Two-pion-exchange in the non-mesonic weak decay of \( \Lambda \)-hypernuclei

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Abstract

The non-mesonic weak decay of \( \Lambda \)-hypernuclei is studied within a one-meson-exchange potential supplemented by a chirally motivated two-pion-exchange mechanism. The effects of final state interactions on the outgoing nucleons are also taken into account. In view of the severe discrepancies between theoretical expectations and experimental data, particular attention is paid to the asymmetry of the protons emitted by polarized hypernuclei. The one-meson-exchange model describes the non-mesonic rates and the neutron-to-proton ratio satisfactorily but predicts a too large and negative asymmetry parameter. The uncorrelated and correlated two-pion mechanisms change the rates moderately, thus maintaining the agreement with experiment. The modification in the strength and sign of some decay amplitudes becomes crucial and produces asymmetry parameters which lie well within the experimental observations.

Key words: Non-Mesonic Weak Decay, Decay Asymmetries of Polarized Hypernuclei, \( \Gamma_n/\Gamma_p \) ratio
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1 Introduction

Hypernuclei may be considered as a powerful “laboratory” for unique investigations of the baryon-baryon strangeness-changing weak interactions. The field of non-mesonic weak decay has indeed experienced a phase of renewed interest in the last few years. On the one hand, different theoretical \([1–7]\) and experimental \([8–11]\) indications have recently appeared in favour of a solution of the well-known and long-standing puzzle on the ratio \( \Gamma_n/\Gamma_p \) \([12, 13]\)
between the rates for the $\Lambda n \rightarrow nn$ and $\Lambda p \rightarrow np$ non-mesonic weak decay processes. Nowadays, values of this neutron-to-proton ratio around 0.3-0.4 for s- and p-shell hypernuclei are common to both theoretical and experimental analyses. An important role in this achievement has been played by a non-trivial interpretation of data, which required analyses of two-nucleon induced decays, $\Lambda NN \rightarrow nNN$, and accurate studies of nuclear medium effects on the weak decay nucleons. On the other hand, considerable concern is rooted in the persistence of another open problem, of more recent origin, which regards an asymmetry in the non-mesonic weak decay of polarized hypernuclei and whose solution is expected to provide new constraints for a deeper understanding of the dynamics of hypernuclear decay. New phenomenological information is in principle accessible since this asymmetry in the angular emission of protons originates from the interference among parity-conserving (PC) and parity-violating (PV) $\bar{\Lambda}p \rightarrow np$ transitions amplitudes [14, 15], while the widely considered decay widths $\Gamma_n$ and $\Gamma_p$ are the result of the incoherent sum of PC and PV amplitudes squared.

While inexplicable inconsistencies appeared between the first asymmetry experiments of Refs. [16, 17], as discussed in Ref. [13], very recent and more accurate data [9, 18, 19] favour small observable asymmetries, compatible with a vanishing value, for both s- and p-shell hypernuclei. On the contrary, theoretical models based on one-meson-exchange potentials [1, 2, 4, 20, 21] (generally including the pseudoscalar and vector mesons $\pi$, $\rho$, $K$, $K^*$, $\omega$ and $\eta$) and/or direct quark mechanisms [2] predicted rather large and negative intrinsic asymmetry values (from $-0.7$ to $-0.5$ in the above quoted studies). It must be noted that, on the contrary, the mentioned models have been able to account fairly well for the total non-mesonic decay rates and for the ratios $\Gamma_n/\Gamma_p$ measured for s- and p-shell hypernuclei. Note also that, when discussing the comparison between theory and experiment for the asymmetric non-mesonic decay, one should distinguish (as we have done above) between intrinsic and observable asymmetries [20]. Due to nucleon final state interactions acting after the decay, the former ($a_\Lambda$) is of theoretical relevance only, while the real observable is the latter ($a_\Lambda^M$). A strong dependence of $a_\Lambda^M$ on final state interactions and proton detection threshold $T_p^{th}$ has been obtained in Ref. [20], with values of $a_\Lambda^M/a_\Lambda$ for standard experimental conditions ($T_p^{th} \simeq 40$ MeV) of about 0.7 for $^5\Lambda$He and 0.6 for $^{11}\Lambda$B and $^{12}\Lambda$C.

Recently, an effective field theory approach based on tree-level pion- and kaon-exchange and leading-order contact interactions has been applied to hypernuclear decay [22]. The coefficients of the considered four-fermion point interaction Lagrangian have been fitted to reproduce the total and partial decay widths for $^5\Lambda$He, $^{11}\Lambda$B and $^{12}\Lambda$C, and the asymmetry parameter for $^5\Lambda$He. In this way, a dominating central, spin-and isospin-independent contact term has been predicted. Such term turned out to be particularly important in order to reproduce a small and positive value of the intrinsic asymmetry for $^5\Lambda$He, as
indicated by the recent KEK experiments. In order to improve the comparison with the observed decay asymmetries in a calculation scheme based on a meson-exchange model, this result could be interpreted dynamically as the need for the introduction of a scalar-isoscalar meson-exchange.

Prompted by the work of Ref. [22], in Ref. [23] a model based on \((\pi + K)\)-exchange and the direct quark mechanism has been supplemented with the exchange of the scalar-isoscalar \(\sigma\)-meson. Rather curiously, to date only a few works considered such a meson as mediator of the non-mesonic decay, despite it could in principle be relevant, being the chiral partner of the pion in QCD and given that its mass is comparable with the exchanged momenta in the \(\Lambda N \rightarrow nN\) process and with the mass of the kaon, a meson whose importance in accounting for the decay rates is, on the contrary, well established [2–4].

The strategy of Ref. [23] has been to determine the weak couplings of the \(\sigma\) by fitting decay data for \(s\)-shell hypernuclei. The \(\pi + K + \sigma + \) direct quark model turned out to reproduce the data for \(\Gamma_{NM} = \Gamma_n + \Gamma_p\), \(\Gamma_n/\Gamma_p\) and \(a_{\Lambda}^M\) for \(^5\Lambda\text{He}\) quite reasonably, while the \(\pi + K + \sigma\) model was unable to account for the experimental value of \(a_{\Lambda}^M(\Lambda\text{He})\) [24]. Although the same model should also be tested for \(p\)-shell hypernuclei, the results of Ref. [23] clearly demonstrate the importance of the \(\sigma\)-exchange in the non-mesonic decay.

A one-meson-exchange potential containing \(\pi, \rho, K, K^*, \omega, \eta\) and \(\sigma\) has been applied more recently [25] to the evaluation of \(\Gamma_{NM}, \Gamma_n/\Gamma_p\) and \(a_{\Lambda}\) for \(^5\Lambda\text{He}\) and \(^{12}\Lambda\text{C}\). The unknown \(\sigma\) couplings have been fixed to reproduce a subset of decay observables \([\Gamma_{NM} \text{ and } \Gamma_n/\Gamma_p \text{ for } ^5\Lambda\text{He}], \text{while the remaining observables have been predicted and compared with data. The authors found that, despite the inclusion of the } \sigma \text{ meson improved the overall agreement with experiment, the asymmetry data for } ^5\Lambda\text{He could not be reproduced} [24].

The contributions of uncorrelated and correlated two-pion-exchange to the non-mesonic weak decay has also been studied in Refs. [3, 5, 26, 27]. Some preliminary papers [26] paved the way for a model that considered, in addition to \(\pi\) and \(\omega\), the exchange of two-pions correlated in the scalar-isoscalar \((2\pi/\sigma)\) and vector-isovector \((2\pi/\rho)\) channels [5], with a phenomenological treatment which is quite similar to the scheme used in the pioneering work of Ref. [28]. The results of Ref. [5] demonstrate how the correlated two-pion-exchange improves the calculation of \(\Gamma_n/\Gamma_p\) over the one-pion-exchange model. After adding the exchange of the kaon, within a \(\pi + K + \omega + 2\pi/\sigma + 2\pi/\rho\) model it was found [27] that the correlated two-pion-exchange also entails some improvement in the evaluation of the asymmetry parameter. In Ref. [3], a meson-exchange potential including pion, kaon, omega and uncorrelated plus correlated two-pions \((\pi + K + \omega + 2\pi + 2\pi/\sigma)\) has been considered. The correlated two-pion-exchange in the scalar-isoscalar channel has been treated in terms of a chiral unitary approach which has revealed to reproduce well \(\pi\pi\) scattering data in the scalar sector and in which the \(\sigma\)-meson appears as a dynamically
generated resonance. The only free parameter of the model is the momentum
cutoff that regularizes the loop integrals, which is fixed to about 1 GeV in order
to produce a σ resonance at the observed mass of 450 MeV and having a large
width. A sizable cancellation between $2\pi$- and $2\pi/\sigma$-exchange was found for
the relevant momenta ($\simeq 410$ MeV) in the non-mesonic decay. Consequently,
the total two-pion-exchange contribution to the decay rates turned out to be
moderate but its effect on the asymmetry parameter was not evaluated.

Motivated by the findings of Refs. [3, 22, 23, 25, 27], in the present paper we in-
vestigate the effects on the non-mesonic decay observables for s- and p-shell hy-
permuceli of uncorrelated two-pion-exchange and correlated two-pion-exchange
in the scalar-isoscalar channel. We employ a finite nucleus approach and pay
special attention to the proton asymmetry. The weak transition potentials for
two-pion-exchange are adopted from Refs. [3] and are added to the exchange
of the pseudoscalar and vector mesons $\pi$, $\rho$, $K$, $K^*$, $\omega$ and $\eta$, with potentials taken from Ref. [1]. Correlated two-pion-exchange in the vector-isovector
channel is not considered here, since the one-meson-exchange potential we use
already includes the $\rho$-meson. We do not take into account the two-nucleon
induced decay mode, $\Lambda NN \rightarrow nNN$ [12, 13, 29]. This channel can be safely
neglected when evaluating the decay asymmetries [20] even if its contribution
to the total non-mesonic decay rate is significant. When comparing with data
for this rate, the following predictions obtained in Refs. [6, 7, 30] should be
taken into account: $\Gamma_2/(\Gamma_n + \Gamma_p) \simeq 0.20$ for $^5\Lambda$He and $\Gamma_2/(\Gamma_n + \Gamma_p) \simeq 0.25$ for
$^{11}\Lambda$B and $^{12}\Lambda$C.

With the novel meson-exchange model introduced in the present paper it turns
out to be possible to reproduce quite reasonably experimental data for both
the decay rates ($\Gamma_{NM}$ and $\Gamma_n/\Gamma_p$) and the observable asymmetry parameter
in s- and p-shell hypernuclei. The introduction of two-pion-exchange reveals
to be of great importance in this achievement.

The paper is organized as follows. The formalism employed for the calculation
of decay rates and asymmetries is outlined in Section 2. Numerical results for
these observables are presented and compared with data in Section 3. Finally,
in Section 4 we draw our conclusions.

2 Formalism

In this Section we briefly present the formalism adopted for the calculation of
the non-mesonic weak decay rates and proton asymmetries.
2.1 Non-mesonic decay rates

The rate associated to a neutron (proton) stimulated decay can be evaluated by the following average:

$$\Gamma_{n(p)} = \frac{1}{2J + 1} \sum_{M_J} \sigma_{n(p)}(J, M_J),$$

(1)

in terms of the intensities $$\sigma_{n(p)}(J, M_J)$$ of neutrons (protons) emitted along the quantization axis in the non-mesonic decay of a hypernucleus with third component $$M_J$$ of the total spin $$J$$.

Following the approach used in Ref. [1], we apply standard nuclear structure techniques that allow us to write these intensities in terms of two-body amplitudes involving a $$\Lambda N$$ pair in the initial state and a $$NN$$ pair in the final state. To do this, one needs to decouple the hyperon (with spin and isospin quantum numbers $$j_\Lambda, m_\Lambda, t_\Lambda = 0, t_{3\Lambda} = 0$$) from the nucleon core ($$J_C, M_C, T_1, T_{31}$$) and the interacting nucleon ($$j_N, m_N, t_N = \frac{1}{2}, t_{3N}$$) from the residual system ($$J_R, M_R, T_R, T_{3R}$$). Moreover, a sum over the quantum numbers of the particles in the final state has to be performed. In terms of the total and relative momenta of the two-nucleon final state, $$\vec{P}_T = \vec{k}_1 + \vec{k}_2$$ and $$\vec{k}_r = (\vec{k}_1 - \vec{k}_2)/2$$, and working in a coupled two-body spin-isospin basis, one has:

$$\sigma_N(J, M_J) = \int \frac{d^3P_T}{(2\pi)^3} \int \frac{d^3k_r}{(2\pi)^3} (2\pi) \delta(M_H - E_R - E_1 - E_2)$$

$$\times \sum_{SM_{j_\Lambda} M_{M_C} T_{3R} T_{3R}} \sum_{J_R M_R T_R T_{3R}} | \langle T_R T_{3R} | \frac{1}{2} t_{3N} | T_T T_{3T} \rangle |^2$$

$$\times \left| \sum_{TT_3} \langle TT_3 | \frac{1}{2} t_{3N} \sum_{m_{\Lambda M_C}} \langle j_\Lambda m_\Lambda, J_C M_C | J M_J \rangle \right|^2$$

(2)

$$\times \sum_{j_N S_{1/2}} \sum_{M_{R M_N}} \langle J_C T_1, J_R T_R, j_N t_{3N} | J_R M_R, j_N m_N | J_C M_C \rangle$$

$$\times \sum_{m_N m_{3N}} \langle j_N m_N | l_N m_{t_N}; \frac{1}{2} m_{s_N} \sum_{m_{\Lambda m_{3\Lambda}}} \langle j_\Lambda m_\Lambda | l_{\Lambda m_{\Lambda}}, \frac{1}{2} m_{s}\rangle$$

$$\times \sum_{S_0 M_{S_0}} \langle S_0 M_{S_0} | \frac{1}{2} m_{s\Lambda}, \frac{1}{2} m_{s_N} \rangle \sum_{T_0 T_{30}} \langle T_0 T_{30} | \frac{1}{2} t_{3N}; \frac{1}{2} t_{3N} \rangle$$

$$\times \frac{1 - (-1)^{L+S+T}}{\sqrt{2}}$$

$$\times t_{\Lambda N \rightarrow nN}(S, M_S, T, T_3, S_0, M_{S_0}, T_0, T_{30}, l_N, l_{\Lambda, \Lambda}, \vec{P}_T, \vec{k}_r)^2.$$  

In Eq. (2), $$M_H$$ stands for the mass of the initial hypernucleus, which is assumed to be at rest, $$E_1, E_2$$ and $$E_R$$ are the asymptotic total energies of the two
nucleons and the residual nucleus in the final state and $S^{1/2}(J_C T_I ; J_R T_R , J_N t_{3n})$ is a nucleon pick-up spectroscopic amplitude. The index $N = n$ or $p$ determines if the decay is induced by a neutron or a proton, with $t_{3n} = -1/2$ and $t_{3p} = 1/2$ correspondingly. The elementary amplitude $t_{\Lambda N \rightarrow nN}$ accounts for the transition from an initial $\Lambda N$ state with spin (isospin) $S_0 (T_0)$ to a final antisymmetric $nN$ state with spin (isospin) $S (T)$. It can be written in terms of other elementary amplitudes which depend on center-of-mass ($N_R, L_R$) and relative ($N_r, L_r$) principal and orbital angular momentum quantum numbers of the $\Lambda N$ and $nN$ systems:

$$t_{\Lambda N \rightarrow nN} = \sum_{N_r L_r N_R L_R} X(N_r L_r, N_R L_R, l_\Lambda l_N) t^{N_r L_r N_R L_R}_{\Lambda N \rightarrow nN},$$

where the dependence on the spin and isospin quantum numbers has to be understood. In Eq. (3), the coefficients $X(N_r L_r, N_R L_R, l_\Lambda l_N)$ are the well known Moshinsky brackets, while:

$$t^{N_r L_r N_R L_R}_{\Lambda N \rightarrow nN} = \frac{1}{\sqrt{2}} \int d^3 R \int d^3 r e^{-i \vec{P}_T \cdot \vec{R}} \Psi^*_{k_r} (\vec{r}) \chi^S \chi^T \mathbf{\hat{V}}_{\sigma, \tau} (\vec{r}) \Phi_{CM}^{N_R L_R} (\vec{R} / (b / \sqrt{2})) \Phi_{rel}^{N_r L_r} (\vec{r} / (\sqrt{2} b)) \chi^{S_0} \chi^{T_0},$$

with $\mathbf{\hat{V}}_{\sigma, \tau} (\vec{r})$ the weak transition potential depending on the relative coordinate between the interacting $\Lambda$ and nucleon, $r$, and their spin, $\sigma$, and isospin, $\tau$, variables. Moreover, $\Phi_{CM}^{N_r L_r} (\vec{r} / (\sqrt{2} b))$ and $\Phi_{rel}^{N_r L_r} (\vec{r} / (b / \sqrt{2}))$ are the relative and center-of-mass harmonic oscillator wave functions describing the $\Lambda N$ system, while $\Psi_{k_r} (\vec{r})$ and $e^{i \vec{P}_T \cdot \vec{R}}$ are the relative and center-of-mass wave functions of the $NN$ final state.

When extracting information on the elementary weak two-body interaction taking place in the medium, it is vital to account for the strong interaction between the hadrons in the initial and final states. The wave function $\Psi_{\vec{k}_r} (\vec{r})$, describing the relative motion of the two nucleons under the influence of a suitable $NN$ interaction, is obtained from the Lippmann-Schwinger equation. For the initial $\Lambda N$ system we start from a mean field approach where the $\Lambda$ and nucleon single particle wave functions are obtained from harmonic oscillator potentials. Their corresponding oscillator parameters have been adjusted to reproduce the $\Lambda$ separation energy in the hypernucleus under consideration and the charge form factor for the corresponding nuclear core. To obtain the correlated $\Lambda N$ wave function one should solve a similar equation as in the $NN$ case, but with the Pauli operator acting on the propagation of the intermediate states properly incorporated ($G-$matrix equation). A simpler approach, which has been tested in the weak decay of $^5\Lambda$He [1], consists in fitting microscopic $G-$matrix calculations with a phenomenological spin-independent
correlation function that multiplies the harmonic oscillator $\Lambda N$ wave function to obtain the correlated one. For the baryon-baryon strong interactions we take the Nijmegen soft-core model, version NSC97f [31], which has been used with success in hypernuclear structure calculations as well as in the decay of hypernuclei.

![Feynman diagrams](image)

Fig. 1. Feynman diagrams corresponding to the weak $|\Delta S| = 1 \Lambda N \rightarrow NN$ transition amplitudes mediated by the exchange of the pseudoscalar $\pi, \eta, K$ mesons and the vector $\rho, \omega, K^*$ mesons. The circle and the square stand for the weak and strong vertices, respectively.

Within our One-Meson-Exchange (OME) model, the weak transition potential is built from the exchange of virtual mesons belonging to the ground state pseudoscalar and vector octets, $\pi, \eta, K, \rho, \omega, K^*$, as depicted in Fig. 1. For the sake of simplicity, we will not derive here the expression for the transition potential starting from the weak and strong vertices entering each Feynman amplitude; details on this calculation can be found in Ref. [1]. Here we only quote the final result for the (non-relativistic) potential, which is:

$$V_{\sigma,\tau}(\vec{r}) = \sum_i \sum_{\alpha} V^{(i)}_{\alpha}(\vec{r}) = \sum_i \sum_{\alpha} V^{(i)}_{\alpha}(r) \hat{O}_{\alpha}(\vec{\sigma}, \vec{r}) \hat{I}^{(i)}_{\alpha}$$

$$= \sum_i [V^{(i)}_{C}(r) \hat{I}_{C}^{(i)} + V^{(i)}_{SS}(r) \vec{\sigma}_1 \cdot \vec{\sigma}_2 \hat{I}_{SS}^{(i)} + V^{(i)}_{T}(r) S_{12}(\hat{r}) \hat{I}_{T}^{(i)}$$

$$+ (n_1 \vec{\sigma}_2 \cdot \vec{r} + (1 - n_1)[\vec{\sigma}_1 \times \vec{\sigma}_2] \cdot \vec{r}) V^{(i)}_{PV}(r) \hat{I}_{PV}^{(i)}],$$

(5)

where $S_{12}(\hat{r}) = 3 \vec{\sigma}_1 \cdot \hat{r} \vec{\sigma}_2 \cdot \hat{r} - \vec{\sigma}_1 \cdot \vec{\sigma}_2$ is the tensor operator and $n_i = 1(0)$ for pseudoscalar (vector) mesons. The index $i$ runs over the different exchanged mesons and the index $\alpha$ over the different transition channels, central spin-independent, central spin-dependent, tensor and parity violating. Again, to avoid an excess of information, we refer to Ref. [1] for the explicit form of $V^{(i)}_{\alpha}(r)$, for the isospin factors $\hat{I}^{(i)}_{\alpha}$, as well as for the numerical values of the coupling constants.

In the present work we complement this OME potential with the contributions of uncorrelated ($2\pi$) and correlated ($2\pi/\sigma$) two-pion-exchange taken from Ref. [3]. In that work, a chiral unitary model has been used to account for the correlated two-pion-exchange in the scalar-isoscalar channel. This scheme was
built originally for the nucleon-nucleon interaction, leading to a $2\pi/\sigma$-exchange potential with a moderate attraction at $r \gtrsim 0.9$ fm and a repulsion at shorter distances [32], in contrast with the attraction at all distances of the standard phenomenological $\sigma$-meson exchange. Once the uncorrelated and correlated two-pion-exchange are added together, an attractive nucleon-nucleon potential is obtained for all distances. Applying an appropriate conversion factor that replaces a strong $\pi NN$ vertex with the weak $\pi \Lambda N$ one, the potential was implemented in the study of the weak decay of hypernuclei [3]. The relevant diagrams for uncorrelated and correlated two-pion-exchange with intermediate $N$ and $\Delta$ states built in Ref. [3] are depicted in Figs. 2 and 3, respectively. Two-nucleon intermediate states are not considered in the uncorrelated (direct and crossed) diagrams in order to avoid double counting when including the $NN$ strong correlations via the solution of a Lippmann-Schwinger equation, which includes the contribution of iterated one-pion-exchange interactions. It was found that, in the range of momenta relevant for the non-mesonic weak decay, the results from the uncorrelated two-pion diagrams with intermediate $\Delta N$ and $\Delta\Delta$ states are largely dominated by the isoscalar piece, which is the only one retained in their final results. As for the $2\pi/\sigma$ correlated contribution, the box in Fig. 3 contains the pion-pion scattering $t$-matrix summed up to all orders in the unitary approach. Diagrams containing intermediate $\Sigma$ and $\Sigma^*$ baryons are not included since their individual contributions approximately cancel each other when these baryons are considered together [3]. We also note that the complete scalar-isoscalar two-pion-exchange potential given in Ref. [3] is of pure parity-conserving nature. The reason is that the parity-violating contribution is strongly reduced by the lack of direct coupling of the $\Lambda$ to $\Delta$ intermediate states. The results of Ref. [3] show a large cancellation between correlated $2\pi/\sigma$ and uncorrelated $2\pi$-exchange at momentum values which are relevant for the non-mesonic decay. Consequently, the total two-pion-exchange contribution to the decay rate turns out to be small.

![Fig. 2. Uncorrelated two-pion exchange diagrams, direct and crossed, for the weak $|\Delta S| = 1 \Lambda N \rightarrow NN$ transition amplitude. Only $\Delta N$ and $\Delta\Delta$ intermediate states are allowed to avoid double counting (see explanation in the text).](image)

Calculations performed with the finite nucleus approach based on Eqs. (1) and (2) and adopting the OME potential previously mentioned reproduced quite well the $\Gamma_n/\Gamma_p$ values determined from data on coincidence nucleon spectra [6]. However, as all the other OME models employed to date, it failed in
accounting for the recent asymmetry data [20]. In the present paper we will see that the implementation of two-pion-exchange contributions modifies the decay widths moderately, as in Ref. [3], but has a tremendous influence on the decay asymmetries, bringing them to values that are in perfect agreement with the recent experimental data.

2.2 Asymmetry parameters

The study of polarized hypernuclei provides us with new and complementary phenomenological insights on the weak hyperon-nucleon interaction. Indeed, while investigations of the decay rates basically serve to clarify the isospin structure of the non-mesonic transitions, asymmetry studies allow us to extract significant information on the strength and the relative phases of the different decay amplitudes.

Using the reaction \( n(\pi^+, K^+)\Lambda \) on a \(^{12}\text{C} (^6\text{Li}) \) target at KEK [16,18] ([17,19]), with a pion momentum of \( \sim 1.05 \text{ GeV} \) and small kaon emission angles, \( 2^\circ \leq \theta_K \leq 15^\circ, \) \(^{12}\text{C} \) and \(^{14}\text{B} (^5\text{He}) \) hypernuclei have been produced with large spin-polarization, aligned preferentially along the axis normal to the reaction plane.

The expression for the intrinsic asymmetry parameter derived below follows the pioneering work developed in Ref. [15]. Neglecting for the moment nucleon final state interactions, the intensity of protons from \( \vec{\Lambda}p \to np \) decays emitted along a direction forming an angle \( \theta \) with respect to the hypernuclear spin-polarization axis is:

\[
I(\theta, J) = I_0(J) \left[ 1 + A(\theta, J) \right],
\]

with

\[
A(\theta, J) = P_y(J)A_y(J) \cos \theta.
\]
In the above, $J$ is the hypernuclear total spin, $I_0$ the isotropic intensity for an unpolarized hypernucleus:

$$I_0(J) = \frac{1}{2J+1} \sum_{M_J} \sigma_p(J, M_J) \equiv \Gamma_p \ .$$  \hspace{1cm} (8)

$P_y$ the hypernuclear polarization, which depends on the kinematics and dynamics of the hypernuclear production reaction, and $A_y$ the hypernuclear asymmetry parameter:

$$A_y(J) = \frac{3}{J+1} \sum_{M_J} \frac{M_J \sigma_p(J, M_J)}{\sum_{M_J} \sigma_p(J, M_J)} \ .$$  \hspace{1cm} (9)

The partial decay rates $\sigma_p(J, M)$ entering Eqs. (8) and (9) are defined by Eq. (2). The shell model weak-coupling scheme, in which the $1s_{1/2} \Lambda$ is assumed to be coupled to the nuclear core ground state with total spin $J_C$, allows rewriting the asymmetry $A$ in terms of the polarization of the hyperon spin, $p_\Lambda$, together with the corresponding intrinsic $\Lambda$ asymmetry parameter, $a_\Lambda$:

$$p_\Lambda(J) = \begin{cases} -\frac{J}{J+1} P_y(J) & \text{if } J = J_C - \frac{1}{2} \\ P_y(J) & \text{if } J = J_C + \frac{1}{2} \end{cases} \ .$$  \hspace{1cm} (10)

$$a_\Lambda = \begin{cases} -\frac{J+1}{J} A_y(J) & \text{if } J = J_C - \frac{1}{2} \\ A_y(J) & \text{if } J = J_C + \frac{1}{2} \end{cases} \ .$$  \hspace{1cm} (11)

In this way, Eq. (7) becomes:

$$A(\theta, J) = p_\Lambda(J) a_\Lambda \cos \theta \ ,$$  \hspace{1cm} (12)

and $a_\Lambda$ can be interpreted as being an intrinsic attribute of the elementary $\bar{\Lambda}p \to np$ process, in which case it should be practically independent of the decaying hypernucleus. Several calculations [1, 4, 12, 13, 15, 20–22, 25] have indeed demonstrated that this asymmetry shows only a moderate dependence on the hypernuclear structure.

Nucleon final state interactions strongly modify the weak decay intensity of Eqs. (6) and (7) and the experimentally accessible quantity is an observable proton intensity of the form [20]:

$$I^M(\theta, J) = I^M_0(J)[1 + p_\Lambda(J) a_\Lambda^M(J) \cos \theta] \ .$$  \hspace{1cm} (13)

The corresponding observable asymmetry is thus obtained from the measured
or calculated intensity as:

$$a^M_N(J) = \frac{1}{\rho_N(J)} \frac{I^M(0^\circ, J) - I^M(180^\circ, J)}{I^M(0^\circ, J) + I^M(180^\circ, J)},$$

(14)

and in general depends on the considered hypernucleus and on experimental conditions such as the adopted proton detection threshold. From Eq. (14) it is evident that to determine experimentally $a^M_N$, a measurement of the hypernuclear polarization $P_y$ [see Eq. (10)] is required. Such a measurement has been possible for $^5\Lambda$He [17], but only theoretical evaluations of $P_y$ are available for $p$-shell hypernuclei [16, 18, 19].

The relation between intrinsic and observable asymmetries has been investigated for the first time in Ref. [20], where the Monte Carlo intranuclear cascade model of Ref. [33] has been used to account for nucleon final state interactions. This code is based on the following basic ingredients through which the kinematics of the emitted nucleons is generated: i) first, according to the values of $\Gamma_n$ and $\Gamma_p$ predicted by the adopted meson-exchange model, a random number generator decides if the decay is neutron or proton-induced, thus determining the charges of the weak decay nucleons; ii) based on the relevant density and momentum probability distributions, the same random number generator selects the positions and momenta of these primary nucleons; iii) next, they propagate under the influence of a local potential and are allowed to collide with the nucleons of the medium, thus producing, among other effects, the emission of secondary nucleons. Each Monte Carlo event generates a certain number of nucleons which leave the nucleus with definite momenta. With enough statistics, one can then build up single nucleon spectra, coincidence spectra or up-down proton asymmetries that are directly comparable with the experimental observations.

## 3 Results

The weak decay observables predicted for $^5\Lambda$He, $^{11}\Lambda$B and $^{12}\Lambda$C are compared with recent data obtained at KEK in Table 1. Concerning the ratio $\Gamma_n/\Gamma_p$, only experimental results from nucleon-nucleon coincidence experiments [9–11] are quoted. These data should be preferred over the ones obtained from single-nucleon studies; the former are less affected by nucleon final state interactions and two-nucleon-induced decays than the latter [7]. In Table 1 we also list theoretical determinations from studies of nucleon final state interactions, with or without the inclusion of the two-nucleon induced decay mode. They have been obtained in Ref. [6] by fitting data on nucleon-nucleon spectra from Refs. [9–11]. Note that the determinations of $\Gamma_n/\Gamma_p$ by Ref. [6] are sometimes significantly smaller than the corresponding experimental results quoted in
Table 1
The non-mesonic weak decay rates (in units of the free Λ decay width) and intrinsic asymmetry parameters predicted for $^5Λ$He, $^{11}Λ$B and $^{12}Λ$C are compared with recent data. See text for details.

| Model         | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_{NM} = \Gamma_n + \Gamma_p$ | $\Gamma_n/\Gamma_p$ | $\alpha_\Lambda$ |
|---------------|-------------|-------------|-------------------------------------|---------------------|------------------|
| OME           | 0.122       | 0.257       | 0.379                               | 0.474               | −0.590           |
| $\pi$         | 0.040       | 0.189       | 0.460                               | 0.095               | −0.231           |
| $\pi + K$     | 0.097       | 0.329       | 0.450                               | 0.510               | −0.544           |
| $\pi + K + 2\pi$ | 0.121     | 0.285       | 0.396                               | 0.390               | +0.114           |
| OME + $2\pi + 2\pi/\sigma$ | 0.111 | 0.275       | 0.388                               | 0.415               | +0.041           |

KEK-E462 [9, 10, 18] 0.424 ± 0.024 0.45 ± 0.11 ± 0.03 0.11 ± 0.08 ± 0.04
KEK-E462 [19] 0.07 ± 0.08 ± 0.06
KEK-E462 [10] 0.39 ± 0.11 (1N)
(analysis of Ref. [6]) 0.26 ± 0.11 (1N + 2N)

| Model         | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_{NM} = \Gamma_n + \Gamma_p$ | $\Gamma_n/\Gamma_p$ | $\alpha_\Lambda$ |
|---------------|-------------|-------------|-------------------------------------|---------------------|------------------|
| OME           | 0.179       | 0.408       | 0.587                               | 0.439               | −0.809           |
| $\pi$         | 0.067       | 0.199       | 0.685                               | 0.108               | −0.353           |
| $\pi + K$     | 0.137       | 0.419       | 0.445                               | 0.147               | −0.773           |
| $\pi + K + 2\pi$ | 0.199     | 0.427       | 0.613                               | 0.438               | −0.074           |
| OME + $2\pi + 2\pi/\sigma$ | 0.202 | 0.425       | 0.627                               | 0.474               | −0.181           |

KEK-E508 [9, 18] 0.256 ± 0.035 0.46 ± 0.11 ± 0.03 0.11 ± 0.08 ± 0.04
KEK-E508 [19] −0.16 ± 0.28 ± 0.00
KEK-E307 [34] 0.26 ± 0.063 ± 0.073

| Model         | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_{NM} = \Gamma_n + \Gamma_p$ | $\Gamma_n/\Gamma_p$ | $\alpha_\Lambda$ |
|---------------|-------------|-------------|-------------------------------------|---------------------|------------------|
| OME           | 0.175       | 0.491       | 0.667                               | 0.357               | −0.698           |
| $\pi$         | 0.066       | 0.371       | 0.817                               | 0.088               | −0.350           |
| $\pi + K$     | 0.134       | 0.581       | 0.505                               | 0.363               | −0.643           |
| $\pi + K + 2\pi$ | 0.195     | 0.521       | 0.776                               | 0.335               | +0.007           |
| OME + $2\pi + 2\pi/\sigma$ | 0.194 | 0.529       | 0.722                               | 0.366               | −0.207           |

KEK-E508 [9, 11, 18] 0.34 ± 0.035 0.51 ± 0.13 ± 0.05 0.20 ± 0.26 ± 0.04
KEK-E508 [19] −0.16 ± 0.28 ± 0.00
KEK-E508 [11] 0.38 ± 0.14 (1N)
(analysis of Ref. [6]) 0.29 ± 0.14 (1N + 2N)
KEK-E307 [34] 0.828 ± 0.056 ± 0.066
Table 1. This signals the importance of final state interactions and two-nucleon induced decays —neglected [10] or accounted for in an approximate way [11] in experimental analyses— even when extracting the ratio from nucleon–nucleon coincidence observables.

We start recalling the results of the OME model, including the exchange of the mesons belonging to the ground state pseudoscalar and vector octets. These results slightly differ from those of Ref. [4] due to the use here of numerically improved correlated $NN$ wave functions. For the three hypernuclei under study, both the neutron-to-proton ratio and the total non-mesonic width are reasonably reproduced within the OME model, especially if one considers that non-negligible two-nucleon induced decay rates $\Gamma_2/(\Gamma_n + \Gamma_p) \simeq 0.20$ for $^5\Lambda$He and $\Gamma_2/(\Gamma_n + \Gamma_p) \simeq 0.25$ for $^{11}_\Lambda$B and $^{12}_\Lambda$C [6,30] should be taken into account as well. The values predicted for the intrinsic asymmetries are large and negative, whereas small results, compatible with zero, have been reported by recent experiments for all three hypernuclei.

The effects of uncorrelated $(2\pi)$ and correlated $(2\pi/\sigma)$ two-pion contributions are better visualized by including them, sequentially, to those of the lighter mesons ($\pi$ and $K$). As it is well know, the dominant tensor component in the one-pion-exchange mechanism disfavors neutron-stimulated decays and produces very small $\Gamma_n/\Gamma_p$ values. The addition of kaon-exchange reduces $\Gamma_{NM}$ by about 40% while increasing $\Gamma_n/\Gamma_p$ to values compatible with data. This result is also well-known, being mainly due to i) the enhancement of the parity-violating $\Lambda N(^3S_1) \rightarrow nN(^3P_1)$ transition contributing especially to neutron-induced decays and ii) the reduction of the tensor component, which for kaon-exchange has opposite sign of the one for pion-exchange. The size of the asymmetry is doubled and practically reaches the large value of the OME model. In fact, for all observables, the pion- plus kaon-exchange contributions already constitute a large fraction of the OME result.

As expected from the size of their respective potentials, see Fig. 14 of Ref. [3], the uncorrelated two-pion-exchange mechanism has a much larger influence than the correlated one. We observe that the $2\pi$ contribution increases $\Gamma_p$ substantially and $\Gamma_n$ more moderately, hence giving rise to a decrease of the $\Gamma_n/\Gamma_p$ ratio with respect to the $\pi + K$ result, which is especially sizable in the case of $^5\Lambda$He. The $2\pi/\sigma$ contribution affects the partial rates mildly, reducing $\Gamma_n$ by less than 10% and $\Gamma_p$ by slightly more than 10%. As it could be reasonably expected on the basis of the masses of the mesons included in the adopted weak transition potential, the exchange of two pions, both uncorrelated and correlated, turns out to be the most relevant mechanism beyond pion- and kaon-exchange, giving an appreciable contribution to the non-mesonic rates.

The effect of the two-pion exchange contribution is larger than that found in [3], where this potential was built and applied to the decay of hypernuclei
within a local density approximation approach. This is probably due to a different implementation of short range correlations. The use of a phenomenological $NN$ correlation function of the type $1 - j(q_r)$, with a cut-off momentum of $q_c = 780$ MeV, reduces the rates considerably [4]. We have checked that, in this situation, the relative changes induced by the two-pion scalar-isoscalar contributions amount to about half of those seen in the results of Table 1, which are obtained using realistic $NN$ wave-functions.

The most spectacular change induced by the uncorrelated two-pion mechanism is seen in the asymmetry parameter, which turns from being large and negative to being small and positive. Incorporating the $2\pi/\sigma$ mechanism brings some additional changes, basically tempering out the above mentioned effects. The remaining heavier mesons produce very moderate changes in the decay widths, while the asymmetry parameter, being built from interferences, shows a much stronger sensitivity. Roughly speaking, the incorporation of the scalar-isoscalar terms to the OME model leaves the rates basically unaltered, while reducing substantially the absolute value of the intrinsic asymmetry in such a way that the predictions for all weak decay observables are in excellent agreement with the measured values.

We note that a proper comparison with the observed asymmetries requires to account for the final state interactions of the weak decay nucleons as they go out of the residual nucleus, as done in Ref. [20]. However, before commenting on these effects below, we analyze the changes on the asymmetry in terms of the modifications induced by our isoscalar-scalar mechanism in the various transition amplitudes.

By applying appropriate projection operators to the elementary $\Lambda N \to nN$ potential, it is possible to select, from the hypernuclear transition amplitude, the contributions coming from specific spin-space transitions, $^{2S+1}L_J \to ^{2S'+1}L'_J$. For $^5$He, the resulting amplitudes are denoted by the capital letters $A, B, C, D, E, F$ in complete analogy with the notation $a, b, c, d, e, f$ used for the same amplitudes in the two-body case. In order to disentangle the contributions to the asymmetry coming from the various interferences, we also perform calculations for specific pairs of transitions, which will be denoted as $AE$, $BC$, $BD$, $CF$ and $DF$. Table 2 shows the size of each proton-induced decay amplitude, including its sign, for the OME and OME + $2\pi + 2\pi/\sigma$ models. The sum of the modulus squared of these amplitudes builds up the corresponding value of $\Gamma_p$. We also show the contribution of all possible interferences between pairs of amplitudes to the asymmetry. The sum of all these interferences produces the final result for the intrinsic asymmetry.

Except for the negligible $B(^1S_0 \to ^3P_0)$ amplitude, the other ones turn out to be of relevance in the determination of the proton decay asymmetry. In the case of the OME model, the parity–conserving amplitudes ($A$, $C$ and $D$) are negative,
Table 2
Hypernuclear amplitudes and interference terms in the proton-induced decay of $^5_\Lambda$He.

|   | Parity | Isospin | OME  | OME $+2\pi +2\pi/\sigma$ |
|---|--------|---------|------|-------------------------|
| A | $^1S_0 \rightarrow ^1S_0$ | PC | 1 | -0.1044 | +0.0835 |
| B | $^1S_0 \rightarrow ^3P_0$ | PV | 1 | +0.0057 | +0.0057 |
| C | $^3S_1 \rightarrow ^3S_1$ | PC | 0 | -0.1399 | +0.1480 |
| D | $^3S_1 \rightarrow ^3D_1$ | PC | 0 | -0.1814 | -0.1814 |
| E | $^3S_1 \rightarrow ^1P_1$ | PV | 0 | +0.3833 | +0.3833 |
| F | $^3S_1 \rightarrow ^3P_1$ | PV | 1 | +0.2234 | +0.2234 |

$\Gamma_p = \sum_{\alpha=A..F} |\alpha|^2$

|   | 0.257 | 0.275 |
|---|-------|-------|
| $AE$ | -0.2854 | +0.2112 |
| $BC$ | +0.0027 | -0.0033 |
| $BD$ | -0.0029 | -0.0027 |
| $CF$ | -0.0856 | +0.0405 |
| $DF$ | -0.2186 | -0.2046 |
| $\alpha_\Lambda$ | -0.590 | +0.041 |

and the parity-violating ones ($B$, $E$ and $F$) positive. We note that the larger contributions to the asymmetry turn out to be negative and correspond to the interferences between the $A$ and $E$ ($AE$), the $D$ and $F$ ($DF$), and the $C$ and $F$ ($CF$) amplitudes. Our two-pion scalar-isoscalar mechanism affects the parity conserving amplitudes which are diagonal in $S$ and $L$, namely $A$ and $C$. As we see, they even change their sign which, in turn, transform the negative interferences $AE$ and $CF$ into positive contributions that largely cancel the negative $DF$ interference. We note that the small reduction in magnitude of the $BD$ and $DF$ contributions to the asymmetry is just a reflection of the slight increase of the $\Gamma_p$ rate. As a consequence of the above mentioned change of sign, the asymmetry of $^5_\Lambda$He turns from being large and negative in the OME model to being slightly positive in the OME plus chiral $2\pi + 2\pi/\sigma$ model, in perfect agreement with the experimental observations.

Complementing the one-pion exchange mechanism in the weak decay of hypernuclei with that of two correlated pions, in the scalar ($2\pi/\sigma$) and vector ($2\pi/\rho$) sectors, was considered for the first time by Itonaga et al. [26]. The model was later extended to incorporate the exchange of the $\omega$ [5] and $K$ [27] mesons. We should note that the approximation scheme employed in these works is purely phenomenological. Their mechanism is such that, at the weak vertex, two pions are emitted via an intermediate $N$ or $\Sigma$ baryon. These two
Pions couple to a $\sigma$ meson, which is absorbed by a nucleon at the $\sigma NN$ strong vertex. The needed $\sigma$ mass, $m_\sigma$, and $\sigma NN$ coupling, $g_{\sigma NN}$, are taken from phenomenological fits of the strong $NN$ interaction, while the remaining unknown $\pi\pi\sigma$ coupling is fitted to reproduce the decay rates of $p$-shell hypernuclei. It is not clear whether this model includes in an effective way the uncorrelated two-pion mechanism. Another phenomenological and even simpler approach to $\sigma$-meson exchange is that adopted in Refs. [23, 25], where the strong $\sigma NN$ coupling constant is taken equal to that of pion, $g_{\sigma NN} = g_{\pi NN} \sim 13.2$, and the weak $\sigma \Lambda N$ vertex is parametrized in terms of a parity-conserving ($A_\sigma$) and parity-violating ($B_\sigma$) coupling constants that are adjusted to reproduce some weak decay observables. In Ref. [23], this $\sigma$ is added to one-pion-exchange, one-kaon-exchange and the direct quark transition induced by an effective four-quark Hamiltonian. In Ref. [25], the $\sigma$ is added to a meson-exchange model which is similar to the one considered here.

In contrast, the two-pion scalar-isoscalar contributions considered in the present work are theoretically well grounded in the sense that all the coupling constants are determined from chiral meson-meson and meson-baryon Lagrangians and by imposing SU(3) symmetry. The regularizing parameter is adjusted such that $\pi\pi$ scattering data is reproduced from threshold to around a center-of-mass energy of around 1 GeV, well beyond the $\sigma$ region. Having such different origin, it becomes difficult to perform a comparative analysis with the above phenomenological models. We will just point out some differences in the results.

The work of Ref. [27] found that to reproduce small and positive values of $a_\Lambda(\frac{5}{2}\Lambda \text{He})$, as experiment indicates, their $2\pi/\sigma$ potential [26] is too strong and must be decreased to half of its calculated value, hence producing an important reduction of the amplitude $A$. This in turn reduces the negative $AE$ term and the asymmetry gets dominated by their positive $F(C + D)$ term. It becomes clear that the way of obtaining a positive asymmetry in Ref. [27] is radically different from what it is found in the present work.

The one-pion and one-kaon exchange potential of Ref. [23] are modified by a $\sigma$-exchange contribution whose coupling constants are fitted to various observables in light nuclei. It is found that small values of the parity-violating coupling constant ($B_\sigma \sim 1$) and a large value of the parity-conserving one ($A_\sigma \sim 4$) would reproduce the measured $\Gamma_{\text{NM}}$ and $\Gamma_{n}/\Gamma_{p}$ in $\frac{5}{2}\Lambda \text{He}$, but the asymmetry would turn positive and large, of the order of 0.6. Another reasonable fit to the rates is found with $A_\sigma \sim -1.5$, but then the asymmetry is very large and negative, close to $-1$. This work concludes that the additional inclusion of the direct quark mechanism permits finding a solution that reproduces both the partial rates and the asymmetry, in which case the values $A_\sigma \sim 4$ and $B_\sigma \sim 6.6$ are found.

Qualitatively similar results are found in Ref. [25] in the sense that their
Table 3
Intrinsic and observable decay asymmetries predicted for $^5_\Lambda$He, $^{11}_\Lambda$B and $^{12}_\Lambda$C.

|                          | $^5_\Lambda$He | $^{11}_\Lambda$B | $^{12}_\Lambda$C |
|--------------------------|----------------|------------------|------------------|
| OME, $T_p^{th} = 0$ MeV  | -0.590         | -0.809           | -0.698           |
| FSI, $T_p^{th} = 0$ MeV  | -0.260         | -0.173           | -0.145           |
| FSI, $T_p^{th} = 30$ MeV | -0.401         | -0.400           | -0.340           |
| FSI, $T_p^{th} = 50$ MeV | -0.455         | -0.554           | -0.468           |
| OME $+ 2\pi + 2\pi/\sigma$, $T_p^{th} = 0$ MeV | +0.041 | -0.181 | -0.207 |
| FSI, $T_p^{th} = 0$ MeV  | +0.021         | -0.038           | -0.048           |
| FSI, $T_p^{th} = 30$ MeV | +0.028         | -0.111           | -0.126           |
| FSI, $T_p^{th} = 50$ MeV | +0.030         | -0.173           | -0.179           |
| EXP [18]                 | $0.11 \pm 0.08 \pm 0.04$ | $-0.20 \pm 0.26 \pm 0.04$ | |
| EXP [19]                 | $0.07 \pm 0.08^{+0.08}_{-0.00}$ | $-0.16 \pm 0.28^{+0.18}_{-0.00}$ | |

$\sigma$-exchange potential added to the full OME exchange model can fit $\Gamma_{NM}$ and $\Gamma_n/\Gamma_p$ but not the asymmetry. However, their solutions are intrinsically very different, since the PV strength of the $\sigma$ meson is dominant in Ref. [25] ($B_\sigma/A_\sigma \sim 10$ to 20) while the PC and PV $\sigma$ amplitudes are of comparable strength in Ref. [23] ($B_\sigma/A_\sigma \sim 1$).

Finally, we present in Table 3 our results for the asymmetry after incorporating the effects of final state interactions (FSI) on the emitted nucleons. These results are then directly comparable with the observed asymmetries. We show predictions for three hypernuclei and for the OME and OME+$2\pi + 2\pi/\sigma$ models. The first line in each case gives the asymmetry in the absence of FSI and without applying any cut on the kinetic energy of the emitted protons. The following lines incorporate the effect of FSI for different energy cuts, namely $T_p^{th} = 0, 30$ and $50$ MeV, to accommodate to the experimental conditions. A cut of $T_p^{th} \sim 30$ to $50$ MeV is applied to $^5_\Lambda$He data, while $T_p^{th} \sim 30$ MeV for $^{11}_\Lambda$B and $^{12}_\Lambda$C. We observe that, as in our OME study of Ref. [20], the incorporation of FSI reduces the magnitude of the observable asymmetry with respect to the intrinsic asymmetry and that, as the kinetic energy cut is increased, $a^{M}_\Lambda$ tends to recover the value of $a_\Lambda$. The results of Table 3 show that the OME model cannot reproduce the measured values of the asymmetries, while the additional incorporation of the $2\pi + 2\pi/\sigma$ mechanism provides asymmetry results in complete agreement with the data for all hypernuclei.
4 Conclusion

We have studied the non-mesonic weak decay of hypernuclei within a one-meson-exchange model supplemented with the contributions of the uncorrelated \( (2\pi) \) and correlated \( (2\pi/\sigma) \) two-pion-exchange mechanisms. These last mechanisms are based on a chiral unitary model which describes \( \pi\pi \) scattering data up to around 1 GeV and are taken from Ref. [3]. Our finite nucleus approach includes realistic strong correlations both in the initial and final states and considers the final state collisions of the nucleons in their way out of the residual nucleus.

We have found that the two-pion-exchange mechanisms modify moderately the partial decay rates but have a tremendous influence on the asymmetry parameter, due to the change of sign of some relevant amplitudes. The one-meson-exchange plus two-pion-exchange model turns out to be able to reproduce satisfactory, not only the total and partial hypernuclear weak decay rates, but also the asymmetries observed in the angular distribution of protons emitted by polarized hypernuclei.

Recent studies on the validity of the \( \Delta I = 1/2 \) isospin rule in the non-mesonic decay [23, 35–37] have been of large interest, especially due to their connections with the determination of \( \Gamma_n/\Gamma_p \) and the asymmetry parameter. Although this kind of studies should be warmly supported, here we have to note that, according to our results, based on pure \( \Delta I = 1/2 \Lambda N \rightarrow nN \) transitions, there appears to be no need for the introduction of \( \Delta I = 3/2 \) contributions to explain the observed non-mesonic decay rates and asymmetries.

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When comparing theoretical values of $a_\Lambda$ with data for $a_{M\Lambda}$ we assume that, as obtained in Ref. [20], $a_{M\Lambda}/a_\Lambda \simeq 0.7$ for $^5\Lambda$He and $\simeq 0.6$ for $^{11}\Lambda$B and $^{12}\Lambda$C.