An EFT for the weak ΛN interaction *

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The nonleptonic weak $|\Delta S| = 1$ ΛN interaction, responsible for the dominant, nonmesonic decay of all but the lightest hypernuclei, is studied in the framework of an effective field theory. The long-range physics is described through tree-level exchange of the SU(3) Goldstone bosons (π and K), while the short-range potential is parametrized in terms of lowest-order contact terms obtained from the most general non-derivative local four-fermion interaction. Fitting to available weak hypernuclear decay rates for ${}^{5}_{\Lambda}$He, ${}^{11}_{\Lambda}$B and ${}^{12}_{\Lambda}$C yields reasonable values for the low-energy constants.

Λ-hypernuclei are bound systems composed of nucleons and one or more Λ hyperons. For the past 50 years, these bound systems have been used to extend our knowledge of both the strong and the weak baryon-baryon interaction from the NN case into the SU(3) sector. Up to date there are no stable hyperon beams, which makes hypernuclear weak decay the only source of information about the $|\Delta S| = 1$ four-fermion (4P) interaction. Since the decay of medium to heavy hypernuclei proceeds mainly via the $\Lambda N \to NN$ reaction, which does not conserve either parity, isospin or strangeness, its study complements the weak $\Delta S = 0$ NN case which allows the study of only the parity-violating (PV) amplitudes, because the parity-conserving (PC) signal is masked by the orders-of-magnitude stronger strong interaction background.

Since the earliest hypernuclear studies it is well known that the one-pion exchange mechanism (OPE), which naturally explains the long-range part of the interaction, is able to fairly reproduce the total NonMesonic Decay (NMD) rate of hypernuclei, but not the partial rates, produced by the proton-induced mechanism ($\Gamma_p : \Lambda p \to np$) and by the neutron-induced mechanism ($\Gamma_n : \Lambda n \to nn$). The ΛN mass difference, on the other hand, produces nucleons with momenta around ≈ 420 MeV, suggesting that the short-range part of the interaction cannot be neglected. These contributions have been described by: 1) the exchange of heavier mesons[1,2] whose production thresholds are too high for the free Λ decay; 2) an effective quark hamiltonian[3]; 3) correlated 2π-exchanges in the form of σ and ρ mesons[4]; 4) K-exchange plus correlated and uncorrelated 2π-exchanges[5].

The remarkable success of effective field-theory techniques based on chiral expansions

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in the SU(2) sector[6,7,8,9] suggests extending this approach to the SU(3) realm, even though stability of the chiral expansion is less clear here due to the significant degree of SU(3) symmetry breaking. A well-known example of the problems facing SU(3) chiral perturbation theory has been the prediction[10,11,12,13] of the four PC p-wave amplitudes in the weak nonleptonic decays of octet baryons, $Y \to N\pi$, with $Y = \Lambda, \Sigma$ or Ξ.

The aim of the present contribution is to build a less model dependent theory for the underlying $|\Delta S| = 1$ ΛN interaction governing hypernuclear decay, and see if a low order Effective Field Theory (EFT) can describe the present available hypernuclear decay data. Studies in this direction have already started [14,15]. In Ref. [14] a Fermi (V-A) interaction was added to the OPE mechanism to describe the weak $\Lambda N \to NN$ transition. Ref. [15] is a former version of the present manuscript, the latter containing a few more results.

In order to build such an EFT we shall allow for all possible contact terms in the four-baryon interaction Lagrangian, and fit the Low Energy Coefficients (LEC), which size those contributions, to the available data in the appropriate energy range. In contrast to the NN case, however, the $\Lambda N \to NN$ transition corresponds to an approximate energy release of 177 MeV at threshold. It is therefore not at all clear if low-energy expansions can be carried out with any validity. In light of this threshold momentum value, $|\vec{p}\,| \approx 417$ MeV, it is thus reasonable to include the pion ($m_\pi \approx 138$ MeV) and the kaon ($m_K \approx 494$ MeV) as dynamical fields. Working within SU(3) also supports treating pion and kaon on equal footing. The last member of the SU(3) Goldstone-boson octet, the η, is usually not included, since the strong ηNN coupling is an order of magnitude smaller than the strong πNN and $K\Lambda N$ ones [16]. Therefore, in the present study we will describe the long-range part of the weak $\Lambda N \to NN$ transition through pion and kaon exchanges, while the short-range interaction will be parametrized through leading-order contact terms. Not considered here is the intermediate-range 2π-exchange. Such a piece is two orders higher in the chiral expansion than the corresponding single pion-exchange piece. The pion exchange is given by the following strong and weak Lagrangians:

$$\begin{aligned}
\mathcal{L}^{\rm S}_{\rm NN\pi} &= -{\rm i}\, g_{\rm NN\pi}\, \overline{\psi}_{\rm N} \gamma_5 \gamma^\mu\, \vec{\tau} \cdot \partial_\mu\, \vec{\phi}^{\,\pi} \psi_{\rm N} \\
\mathcal{L}^{\rm W}_{\rm \Lambda N\pi} &= -{\rm i} G_F m_\pi^2 \overline{\psi}_{\rm N} (A_\pi + B_\pi \gamma_5 \gamma^\mu) \vec{\tau} \cdot \partial_\mu\, \vec{\phi}^{\,\pi} \psi_\Lambda \, \binom{0}{1} .
\end{aligned} \tag{1}$$

Here, $G_F m_\pi^2 = 2.21 \times 10^{-7}$ is the weak coupling constant, $g_{\rm NN\pi} = 13.16$, and the empirical constants $A_\pi = 1.05$ and $B_\pi = -7.15$, adjusted to the observables of the free Λ decay, determine the strength of the parity-violating and parity-conserving amplitudes, respectively. The iso-spurion $\binom{0}{1}$ is included to enforce the empirical $\Delta I = 1/2$ rule observed in the decay of a free Λ. Performing a nonrelativistic reduction of the resulting Feynman amplitude one obtains the OPE transition potential in momentum space:

$$V_{\rm OPE}(\vec{q}\,) = -G_F m_\pi^2 \frac{g_{\rm NN\pi}}{2M_N} \left(A_\pi + \frac{B_\pi}{2\overline{M}} \vec{\sigma}_1 \cdot \vec{q} \right) \frac{\vec{\sigma}_2 \cdot \vec{q}}{\vec{q}^{\,2} + {m_\pi}^2}\, \vec{\tau}_1 \cdot \vec{\tau}_2, \tag{2}$$

where $\vec{q}$ represents the momentum transfer directed towards the strong vertex [17], M_N is the nucleon mass and $\overline{M} = (M_N + M_\Lambda)/2$ is the average of the N and Λ masses.

The corresponding Lagrangians for the one-kaon-exchange (OKE) mechanism read:

$$\begin{aligned}
\mathcal{L}^{\rm S}_{\rm \Lambda NK} &= -{\rm i}\, g_{\rm \Lambda NK}\, \overline{\psi}_{\rm N} \gamma_5\, \gamma^\mu\, \partial_\mu\ \phi^{\rm K} \psi_\Lambda \\
\mathcal{L}^{\rm W}_{\rm NNK} &= -{\rm i}\, G_F m_\pi^2 \left[\overline{\psi}_{\rm N} \binom{0}{1}\ (C^{\rm PV}_{\rm K} + C^{\rm PC}_{\rm K} \gamma_5\, \gamma^\mu\, \partial_\mu\ ) (\phi^{\rm K})^\dagger \psi_{\rm N} \right.
\end{aligned}$$

$$+\overline{\psi}_{\rm N}\psi_{\rm N}\ (D_{\rm K}^{\rm PV}+D_{\rm K}^{\rm PC}\gamma_5\,\gamma^\mu\,\partial_\mu\ )(\phi^{\rm K})^\dagger\ \binom{0}{1}\Big]\,. \tag{3}$$

The weak C_K and D_K coefficients are related to the coupling constants at the p$\overline{\rm n}$K$^+$ and p$\overline{\rm p}$K^0 vertices, respectively. Like the $NN\pi$ coupling constant, the ΛNK and ΣNK ones are a fundamental input into our calculation. Unlike the $NN\pi$ coupling, however, their values are less well known, with $g_{\Lambda{\rm NK}} = 13-17$ and $g_{\Sigma{\rm NK}} = 3-6$. Here, we choose the values given by the Nijmegen Soft-Core model f, NSC97f [18], $g_{\Lambda{\rm NK}} = -17.66$ and $g_{\Sigma{\rm NK}} = -5.38$. The corresponding nonrelativistic OKE potential is analogous to Eq. (2), but with the replacements [1]:

$$\begin{aligned}
g_{{\rm NN}\pi} &\rightarrow g_{\Lambda{\rm NK}},\ \ m_\pi \rightarrow m_{\rm K}\\
A_\pi\,\vec\tau_1\cdot\vec\tau_2 &\rightarrow \left(\frac{C_{\rm K}^{\rm PV}}{2}+D_{\rm K}^{\rm PV}+\frac{C_{\rm K}^{\rm PV}}{2}\vec\tau_1\cdot\vec\tau_2\right)\frac{\overline{M}}{M_N}\\
B_\pi\,\vec\tau_1\cdot\vec\tau_2 &\rightarrow \left(\frac{C_{\rm K}^{\rm PC}}{2}+D_{\rm K}^{\rm PC}+\frac{C_{\rm K}^{\rm PC}}{2}\vec\tau_1\cdot\vec\tau_2\right)\,.
\end{aligned} \tag{4}$$

Of course, the weak NNK couplings are not accessible experimentally, so we obtain numerical values by making use of SU(3) and chiral algebra considerations [1,2]. In addition, our OPE and OKE potentials will be regularized by using monopole form factors at each vertex [17]. We note that all the strong model dependent ingredients used in the present calculation (as cut-off parameters or strong coupling constants) have been taken from the NSC97f interaction model [18].

Table 1
$\Lambda N \rightarrow NN$ partial waves.

partial wave	operator	order	I
$a :^1 S_0 \rightarrow^1 S_0$	$\hat{1}$, $\vec\sigma_1\vec\sigma_2$	1	1
$b :^1 S_0 \rightarrow^3 P_0$	$(\vec\sigma_1-\vec\sigma_2)\vec q$, $(\vec\sigma_1\times\vec\sigma_2)\vec q$	q/M_N	1
$c :^3 S_1 \rightarrow^3 S_1$	$\hat{1}, \vec\sigma_1\vec\sigma_2$	1	0
$d :^3 S_1 \rightarrow^1 P_1$	$(\vec\sigma_1-\vec\sigma_2)\vec q$, $(\vec\sigma_1\times\vec\sigma_2)\vec q$	q/M_N	0
$e :^3 S_1 \rightarrow^3 P_1$	$(\vec\sigma_1+\vec\sigma_2)\vec q$	q/M_N	1
$f :^3 S_1 \rightarrow^3 D_1$	$(\vec\sigma_1\times\vec q)(\vec\sigma_2\times\vec q)$	$q^2/{M_N}^2$	0

If no model is assumed, the low energy $\Lambda N \rightarrow N_1N_2$ process can be parametrized through the 6 partial waves listed in Table I [1]. The PC a and c transitions can only be produced by the $\hat{1}\cdot\delta^3(\vec r\ )$ and $\vec{\sigma_1}\cdot\vec\sigma_2\cdot\delta^3(\vec r\ )$ operators, where $\delta^3(\vec r\ )$ represents the contact interaction. The PV b and d transitions proceed through the combination of the spin-nonconserving operators: $(\vec\sigma_1-\vec\sigma_2)$ and i$(\vec\sigma_1\times\vec\sigma_2)$ with the following operators: $\{\vec{p_1}-\vec{p_2}\,,\,\delta^3(\vec r\ )\,\}$ and $[\,\vec{p_1}-\vec{p_2}\,,\,\delta^3(\vec r\ )\,]$, where $\vec{p_i}$ is the derivative operator acting on the

[1]This illustrative argument is only valid for the $L=0$ part of the ΛN wave function.

"ith" particle [2]. The e transition is allowed by the combination of the spin-conserving operator $(\vec{\sigma}_1 + \vec{\sigma}_2)$ with the previous commutator and anti-commutator, while only two-derivative operators can produce the last (tensor) transition. In order to write the general four-fermion interaction we will retain those terms depending on the momentum transfer $\vec{q} = \vec{p}_\Lambda - \vec{p}_{N_1} = \vec{p}_{N_2} - \vec{p}_N = \vec{k}_i - \vec{k}_f$, with $\vec{k}_i$ $(\vec{k}_f)$ being the initial (final) relative momentum in the c.m. frame [3]. The 4P potential reads (in units of $G_F = 1.166 \times 10^{-11}$ MeV^{-2}):

$$\begin{aligned} V_{4P}(\vec{q}) &= C_0^0 + C_0^1\, \vec{\sigma}_1 \cdot \vec{\sigma}_2 \; + C_1^0\, \frac{\vec{\sigma}_1 \cdot \vec{q}}{2\overline{M}} + C_1^1\, \frac{\vec{\sigma}_2 \cdot \vec{q}}{2M} + \mathrm{i}\, C_1^2\, \frac{(\vec{\sigma}_1 \times \vec{\sigma}_2)\, \cdot \vec{q}}{2\tilde{M}} \\ &+ C_2^0\, \frac{\vec{\sigma}_1 \cdot \vec{q}\, \vec{\sigma}_2 \cdot \vec{q}}{4M\overline{M}} + C_2^1\, \frac{\vec{\sigma}_1 \cdot \vec{\sigma}_2\, \vec{q}^{\,2}}{4M\overline{M}} + C_2^2\, \frac{\vec{q}^{\,2}}{4M\tilde{M}} \end{aligned} \tag{5}$$

where $\tilde{M} = \dfrac{3M + M_\Lambda}{4}$ is a weighted average of N, Λ masses and C_i^j is the jth LEC at ith order. While the form of the contact terms is model independent, the size of the coefficients depends on how the theory is formulated, and they are expected to be of the order of the other couplings in the problem. These couplings provide a very simple representation of the short distance contributions to the process at hand. In a complete model, they would be represented by specific dynamical contributions, such as $\rho, \omega, etc.$-exchange. However, we eschew the temptation to be more specific—in fact this generality is one of the strengths of our approach. We evaluate the coefficients purely phenomenologically and leave the theoretical interpretation of the pieces to future investigations. Of course, the specific size of such coefficients depends upon the chiral order to which we are working. However, if the expansion is convergent, then the values of these effective couplings should be relatively stable as NLO or higher effects are included.

To reduce the number of free parameters we use power counting, discarding operators of order $q^2/{M_N}^2$. To obtain the 4P potential in configuration space we must Fourier transform $V_{4\mathrm{P}}(\vec{q})$, smearing the resulting delta functions by using a normalized Gaussian form for the 4-fermion contact potential, $f_{ct}(r) = e^{-\frac{r^2}{\delta^2}}/(\delta^3\pi^{3/2})$, where δ is taken to be of the order of a typical vector-meson range, $\delta \sim \sqrt{2}m_\rho^{-1} \approx 0.36$ fm. The leading order $V_{4\mathrm{P}}(\vec{r})$ potential for both PV and PC terms can then be written as:

$$\begin{aligned} V_{4P}(\vec{r}) &= \left\{ C_0^0 + C_0^1\, \vec{\sigma}_1 \cdot \vec{\sigma}_2 + \frac{2r}{\delta^2} \left[C_1^0\, \frac{\vec{\sigma}_1 \cdot \hat{r}}{2\overline{M}} + C_1^1\, \frac{\vec{\sigma}_2 \cdot \hat{r}}{2M} + C_1^2\, \frac{(\vec{\sigma}_1 \times \vec{\sigma}_2) \cdot \hat{r}}{2\tilde{M}} \right] \right\} \\ &\times \; f_{ct}(r) \; \times \left[C_{IS}\, \hat{1} + C_{IV}\, \vec{\tau}_1 \cdot \vec{\tau}_2 \right] , \end{aligned} \tag{6}$$

where the last factor represents the isospin part of the 4-fermion interaction. Note that we only allow for $\Delta I = 1/2$ transitions.

It is well known that the high momentum transferred in the $\Lambda N \to NN$ reaction makes this process sensitive to the short range physics which is characterized by our contact coefficients. Moreover, since the $|\Delta S| = 1$ reaction takes place in a finite nucleus, extracting

[2] Note that we are assuming that $\vec{p}_1 - \vec{p}_2$ is small enough to disregard higher powers of the derivative operators $\vec{p}_1 - \vec{p}_2$.

[3] This is a reasonable assumption for the weak $\Lambda N \to N_1 N_2$ process, where the two particles in the initial state (bound in a hypernucleus) have very low momentum, specially when compared to the momentum of the two outgoing nucleons. That means that we can neglect terms containing $\vec{k}_i$ in the expansion, and at the same time approximate $\vec{k}_f$ by $-\vec{q}$.

information of the elementary weak two-body interaction requires a careful investigation of the many-body nuclear effects present in the hypernucleus. The nonmesonic decay rate is written as:

$$\Gamma_{\rm nm} = \int \frac{d^3k_1}{(2\pi)^3} \int \frac{d^3k_2}{(2\pi)^3} \sum_{\substack{M_J\{R\}\\ \{1\}\{2\}}} (2\pi)\, \delta(M_H - E_R - E_1 - E_2) \frac{1}{(2J+1)} \mid \mathcal{M}_{fi} \mid^2 ,$$

where k_1 and k_2 represent the momenta of the two outgoing nucleons, J and M_J the total spin and spin projection of the initial hypernucleus, R the quantum numbers of the residual nucleus, M_H the mass of the hypernucleus, and $\mathcal{M}_{fi}$ the hypernuclear transition amplitude, given by $\mathcal{M}_{fi} \sim \langle \vec{k}_1 m_1 \vec{k}_2 m_2; \Psi_{\rm R}^{A-2} \mid \hat{O}_{\Lambda {\rm N} \to {\rm NN}} \mid^A_\Lambda Z\rangle$.

Table 2
Experimental data for the decay of ${}^5_\Lambda$He, ${}^{11}_\Lambda$B and ${}^{12}_\Lambda$C. The values underlined are the ones used in our fits.

	Γ	$\Gamma_{\rm n}/\Gamma_{\rm p}$	$\Gamma_{\rm p}$	$\mathcal{A}$
${}^5_\Lambda$He	$\underline{0.41 \pm 0.14}$ [19]	0.93 ± 0.55 [19]	$\underline{0.21 \pm 0.07}$ [19]	$\underline{0.24 \pm 0.22}$ [20]
	$\underline{0.50 \pm 0.07}$ [21]	1.97 ± 0.67 [21]		
		$\underline{0.50 \pm 0.10}$ [22]		
${}^{11}_\Lambda$B	$\underline{0.95 \pm 0.14}$ [21]	$\underline{1.04^{+0.59}_{-0.48}}$ [19]	$0.30^{+0.15}_{-0.11}$ [21]	-0.20 ± 0.10 [23]
		$2.16 \pm 0.58^{+0.45}_{-0.95}$ [21]		
		$0.59^{+0.17}_{-0.14}$ [24]		
${}^{12}_\Lambda$C	$\underline{0.83 \pm 0.11}$ [25]	$1.33^{+1.12}_{-0.81}$ [19]	$0.31^{+0.18}_{-0.11}$ [21]	-0.01 ± 0.10 [23]
	$\underline{0.89 \pm 0.15}$ [21]	$1.87 \pm 0.59^{+0.32}_{-1.00}$ [21]		
	$\underline{1.14 \pm 0.2}$ [19]	$0.59^{+0.17}_{-0.14}$ [24]		
		$\underline{0.87 \pm 0.23}$ [26]		

In the present calculation, we use a shell-model for the initial hypernucleus and assume a weak coupling scheme for the hyperon (the Λ particle will only couple the ground state core wave function). Moreover, using coefficients of fractional parentage we can decouple the weakly interacting nucleon from the antisymetrized core wave function, leaving also an antisymmetric wave function for the residual system. The single-particle Λ and N orbits are taken to be solutions of harmonic oscillator mean field potentials with parameters adjusted to experimental separation energies and charge form factor of the hypernucleus under study. The strong hyperon-nucleon interaction at short distances, absent in mean-field models, is accounted for by replacing the mean-field two-particle ΛN wave function by a correlated [27] one inspired in a microscopic finite-nucleus G-matrix calculation [28] which uses the soft-core and hard-core Nijmegen models [29]. The NN wave function is obtained by solving the Lippmann-Schwinger (T-matrix) equation with the input of the

Nijmegen Soft Core NSC97f potential model (details of the calculation can be found in the Appendix of Ref. [17]).

Table 3
Results for the weak decay observables, when a fit to the Γ and n/p for ${}^{5}_{\Lambda}$He, ${}^{11}_{\Lambda}$B and ${}^{12}_{\Lambda}$C is performed. The values in parentheses have been obtained including $\alpha_\Lambda({}^{5}_{\Lambda}\mathrm{He})$ in the fit.

	π	$+K$	+ LO PC	+ LO PV	EXP:
$\Gamma({}^{5}_{\Lambda}\mathrm{He})$	0.46	0.29	0.44	0.44 (0.44)	0.41 ± 0.14 [19], 0.50 ± 0.07 [21]
$n/p({}^{5}_{\Lambda}\mathrm{He})$	0.09	0.51	0.56	0.54 (0.54)	0.93 ± 0.55 [19], 0.50 ± 0.10 [22]
$\alpha_\Lambda({}^{5}_{\Lambda}\mathrm{He})$	-0.23	-0.54	-0.72	-0.33 (0.24)	0.24 ± 0.22 [20]
$\Gamma({}^{11}_{\Lambda}\mathrm{B})$	0.67	0.43	0.87	0.87 (0.87)	0.95 ± 0.14 [21]
$n/p({}^{11}_{\Lambda}\mathrm{B})$	0.11	0.45	0.82	0.99 (0.99)	$1.04^{+0.59}_{-0.48}$ [19]
$\mathcal{A}({}^{11}_{\Lambda}\mathrm{B})$	-0.11	-0.24	-0.20	-0.02 (0.12)	-0.20 ± 0.10 [23]
$\Gamma({}^{12}_{\Lambda}\mathrm{C})$	0.80	0.49	0.95	0.93 (0.93)	1.14 ± 0.2 [19], 0.89 ± 0.15 [21]
					0.83 ± 0.11 [25]
$n/p({}^{12}_{\Lambda}\mathrm{C})$	0.09	0.36	0.64	0.82 (0.81)	0.87 ± 0.23 [26]
$\mathcal{A}({}^{12}_{\Lambda}\mathrm{C})$	-0.03	-0.06	-0.05	-0.006 (0.03)	-0.01 ± 0.10 [23]
$\hat{\chi}^2$			0.98	1.49 (1.12)	

Before we start with the discussion of the results, we should make a remark on the data. One might wonder if there can be only three independent data points in the nonmesonic decay: the proton-induced and neutron-induced rates Γ_p and Γ_n, and the asymmetry $\mathcal{A}$ (associated with the proton-induced decay), relating observables from one hypernucleus to another through hypernuclear structure coefficients. While one may indeed expect measurements from different p-shell hypernuclei, say, A=12 and 16, to provide the same constraint, the situation is different when including data from s-shell hypernuclei like A=5. For the latter, the initial ΛN pair can only be in a relative s-state, while for the former, relative p-states are allowed as well. We therefore include in our fits the data shown in Table 2 from the A=5,11 and 12 hypernuclei. Only more recent measurements from the last 12 years were used, however, we excluded those recent data of the ratio $\Gamma_n/\Gamma_p = n/p$ whose error bars were larger than 100%. We also have taken the data at face value and have not applied any corrections due to, e.g., the two-nucleon induced mechanism [30]. Given the sizable error bars of the data, this omission does not change any of our conclusions. New exclusive measurements will hopefully be able to verify the magnitude of such contributions.

No parameters were fitted for the results with only π and K exchange, shown in Table 3. As has been known for a long time, π exchange alone describes reasonably well the observed total rates, while dramatically underestimating the n/p ratio. Incorporation of kaon exchange gives a destructive interference between both mechanisms (OPE and OKE) in the PC amplitudes, while the interference is constructive in the PV ones. The tensor PC channel dominates the proton-induced rate while it is absent in the $L = 0$ neutron-induced one. As a consequence, n/p is enhanced by about a factor of five, within reach of the lower bounds of the experimental measurements, and the total rate underpredicts the observed value by about a factor of two. It also leads to values for the asymmetry that are close to experiment for the p-shell hypernuclei, but far off for A=5 [4]. Since the contributions of both η-exchange and two-pion exchange are negligible, these discrepancies illustrate the need for short-range physics.

Table 4
LEC coefficients corresponding to the LO calculation. The values in parentheses have been obtained including $\alpha_\Lambda({}^5_\Lambda\mathrm{He})$ in the fit.

	+ LO PC	+LO PV
C_0^0	-1.46 ± 0.44	-1.12 ± 0.52 (-0.95 ± 0.31)
C_0^1	-0.85 ± 0.27	-0.99 ± 0.74 (-0.53 ± 0.24)
C_1^0	$---$	-5.71 ± 4.16 (-4.99 ± 1.62)
C_1^1	$---$	2.95 ± 2.90 (2.71 ± 1.91)
C_1^2	$---$	-6.56 ± 1.78 (-5.24 ± 1.93)
C_{IS}	4.85 ± 1.45	4.69 ± 0.11 (5.97 ± 0.71)
C_{IV}	1.43 ± 0.45	1.25 ± 0.11 (1.59 ± 0.21)

Allowing contact terms of order unity (leading-order PC operators) to contribute leads to four free parameters, C_0^0, C_0^1, C_{IS} and C_{IV}. Data on the total and partial decay rates for all three hypernuclei are included in the fit, but no asymmetry measurements. The inclusion of the contact terms roughly doubles the values for the total decay rates, thus restoring agreement with experiment. The impact on the n/p ratio is noteworthy: the value for ${}^5_\Lambda$He increases by 10% while the n/p ratios for ${}^{11}_\Lambda$B and ${}^{12}_\Lambda$C almost double. This is an example of the differing impact certain operators can have for s- and p-shell hypernuclei. The effect on the asymmetry is opposite, almost no change for A=11 and 12, but a 30% change for A=5. The magnitudes of the four parameters, C_0^0, C_0^1, C_{IS} and C_{IV}, listed in Table 4, are all around their natural size of unity, with the exception of C_{IS} which is a factor of five or so larger. Note the substantial error bars on all the parameters, reflecting the uncertainties in the measurements.

Three new parameters are admitted when we allow the leading-order PV terms (of order q/M_N) to contribute with the coefficients C_1^0, C_1^1, and C_1^2. As shown in Table 4, the parameters for the PV contact terms are larger than the ones for the PC terms, and in fact, compatible with zero. Including the three new parameters does not substantially alter

[4]Note that α_Λ stands for the intrinsic Λ asymmetry parameter, characteristic of the elementary $\Lambda N \to NN$ reaction. For helium, this quantity can be extracted directly from experiment without the model dependent input of the Λ (or hypernuclear) polarization [20].

the previously fitted ones, thus supporting the validity of our expansion. Regarding their impact on the observables, the PV contact terms barely modify the total and partial rates but significantly affect the asymmetry, as one would expect for an observable defined by the interference between PV and PC amplitudes. The calculated asymmetry considerably decreases in size for all three hypernuclei, giving p-shell values close to zero, but still negative. In order to further understand this behavior, we have performed a number of fits including the asymmetry data points of either ${}^{5}_{\Lambda}$He or ${}^{11}_{\Lambda}$B or both. Tables 3 and 4 display the result of one of those fits. Inclusion of the ${}^{5}_{\Lambda}$He (intrinsic $\Lambda-$) asymmetry helps in constraining the values of two of the LO PV parameters. We find that the two present experimental values for A=5 and A=11 cannot be fitted simultaneously with this set of contact terms. Future experiments will have to settle this issue.

Table 5
Sensitivity of the calculation to the strong interaction model. The numbers in the left (right) column use final NN wave functions obtained from the NSC97f (NSC97a) interaction[18]. The $\alpha_\Lambda({}^{5}_{\Lambda}\mathrm{He})$ has been included in the fit.

$\pi + K +$ LO PC $+$ LO PV			$\pi + K +$ LO PC $+$ LO PV		
	NSC97a	NSC97f		NSC97a	NSC97f
$\Gamma({}^{5}_{\Lambda}\mathrm{He})$	0.44	0.44	C^0_0	-0.67 ± 0.42	-0.95 ± 0.31
$n/p({}^{5}_{\Lambda}\mathrm{He})$	0.55	0.54	C^1_0	-0.34 ± 0.32	-0.53 ± 0.24
$\alpha_\Lambda({}^{5}_{\Lambda}\mathrm{He})$	0.24	0.24			
$\Gamma({}^{11}_{\Lambda}\mathrm{B})$	0.86	0.87	C^0_1	-5.85 ± 1.40	-4.99 ± 1.62
$n/p({}^{11}_{\Lambda}\mathrm{B})$	0.95	0.99	C^1_1	3.65 ± 1.66	2.71 ± 1.91
$\mathcal{A}({}^{11}_{\Lambda}\mathrm{B})$	0.08	0.12	C^2_1	-6.47 ± 1.64	-5.24 ± 1.93
$\Gamma({}^{12}_{\Lambda}\mathrm{C})$	0.94	0.93	C_{IS}	5.78 ± 0.86	5.97 ± 0.71
$n/p({}^{12}_{\Lambda}\mathrm{C})$	0.77	0.81	C_{IV}	1.65 ± 0.27	1.59 ± 0.21
$\mathcal{A}({}^{12}_{\Lambda}\mathrm{C})$	0.02	0.03			
$\hat{\chi}^2$	1.19	1.12			

We have also performed fits allowing a contribution from an isospin $\Delta I = 3/2$ transition operator in Eq. 6. The resulting fit shifts strength from the isoscalar contribution to the new $\Delta I = 3/2$ one, leaving the other parameters unchanged. In any case, as shown in Table 3, we can clearly get an excellent fit to all observables without such transitions while obtaining constants of reasonable size. Table 5 demonstrates that our conclusions are basically independent of the model we use for the strong force used for describing Final State Interactions in the transition. Employing NN wave functions that are obtained with either the NSC97f or the NSC97a model, one can perfectly fit the total and partial rates, as well as the helium asymmetry, while the predicted $p-$shell asymmetries show a variation range of 50%. The constants can easily absorb the changes (the largest change affects the LO PC C_0^1, which becomes compatible with zero when the NSC97a model is used) and remain compatible within their error bars. Similarly, the results in Table 6 show the insensitivity of the predicted observables to realistic values of the δ range used to smear the delta function in the contact potential.

Table 6
Sensitivity of the calculation to the 4-fermion contact range δ in the smeared delta function. The $\alpha_\Lambda({}^5_\Lambda\text{He})$ is included in the fit.

	$\delta \approx 0.3$fm	$\delta \approx 0.36$fm	$\delta \approx 0.4$fm
$\mathcal{A}({}^{11}_\Lambda\text{B})$	0.12	0.12	0.12
$\mathcal{A}({}^{12}_\Lambda\text{C})$	0.03	0.03	0.03
C_0^0	-1.81 ± 0.43	-0.95 ± 0.31	-0.69 ± 0.27
C_0^1	-1.03 ± 0.37	-0.53 ± 0.24	-0.37 ± 0.20
C_1^0	-7.10 ± 2.38	-4.99 ± 1.62	-4.13 ± 1.35
C_1^1	4.00 ± 2.93	2.71 ± 1.91	2.15 ± 1.56
C_1^2	-7.44 ± 3.01	-5.24 ± 1.93	-4.33 ± 1.56
C_{IS}	6.43 ± 0.60	5.97 ± 0.71	5.69 ± 0.79
C_{IV}	1.82 ± 0.19	1.59 ± 0.21	1.44 ± 0.23
$\hat{\chi}^2$	1.12	1.12	1.13

In conclusion, we have studied the nonmesonic weak decay using an Effective Field Theory framework for the weak interaction. The long-range components were described with pion and kaon exchange, while the short-range part is parametrized in leading-order PV and PC contact terms. We find coefficients of natural size with significant error bars, reflecting the level of experimental uncertainty. We found a large contribution from an isoscalar, spin-independent central operator. There is no indication of any contact terms violating the $\Delta I = 1/2$ rule. In this study we have not speculated on the dynamical origin of these contact contributions but our aim was to ascertain their size and verify the validity of the EFT framework for the weak decay. The next generation of data from recent high-precision weak decay experiments currently under analysis holds the promise to provide much improved constraints for studies of this nature.

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