleon Interaction

Because the hyperon masses differ markedly from those of the neutron and proton, we know that $SU(3)$ symmetry is broken. Just how it is broken is a question of fundamental importance to our understanding of the baryon-baryon interaction in the nonperturbative realm of QCD. The investigation of strangeness in hypernuclei will play a significant role in our understanding of the strong force. For example, several features of the hyperon-nucleon (YN) interaction clearly come into play in hypernuclear physics. The Λ ($T = 0$) and N ($T = 1/2$) cannot exchange a π ($T = 1$) in first order, so that there is no dominant OPE tensor force in ΛN scattering. Shorter range properties of the baryon-baryon interaction play a more important role than in NN scattering. The longest range components are due to the exchange of two pions or one kaon. The shorter range K -exchange potential does admit a tensor-force component, but it is largely cancelled by that from K^* -exchange. Thus, tensor-force effects in the $N\Lambda$ interaction are expected to be smaller than those in the NN interaction [33–39]. On the other hand, explicit $N\Lambda - N\Sigma$ (octet-octet) mixing in the hyperon-nucleon sector appears to play a much larger role in hypernuclear physics than does $NN - N\Delta$ (octet-decuplet) mixing in the nonstrange sector. The $m_\Sigma - m_\Lambda$ mass difference is only some 80 MeV, and the width of the Σ is small compared to that of the Δ .

The experimental data on ΛN and ΣN scattering consist of some 600 events in the low energy (momenta of 200–300 MeV/c) region [14–16]. There exist another 250 events in the 300–1500 MeV/c momentum range [17]. The low energy data do not adequately

define even the dominant spin-singlet and spin-triplet scattering lengths and effective ranges. Nonetheless, it is clear that no bound two-body (YN) state exists. Therefore, the loosely bound hypertriton ($^3_{\Lambda}\text{H}$) plays a role in hypernuclear physics similar to that of the deuteron in the conventional arena.

4.2. The Hypertriton

Explicit $N\Lambda - N\Sigma$ mixing was demonstrated in Ref. [40] to play a crucial role in driving the Λ separation energy for $^3_{\Lambda}\text{H}$ from 2/3 MeV toward 0.1 MeV. Moreover, these separable potential model Faddeev calculations demonstrated that the small binding of possibly the world's largest halo nucleus — the Λ separation energy is about 2% of that of the neutron in ^3H and the Λ resides about 7 times as far from the ^2H center-of-mass as does that additional neutron in ^3H — was due to the existence of an attractive $NN\Lambda$ 3BF, when the $N\Sigma$ channel was formally eliminated. The dispersive energy dependence of the two-body force, which comes from embedding the $N\Lambda-N\Sigma$ potential in a three-body system is repulsive, as Peter Sauer has often noted. Gloeckle and coworkers [41] have since established that the $S = -1$ Nijmegen soft-core potential yields a value for $B_{\Lambda}(^3_{\Lambda}\text{H})$ which agrees with experiment. Miyagawa [42] has found similar results for the 1997 Nijmegen soft core model f [39]. Furthermore, he attributes 100 keV of the binding to the effective three-body force. Earlier calculations [43] using the Juelich potential models [36–38] were not in such good agreement with the experimental Λ separation energy.

4.3. The $A = 4$ Isodoublet

The $^4_{\Lambda}\text{He}-^4_{\Lambda}\text{H}$ isodoublet provides a strong test of our ability to model correctly the YN interaction. The quality of the calculations for both the ground states and the “spin-flip” excited states of this mirror pair should approach that demonstrated for the α particle. Even greater precision should be possible for the charge-symmetry-breaking (CSB) difference

$$\Delta B_{\Lambda} = B_{\Lambda}(^4_{\Lambda}\text{He}) - B_{\Lambda}(^4_{\Lambda}\text{H})$$

The nominal $\Delta B_{\Lambda} \simeq 350$ keV is much larger than the $\simeq 100$ keV CSB effect seen in the $^3\text{He}-^3\text{H}$ binding energy difference after correcting for the pp Coulomb energy in ^3He . A key question is whether the CSB can be understood in terms of the free $N\Lambda$ interaction. The question is not trivial [44]. For example, $N\Lambda-N\Sigma$ mixing introduces the 9 MeV mass difference of the charged Σ s. Mass mixing among the exchanged mesons in a potential model actually accounts for a significant aspect of CSB in the $^3\text{He}-^3\text{H}$ pair. Furthermore, a correct ordering of the $A=4$ isodoublet 0^+ and 1^+ states appears to require explicit $N\Lambda - N\Sigma$ mixing [45]. Simple single-channel four-body calculations would produce a ground state with 1^+ quantum numbers, because the low energy scattering parameters for the free $N\Lambda$ interaction indicate that the spin-triplet interaction is stronger than the spin-singlet interaction. Finally, Monte Carlo calculations [46] have indicated that suppression of $\Lambda \otimes ^4\text{He} \leftrightarrow \Sigma \otimes ^4\text{He}^*$ mixing, because of the large excitation energy of the $T = 1$ even parity $^4\text{He}^*$ states that result when the $T = 0\Lambda$ converts to a $T = 1\Sigma$, can account for the anomalously low value of $B_{\Lambda}(^5_{\Lambda}\text{He}) = 3.1$ MeV. More recently, Hiyama [47] has reported that, using a central potential approximation to the Nijmegen model D, neglecting $NNN\Lambda-NNN\Sigma$ mixing underbinds her full model calculation for the 0^+ states by about 1 1/2 MeV and underbinds her 1^+ state by at least an MeV.

5. Weak decay

The weak decay of the Λ , which decays freely into $p\pi^-$ (64%) and into $n\pi^0$ (36%), is first cousin to the familiar neutron β decay. The pion (mesonic) decay mode was first observed in the 1950s. However, the mesonic decay rates for hypernuclear weak decay are not well understood in terms of the underlying weak Hamiltonian; that is, we are not yet certain that the pionic decay rates are consistent with our parametrization of the free $\Lambda \rightarrow N + \pi$ and our microscopic models of the nuclear wave functions. This is in part because nonmesonic decay ($N + \Lambda \rightarrow N + N$) is the dominant decay mode for all but the lightest hypernucleus. The nonmesonic decay process provides our primary means of investigating the four-Fermion, strangeness-changing weak interaction. For example, the weak $\Lambda N \rho$ vertex can be investigated by no other means. Significant open issues in hypernuclear weak decay include:

- Does the $\Delta I = 1/2$ rule (an observed order of magnitude enhancement of the $\Delta I = 1/2$ amplitude over the $\Delta I = 3/2$ amplitude [48,49]), which governs mesonic K decay as well as the pionic decay modes of the Λ , also apply to the nonmesonic weak decay modes of hypernuclei?
- In the nonmesonic decay process, why should the rate for neutron stimulated emission and that for proton stimulated emission (Γ_{nn} and Γ_{np}) be essentially equal [20] when theoretical models [50,51] suggest that the $n + \Lambda \rightarrow n + n$ branch should be small?

The Λ decay into a π^+ in free space is forbidden. Yet the branching ratio for π^+ decay of $^4\Lambda$ He is approximately 5% [52]. Second order pion processes such as charge exchange ($\pi^- pp \rightarrow \pi^+ nn$) are too small [53] to explain more than about 1%. The free decay of a virtual Σ^+ was also found to be too small [54]. The virtual transition $p\Lambda \rightarrow n\Sigma^+$ followed by $\Sigma^+ N \rightarrow \pi^+ nN$ decay appears to play a significant role [55]. Experimentally the π^+ spectrum has a flat energy distribution and an isotropic angular distribution, both indicative of such a three-body decay mechanism. Thus, $^4\Lambda$ He $\rightarrow \pi^+ + X$ may provide observable evidence for the existence of a virtual Σ^+ in a Λ hypernucleus.

6. The $S = -2$ Puzzle

Because octet-octet mixing appears to play a key role in $S = -1$ physics, let us turn to the interesting puzzle that the single reported $^6\Lambda\Lambda$ He event [32] presents. [If the existence of $^6\Lambda\Lambda$ He is confirmed (there is a search underway at the BNL AGS), the mass of the $\Lambda\Lambda$ dibaryon would likely be limited to something like $2m_\Lambda - 10$ MeV; that is, the $\Lambda\Lambda$ would resemble a weakly bound di- Λ , a deuteron-like state.] Assuming that the $\Lambda\Lambda$ separation energy $B_{\Lambda\Lambda}(^6\Lambda\Lambda\text{He}) = B(^6\Lambda\Lambda\text{He}) - B(^4\text{He}) \simeq 10.9$ MeV is accurate (such an interpretation is consistent with the other accepted $\Lambda\Lambda$ events [13,56]), we see that the matrix element $\langle V_{\Lambda\Lambda} \rangle_{A=6}$ is relatively small:

$$-\langle V_{\Lambda\Lambda} \rangle = B_{\Lambda\Lambda}(^6\Lambda\Lambda\text{He}) - 2 \times B_\Lambda(^5\Lambda\text{He}) \simeq 10.9 - 2(3.1) = 4.7 \text{ MeV}.$$

This value is comparable with that of the $N\Lambda$ interaction: $\langle V_{\Lambda\Lambda} \rangle_{A=6} \simeq \langle V_{N\Lambda} \rangle_{A=4}$. Thus, both the $\Lambda\Lambda$ and $N\Lambda$ matrix elements are significantly smaller [57–59] than that of

the nn interaction: $\langle V_{nn} \rangle \simeq -7$ MeV. However, $\Lambda\Lambda$ and nn are analogs, belonging to the same 1S_0 multiplet. Why the disparity? Why should $\langle V_{\Lambda\Lambda} \rangle_{A=6}$ be much smaller?

Can we observe $\Lambda\Lambda$ scattering? "Yes, indirectly." Two examples of Ξ^- capture reactions that could measure $a_{\Lambda\Lambda}$ are $\Xi^-d \rightarrow \Lambda\Lambda n$ and $\Xi^-{}^7\text{Li} \rightarrow \Lambda\Lambda {}^6\text{He}$. The spectator particle would be detected in analogy with the a_{nn} measurements from $nd \rightarrow nn p$ and ${}^3\text{H} {}^3\text{H} \rightarrow nn {}^4\text{He}$. The $\Xi^-d \rightarrow \Lambda\Lambda n$ reaction has been investigated by Carr [60].

Lacking such data one can ask about the constraints that can be obtained from ${}_{\Lambda\Lambda} {}^6\text{He}$. In an $\alpha\Lambda\Lambda$ model analysis by Carr *et al.* [61], octet-octet mixing is essential. For an effective $\Lambda\Lambda$ potential whose strength is comparable to that of the NN force ($\bar{V}_{\Lambda\Lambda} \simeq V_{nn}$), overbinding of ${}_{\Lambda\Lambda} {}^6\text{He}$ results. In contrast, a coupled-channel ($\Lambda\Lambda - N\Xi$) potential

$$\begin{pmatrix} V_{\Lambda\Lambda} & V_{\Lambda\Lambda-N\Xi} \\ V_{\Lambda\Lambda-N\Xi} & V_{N\Xi} \end{pmatrix}$$

of similar overall strength yields binding comparable to experiment, because of Pauli blocking. The α core saturates the $(1s)^4$ shell, forcing a 5th nucleon into a higher shell and significantly weakening the effect of the $N\Xi$ component of the force. That is, by including $\Lambda\Lambda - N\Xi$ coupling explicitly, one can accommodate a weak $\langle V_{\Lambda\Lambda} \rangle_{A=6}$ even though the free space $\Lambda\Lambda - N\Xi$ potential is comparable in strength to the nn potential. In contrast, ${}_{\Lambda\Lambda} {}^5\text{H}$ should show evidence of enhanced binding. Moreover, observation of ${}_{\Lambda\Lambda} {}^4\text{H}$ may provide the optimum view of the free interaction. Whereas the ${}^4\text{He}$ core in ${}_{\Lambda\Lambda} {}^6\text{He}$ must be excited by 40 MeV to support the $\Lambda\Lambda \rightarrow N\Xi$ transition, the ${}^2\text{H}$ (${}^3\text{H}$) core in ${}_{\Lambda\Lambda} {}^4\text{He}$ (${}_{\Lambda\Lambda} {}^5\text{H}$) is bound by an additional 6 (20) MeV following $\Lambda\Lambda \rightarrow N\Xi$ ($\Lambda\Lambda \rightarrow p\Xi^-$) conversion, to form the ${}^3\text{H}$ or ${}^3\text{He}$ (${}^4\text{He}$) core that couples to the Ξ .

7. Conclusions

In summary, hypernuclear physics has been a source of surprising, new phenomena during its brief history. A number of novel and exciting features remain to be understood. New questions continue to arise. As a testing ground for $S = 0$ based concepts, hypernuclei are unsurpassed.

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