Violation of the $\Delta I=1/2$ rule in the nonmesonic weak decay of hypernuclei

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Abstract

Violations of the $\Delta I=1/2$ rule are investigated in the nonmesonic weak hypernuclear decay using a weak $\Lambda N\rightarrow NN$ transition potential based on meson exchange. While the weak $\Delta I=3/2$ couplings of baryons to pseudoscalar mesons are known to be very small, the analogous couplings of baryons to vector mesons are not known experimentally. These couplings have been evaluated using the factorization approximation and are found to produce potentially significant changes in the predictions for hypernuclear decay observables in the case of $^{12}_\Lambda C$. Within the uncertainties of the factorization approximation we find that the proton-induced rate is affected by at most 10%, while the

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neutron-induced rate can change by up to a factor of two. The asymmetry parameter is strongly affected as well.
The observed dominance of strangeness-changing non-leptonic weak interactions by \( \Delta I=1/2 \) transition operators in the two processes studied to date (hyperon decay and \( K \to 2\pi, 3\pi \)) has often lead to the assumption that the \( \Delta I=1/2 \) rule will turn out to be a universal feature of all such non-leptonic interactions. It is, however, not clear that this is necessarily the case. Ref. [1], for example, from a quark model perspective, suggests the possibility of significant \( \Delta I=3/2 \) contributions to transitions from the \( ^1S_0 \) initial state in \( \Lambda N \to NN \). Similarly, relevant to meson exchange model treatments of \( \Lambda N \to NN \), it has been shown [2] that there may be significant \( \Delta I=3/2 \) contributions to the weak vector meson-baryon couplings. It is important to stress, in this regard, that there is a qualitative physical distinction between the weak baryon-pseudoscalar and weak baryon-vector meson couplings. This follows from the fact that the factorization contributions generated by those operators in the QCD-enhanced effective weak Hamiltonian associated with penguin graphs receive large (\( \approx 10 \)) enhancements in the case of the weak baryon-pseudoscalar couplings (due to the different chiral structure of the penguin-induced operators), in contrast to the case, for example, of the weak \( \Lambda N \rho \), \( \Sigma N \rho \), and \( \Sigma N \omega \) couplings, for which the corresponding factorization contributions actually vanish [2]. The observed \( \Delta I=1/2 \) dominance of the weak baryon-pseudoscalar couplings should, thus, not be taken as providing any evidence for a similar dominance of the baryon-vector meson couplings by \( \Delta I=1/2 \) operators. We will, in consequence, in what follows, attempt to make estimates of the sizes of the weak \( \Delta I=3/2 \) contributions to the \( \Lambda N \rho \), and \( NNK^* \) couplings relevant to \( \Lambda N \to NN \), and explore the influence of such terms on various hypernuclear decay observables.

The starting point for our estimate of the \( \Delta I=3/2 \) coupling contributions is the factorization approximation. As is well-known, QCD corrections to the basic weak interactions produce an effective weak Hamiltonian, \( H_{\text{eff}} \), which can be evaluated using perturbative QCD, down to a scale \( \sim 1 \) GeV. The form of \( H_{\text{eff}} \) for the non-leptonic strangeness changing weak interactions is then found to be [3,4]:

\[
H_{\text{eff}} = -\sqrt{2}G_F \sin \theta_C \cos \theta_C \sum_{i=1}^{6} c_i \mathcal{O}_i
\]

(1)
where $G_F$ is the Fermi constant, $\theta_C$ the Cabbibo angle and the operators $O_i$ have the form:

\[
O_1 = \bar{d}_L \gamma_\mu s_L \bar{u}_L \gamma^\mu u_L - \bar{u}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu u_L \\
O_2 = \bar{d}_L \gamma_\mu s_L \bar{u}_L \gamma^\mu u_L + \bar{u}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu u_L + 2 \bar{d}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu d_L + 2 \bar{d}_L \gamma_\mu s_L \bar{s}_L \gamma^\mu s_L \\
O_3 = \bar{d}_L \gamma_\mu s_L \bar{u}_L \gamma^\mu u_L + \bar{u}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu u_L + 2 \bar{d}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu d_L - 3 \bar{d}_L \gamma_\mu s_L \bar{s}_L \gamma^\mu s_L \\
O_4 = \bar{d}_L \gamma_\mu s_L \bar{u}_L \gamma^\mu u_L + \bar{u}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu u_L - \bar{d}_L \gamma_\mu s_L \bar{d}_L \gamma^\mu d_L \\
O_5 = \bar{d}_L \gamma_\mu \lambda^a s_L (\bar{u}_R \gamma^\mu \lambda^a u_R + \bar{d}_R \gamma^\mu \lambda^a d_R + \bar{s}_R \gamma^\mu \lambda^a s_R) \\
O_6 = \bar{d}_L \gamma_\mu \lambda^a s_L (\bar{u}_R \gamma^\mu u_R + \bar{d}_R \gamma^\mu d_R + \bar{s}_R \gamma^\mu s_R). 
\]

The Wilson coefficients, $c_i$, are scale-dependent and calculable perturbatively. The operators $O_1, \cdots, O_6$ in Eq. (2) have the specific (flavor, isospin) quantum numbers $(8, 1/2), (8, 1/2), (27, 1/2), (27, 3/2), (8, 1/2)$ and $(8, 1/2)$, respectively. The operators $O_5, O_6$, with LR chiral structure are generated by QCD penguin-type corrections and, as noted above, have different chiral structure than do the remaining operators. Of the operators, $O_1, \cdots, O_6$, only $O_4$ is $\Delta I=3/2$. An important feature of $O_4$ is that it is symmetric in the colors of the quark fields. Thus, the contributions to a $B'B M$ vertex (where $M$ is a meson and $B, B'$ are baryons) in which two quark fields from $O_4$ contract with two quarks from either the initial or final state baryon vanish in the quark model limit for baryon structure, in which the quark colors are antisymmetric. In this limit, therefore, one would expect the $\Delta I=3/2$ couplings to be saturated by the so-called “factorization contribution” graphs, in which one quark and one antiquark line from $O_4$ end up in the meson and the remaining quark and antiquark lines contract with a quark in the initial and a quark in the final state baryon, respectively. (This statement is not true for the $\Delta I=1/2$ operators, even in the quark model limit, and their matrix elements, therefore, have not only “factorization” contributions, but also non-factorization or “internal” contributions.) The factorization contributions, thus, involve baryon-to-baryon and vacuum-to-meson matrix elements of quark currents and these matrix elements can, of course, be evaluated either in terms of experimentally determined baryon form factors and meson decay constants, or in terms of form factors and decay constants.
related to those actually measured by $SU(3)_F$ arguments.

The “factorization approximation”, which we will employ in what follows, consists of assuming that, as in the quark model limit, the $\Delta I=3/2$ amplitudes of a non-leptonic strangeness changing process are given solely by their factorization contributions. The probable degree of accuracy of this approximation can be investigated by comparing “experimental” and factorization approximation values of the known hyperon decay amplitudes. We have put the word “experimental” in quotes here because it is, in fact, a non-trivial matter to go from the experimental amplitudes to the actual $\Delta I=3/2$ weak transition amplitudes associated with the operator $O_4$. This is because the strong interactions, as a consequence of the difference of $u$ and $d$ quark masses, themselves violate isospin, and this, combined with the fact that the basic $\Delta I=1/2$ transition amplitude is $\sim 20$ times larger than that for $\Delta I=3/2$ means that strong interaction mixing effects of only a few % could lead to significant enhancements or suppressions in the $\Delta I=3/2$ amplitudes. An example of such an effect is that associated with strong-interaction induced $\pi$-$\eta$ and $\Sigma^0$-$\Lambda$ mixing, which drastically alters the $\Delta I=3/2$ amplitudes for the p-wave $\Lambda \to N\pi$ and $\Xi \to \Lambda\pi$ transitions (by $\sim 400\%$ and $\sim 100\%$, respectively [5]). There would also be, in addition, strong isospin breaking effects associated with final state interactions which have not, to the best of our knowledge, been explored to date. Allowing only for particle mixing corrections, one finds that, for a scale of $\sim 1$ GeV, the factorization approximation for the $\Delta I=3/2$ amplitudes yields [5]:

1) good fits to the experimental values for the s- and p-wave $\Xi$ amplitudes and the s-wave $\Sigma$ triangle discrepancy;

2) an underestimate of the p-wave $\Sigma$ triangle discrepancy by a factor of 3-4;

3) an overestimate of the s- and p-wave $\Lambda$ amplitudes by a factor of 3-4.

It is important to stress, first, that the experimental errors on the $\Delta I=3/2$ amplitudes are large (the discrepancies above are $2\sigma$ at most) and, second, that even if we take these discrepancies seriously, we do not know whether the shortcoming represents a problem with the quark model approximation to baryon structure (i.e. the omission of “internal” $\Delta I=3/2$ contributions) or the effect of yet-to-be-corrected-for strong isospin breaking. We will, there-
fore, take the view that the factorization approximation provides an estimate for the \( \Delta I = \frac{3}{2} \) couplings of the vector mesons which can be expected to be reliable to within what we hope is a conservatively estimated error of a factor of 3-4.

From Eqs. (1) and (2) it is straightforward to obtain the factorization contributions to the \( \Delta I = \frac{3}{2} \) amplitudes, \( A^{(3)}(B \rightarrow B'V) \), relevant to the process \( \Lambda N \rightarrow NN \):

\[
A^{(3)}(\Lambda \rightarrow p\rho^-) = \frac{4}{3} c_4 K \langle \rho^- | \bar{u} L \gamma \mu d_L | 0 \rangle \langle p | \bar{u} L \gamma \mu s_L | \Lambda \rangle
\]

\[
A^{(3)}(p \rightarrow nK^{**}) = \frac{4}{3} c_4 K \langle K^{**} | \bar{s} L \gamma \mu d_L | 0 \rangle \langle n | \bar{d} L \gamma \mu u_L | p \rangle
\] (3)

and

\[
A^{(3)}(\Lambda \rightarrow n\rho^0) = -A^{(3)}(\Lambda \rightarrow p\rho^-) / \sqrt{2}
\]

\[
A^{(3)}(n \rightarrow nK^*) = -A^{(3)}(p \rightarrow nK^{**})
\]

\[
A^{(3)}(p \rightarrow pK^0) = A^{(3)}(p \rightarrow nK^{**})
\] (4)

where \( K = \sqrt{2} G_F \sin(\theta_C) \cos(\theta_C) \).

In Eqs. (3) and (4) we have omitted the weak \( \Lambda N \omega \) couplings since the \( \Delta I = \frac{3}{2} \) factorization contributions to these couplings vanish. The weak \( \Lambda N \rho \) couplings were obtained previously \[2\]; the \( K^* \) couplings are new. The relations in Eq. (4) follow from isospin Clebsch-Gordan coefficients. To complete our estimates for the \( \Delta I = \frac{3}{2} \) couplings we, therefore, need only to relate the baryon and meson matrix elements in Eq. (3) to observable meson decay constants and baryon form factors.

The meson-to-vacuum matrix elements are straightforward. We have:

\[
\langle 0 | V^3_\mu \rho(\epsilon, k) \rangle = f_\rho m^2_\rho \epsilon^\rho_\mu(\epsilon, k)
\]

\[
\langle 0 | V^{4-i5}_\mu K^*(\epsilon, k) \rangle = f_{K^*} m^2_{K^*} \epsilon^K_\mu(\epsilon, k)
\] (5)

where \( V^a_\mu \) is the usual \( SU(3)_F \) octet vector current and \( f_V \) the usual dimensionless vector meson decay constant. From \( \rho^0 \rightarrow e^+e^- \), \( f_\rho = 0.2 \) \[1\], while from \( \tau^- \rightarrow \nu_\tau K^{*-} \), \( f_{K^*} m^2_{K^*} \simeq 2f_\rho m^2_\rho \).
Similarly, dropping the second class baryon form factors \( f_3 \) and \( g_2 \), and the form factor \( g_3 \), which yields vanishing contribution when contracted against the transverse vector meson polarization vector, one finds that the baryon matrix elements can be written as:

\[
\langle B'(p')|V_\mu - A_\mu |B(p)\rangle \rightarrow \bar{u}_{B'}(p') \left[ f_1^{BB'} \gamma_\mu - \frac{i}{2m_N} f_2^{BB'} + g_1^{BB'} \gamma_\mu \gamma_5 \right] u_B(p). \tag{6}
\]

To obtain values for the baryon form factors not measured experimentally, we rely on \( SU(3)_F \) arguments. The form factors \( f_1^{BB'} \), \( f_2^{BB'} \) are then determined by CVC (the former in terms of the \( SU(3) \) structure constants, the latter in terms of the neutron and proton anomalous magnetic moments); the \( g_1^{BB'} \) form factors are similarly determined from the \( SU(3)_F \) analysis used for the axial vector hyperon decay amplitudes, which produces \( D, F \) factors \( D_A = 0.79 \) and \( F_A = 0.47 \). (For a clear discussion of this analysis, see, for example, Ref. [7].)

Combining all aspects of the above analysis and defining the weak \( \Delta I=3/2 \) \( B \rightarrow B'V \) vertex form factors via

\[
\langle B'(p')V(\epsilon,k)|\mathcal{H}_{eff}^{\Delta I=3/2}|B(p)\rangle = \epsilon_\mu^{(V)} \bar{u}_{B'}(p') \left[ f_1 \gamma_\mu - \frac{i}{2m_N} f_2 + g_1 \gamma_\mu \gamma_5 \right] u_B(p), \tag{7}
\]

in the factorization approximation, we find the following values for the weak form factors:

\[
f_1(\Lambda \rightarrow pp^-) = -\frac{1}{\sqrt{3}} c_4 K f_\rho m_\rho^2 = -1.2 \times 10^{-7}
\]
\[
f_2(\Lambda \rightarrow pp^-) = 1.63 f_1(\Lambda \rightarrow pp^-) = -1.9 \times 10^{-7}
\]
\[
g_1(\Lambda \rightarrow pp^-) = -0.72 f_1(\Lambda \rightarrow pp^-) = 0.85 \times 10^{-7} \tag{8}
\]

for the process \( \Lambda \rightarrow pp^- \), and:

\[
f_1(p \rightarrow nK^{*+}) = \frac{1}{3} c_4 K f_{K^{*+}} m_{K^{*+}}^2 = 1.4 \times 10^{-7}
\]
\[
f_2(p \rightarrow nK^{*+}) = 3.7 f_1(p \rightarrow nK^{*+}) = 5.1 \times 10^{-7}
\]
\[
g_1(p \rightarrow nK^{*+}) = 1.26 f_1(p \rightarrow nK^{*+}) = 1.7 \times 10^{-7} \tag{9}
\]

for the \( p \rightarrow n \) \( K^{*+} \) one, where, in obtaining the numerical values quoted, we have used \( c_4(1 \text{ GeV}) = 0.49 \) from Refs. [3,4].
In light of the discussion above, in exploring the potential consequences of including ∆I=3/2 contributions to the couplings, we will allow a variation of up to a factor of 3 about the central values quoted in Eqs. (8) and (9).

The ∆I=3/2 couplings derived above are incorporated in the hypernuclear weak decay model described in Ref. [8]. We then apply the modified model to the decays of $^3\Lambda\text{C}$ below. Enforcing the ∆I=1/2 rule, Ref. [8] developed a nonrelativistic ∆S=1 potential for the AN→NN transition in Λ-hypernuclei on the basis of a meson-exchange model. In addition to the long-ranged pion, the exchange of the other pseudoscalar mesons belonging to the octet, the η and the K, was included along with the vector mesons $\rho$, $\omega$ and $K^*$. The weak coupling constants for the parity-violating (PV) amplitudes were obtained by making use of soft-meson techniques plus SU(3) symmetry, which allows one to relate the unphysical amplitudes of the nonleptonic hyperon decays involving η’s and K’s to the physical B → B’ + π amplitudes. SU(6)$_W$, on the other hand, was used to relate the vector meson amplitudes to the pseudoscalar meson ones. For the evaluation of the parity-conserving (PC) amplitudes the pole model was employed.

The nonmesonic decay rate is directly related to the hypernuclear transition amplitude, in a manner which facilitates the transition from an initial hypernuclear state to a final state composed of two nucleons and a residual (A-2)-particle system. Using standard nuclear structure methods, the hypernuclear amplitude can be expressed in terms of two-body amplitudes, AN→NN. Short range AN correlations were accounted for in Ref. [8] through an appropriate correlation function based on the Nijmegen ΛN interaction [9], while the NN wave function was obtained by solving the scattering problem of two nucleons moving under the influence of the strong interaction.

The nonrelativistic reduction of the free space Feynman amplitude for the virtual meson exchange gives the transition potential in momentum space. For vector mesons the ∆I=1/2 potential reads:

$$V(q) = G_F m^2 \left( F_1 \hat{\alpha} - \frac{(\hat{\alpha} + \hat{\beta})(F_1 + F_2)}{4MM}(\sigma_1 \times q)(\sigma_2 \times q) \right)$$
\[ +i \frac{\hat{\epsilon}(F_1 + F_2)}{2M} (\sigma_1 \times \sigma_2) q \frac{1}{q^2 + \mu^2} \]  

(10)

where \( \mu \) is the mass of the meson, \( F_1 \) and \( F_2 \) are the strong vector and tensor couplings, for which we take those of the Nijmegen soft-core potential [3], and \( \hat{\alpha}, \hat{\beta} \) and \( \hat{\epsilon} \) operators which contain, in addition to the weak coupling constants listed in Table I (in units of \( G_F m^2 = 2.21 \times 10^{-7} \)) for the \( \rho \) and \( K^* \) mesons, the particular isospin structure. For the isovector \( \rho \)-meson the isospin dependence reads:

\[
\hat{\alpha}_\rho = f_1 \, \tau_1 \tau_2 \\
\hat{\beta}_\rho = f_2 \, \tau_1 \tau_2 \\
\hat{\epsilon}_\rho = g_1 \, \tau_1 \tau_2.
\]

(11)

where now \( f_1, f_2 \) and \( g_1 \) (denoted by \( \alpha_\rho, \beta_\rho \) and \( \epsilon_\rho \) in Ref. [8]) correspond to \( \Delta I=1/2 \) weak \( \Lambda N \rho \) coupling constants. For the isodoublet \( K^* \)-meson there are contributions proportional to \( \hat{1} \) and to \( \tau_1 \tau_2 \) that depend on the coupling constants, leading to the following dependence:

\[
\hat{\alpha}_{K^*} = \frac{C_{K^*}^{PC, V}}{2} + D_{K^*}^{PC, V} + \frac{C_{K^*}^{PC, T}}{2} \tau_1 \tau_2 \\
\hat{\beta}_{K^*} = \left( \frac{C_{K^*}^{PC, T}}{2} + D_{K^*}^{PC, T} \right) + \frac{C_{K^*}^{PC, T}}{2} \tau_1 \tau_2 \\
\hat{\epsilon}_{K^*} = \frac{C_{K^*}^{PV}}{2} + D_{K^*}^{PV} + \frac{C_{K^*}^{PV}}{2} \tau_1 \tau_2.
\]

(12)

Following Ref. [8] we also included a monopole form factor at each vertex.

The \( \Delta I=3/2 \) amplitudes can be straightforwardly included by assuming the \( \Lambda \) to behave like an isospin \( |3/2 - 1/2\rangle \) state and introducing an isospin dependence in the \( \Delta I=3/2 \) transition potential of the type \( \tau_3/2 \tau \), where \( \tau_3/2 \) is the \( 1/2 \rightarrow 3/2 \) isospin transition operator whose spherical components have the matrix elements:

\[
\langle 3/2 \, m' | \tau_3^{(i)} | 1/2 \, m \rangle = \langle 1/2 \, m \, 1 \, i \mid 3/2 \, m' \rangle \quad i = \pm 1, 0.
\]

(13)

The \( \Delta I=3/2 \) transition potential for the \( \rho \) and \( K^* \) mesons is then of the same form as the \( \Delta I=1/2 \) one but replacing in Eq. (10) the isospin operators by:
\[
\hat{\alpha} \rightarrow f_1 \, \tau_{3/2} \tau \\
\hat{\beta} \rightarrow f_2 \, \tau_{3/2} \tau \\
\hat{\varepsilon} \rightarrow g_1 \, \tau_{3/2} \tau ,
\]
\(\text{(14)}\)

for both the K* and the \(\rho\), where the weak \(\Delta I=3/2\) couplings \(f_1, f_2\) and \(g_1\), listed in Table I, are determined by comparing the values obtained from the operators in Eq. (14) with those for the specific transitions quoted in Eqs. (8) and (9).

Specifically, for the \(\rho\) meson, the operators in Eq. (14) give rise to the following product of isospin factors, coming from the weak and strong vertices:
\[
\sqrt{2/3} G_w \times 1 \quad \text{for the } \Lambda p \rightarrow np \text{ transition} \\
\frac{1}{\sqrt{3}} G_w \times \sqrt{2} \quad \text{for the } \Lambda p \rightarrow pn \text{ transition} \\
\sqrt{2/3} G_w \times (-1) \quad \text{for the } \Lambda n \rightarrow nn \text{ transition} ,
\]
\(\text{(15)}\)

where \(G_w\) stands for either \(f_1, f_2\) or \(g_1\). Eq. (15) demonstrates that, for example, the weak coupling contribution for the exchange term \(\Lambda p \rightarrow pn\) is given by \(G_w/\sqrt{3}\), indicating that the constants \(f_1, f_2\) and \(g_1\) (represented by \(G_w\)) must be defined as those in Eq. (8), corresponding to the process \(\Lambda \rightarrow p\rho^-\), but containing an additional factor \(\sqrt{3}\).

For the K* meson, simple Clebsch-Gordan algebra relations allow us to see that the transitions \(\Lambda p \rightarrow np\) and \(\Lambda p \rightarrow pn\) are identical, while \(\Lambda n \rightarrow nn\) has a relative change of sign. This can be easily seen from the relations (4) for the weak vertex and from the equality of the ANK* couplings in the strong sector. The operator \(\tau_{3/2} \tau\) of Eq. (14), connecting the initial and final states taking the \(\Lambda\) as a \(|3/2 - 1/2\rangle\) state, yields precisely this structure introducing an extra factor \(\sqrt{2/3}\). This global factor will be compensated by using weak couplings \((f_1, f_2\) and \(g_1))\) which are those of Eq. (9) multiplied by \(\sqrt{3/2}\).

Our results for \(^{12}\Lambda C\) are summarized in Table II. The observables include the nonmesonic decay rate, \(\Gamma_{nm}\) (in units of the free \(\Lambda\) decay rate, \(\Gamma_{\Lambda} = 3.8 \times 10^9 \text{s}^{-1}\)), the ratio of the neutron to proton induced decay rates, \(\Gamma_n/\Gamma_p\), which are also given separately, and the intrinsic asymmetry parameter, \(a_{\Lambda}\), which is related via simple angular momentum coupling
to the asymmetry parameter, \( A_p \). As mentioned above, the weak \( \Delta I=3/2 \) coupling constants are allowed to vary by up to a factor of \( \pm 3 \) to account for the limitations of the factorization model. Since the relative sign between the \( \Delta I=1/2 \) and \( \Delta I=3/2 \) amplitudes is not predicted within the factorization approach we allow for both possibilities.

As shown in Table II the total decay rate changes by at most 6% and lies within the error bars of the more recent experimental result [11]. However, more significant changes are seen for the ratio \( \Gamma_n/\Gamma_p \) which is enhanced by a factor of two for the combination \( (s_\rho, s_{K^*}) = (3, 3) \) and decreases by the same amount for the combination \( (s_\rho, s_{K^*}) = (-3, -3) \), where \( s_\rho \) and \( s_{K^*} \) denote the scaling factors mentioned above. This dramatic effect on the neutron-to proton-induced ratio is precisely of the nature expected once \( \Delta I=3/2 \) amplitudes are included. While not affecting the total rate the relative contributions of the vector mesons in a particular isospin channel are changed, thus modifying the \( \Gamma_n/\Gamma_p \) ratio. From Table II it is clear that the largest effect comes from the \( \Delta I=3/2 \) amplitudes of the \( K^* \) meson. This comes as no surprise since Ref. [8] already found the \( K^* \) to be the most important vector meson contribution.

Analyzing the changes in the neutron- to proton-induced ratio in more detail we find that its modifications are due almost entirely to changes in the neutron-induced decay rate, while the proton-induced channel is barely affected. Thus, the inclusion of \( \Delta I=3/2 \) amplitudes does not alleviate the discrepancy between theory and experiment for \( \Gamma_p \). The measured uncertainties in the \( \Gamma_n/\Gamma_p \) ratio are too large to draw any conclusions.

While the effects of the \( \Delta I=3/2 \) amplitudes on the ratio are significant, their effect on the asymmetry is even more dramatic. The two extreme values for the asymmetry parameter, which occur at \( (s_\rho, s_{K^*}) = (-3, -3) \) and \( (s_\rho, s_{K^*}) = (3, 3) \), differ by a factor of 7. In contrast to the neutron-to proton-induced ratio, both the \( K^* \) and the \( \rho \) contribute about equally to the change of this observable. Again, this is consistent with our findings of Ref. [8] where both the \( \rho \) and the \( \omega \) significantly affected the asymmetry. However, all the values of the hypernuclear asymmetry at \( 0^\circ \), \( A(0^\circ) \), obtained by multiplying \( a_\Lambda \) with the \( \Lambda \) polarization in \(^{12}\Lambda\)C and shown in the last column, are within the error bars of the experimental value.
Improved measurements using targets with higher hypernuclear polarization would be very beneficial to constrain the factorization approach used here.

We mention in passing that Ref. [14] has studied $\Delta I=3/2$ contributions to the nonmesonic decay in a direct quark mechanism, induced by four-point quark vertices in the effective weak Hamiltonian. As in Ref. [1], large $\Delta I=3/2$ contributions are found. These help to reproduce the neutron-induced decay rate. However, the sum of one-pion exchange plus the direct quark mechanism overpredicts the proton-induced rate. In a constituent quark picture such as that employed in Ref. [14], moreover, one would expect to have to include the long-range forces mediated by the exchange of all pseudoscalar mesons. As was pointed out long ago by Dubach and collaborators [15,16], and recently confirmed in a rigorous finite-nucleus calculation [8], there is considerable cancellation between one-pion-exchange and one-kaon-exchange contributions (the latter appearing through diagrams involving the weak NNK vertex, which contribution is omitted in Ref. [14]). This cancellation was also seen to be important in the results of Ref. [1], which included both one-pion-exchange and one-kaon-exchange, in addition to the direct quark contributions also present in Ref. [14].

In conclusion, we have studied the effect of the $\Delta I=3/2$ amplitudes in the meson-exchange model of Ref. [8]. Our results clearly indicate that both the neutron-induced rate and the asymmetry are very sensitive to the presence of $\Delta I=3/2$ amplitudes. Thus, better data for these observables will help in assessing the validity of the $\Delta I=1/2$ rule for the weak vector meson sector. Hopefully, such data will be forthcoming soon from the FINUDA facility at DAPHNE.

AP and AR acknowledge Joan Soto for clarifying discussions regarding the effective weak Hamiltonian. The work of CB was supported by US-DOE grant no. DE-FG02-95-ER40907 while the work of AP and AR was supported by DGICYT contract no. PB95-1249 (Spain) and by the Generalitat de Catalunya grant no. GRQ94-1022. This work has received support from the NATO Grant CRG 960132. AP acknowledges support from a doctoral fellowship of the Ministerio de Educación y Ciencia (Spain). KM acknowledges the ongoing support of the Natural Sciences and Engineering Research Council of Canada, and the hospitality...
of the Special Research Center for the Subatomic Structure of Matter of the University of Adelaide.
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TABLES

TABLE I. Weak coupling constants for the $\rho$ and K$^*$ exchange potential in units of $G_F m^2 = 2.21 \times 10^{-7}$

|       | $\Delta I = 3/2$ | $\Delta I = 1/2$ | $\Delta I = 3/2$ | $\Delta I = 1/2$ |
|-------|------------------|------------------|------------------|------------------|
| $f_1$ | -0.93            | -3.50            | 0.77             | $C_{K^*}^{PC,V} = -3.61$ |
|       |                  |                  |                  | $D_{K^*}^{PC,V} = -4.89$ |
| $f_2$ | -1.52            | -6.11            | 2.83             | $C_{K^*}^{PC,T} = -17.9$ |
|       |                  |                  |                  | $D_{K^*}^{PC,T} = 9.30$ |
| $g_1$ | 0.67             | 1.09             | 0.94             | $C_{K^*}^{PV} = -4.48$ |
|       |                  |                  |                  | $D_{K^*}^{PV} = 0.60$ |
TABLE II. Weak decay observables of $^{12}\Lambda$C for the full meson exchange potential

| $s_\rho$ $s_{K^*}$ | $\Gamma_{mm}/\Gamma_\Lambda$ | $\Gamma_n/\Gamma_p$ | $\Gamma_p/\Gamma_\Lambda$ | $a_\Lambda$ | $A(0^\circ)$ |
|---------------------|-----------------------------|---------------------|-----------------------------|-------------|-------------|
| $-3$ $-3$            | 0.783                       | 0.034               | 0.026                       | 0.758       | $-0.522$    | $-0.050$    |
| $-3$ 0               | 0.757                       | 0.061               | 0.043                       | 0.714       | $-0.414$    | $-0.039$    |
| $-3$ 3               | 0.789                       | 0.115               | 0.082                       | 0.708       | $-0.280$    | $-0.030$    |
| 0 $-3$              | 0.779                       | 0.040               | 0.030                       | 0.748       | $-0.434$    | $-0.041$    |
| 0 0                 | 0.753                       | 0.068               | 0.048                       | 0.705       | $-0.316$    | $-0.030$    |
| 0 3                 | 0.786                       | 0.123               | 0.086                       | 0.700       | $-0.178$    | $-0.017$    |
| 3 $-3$              | 0.787                       | 0.053               | 0.039                       | 0.747       | $-0.338$    | $-0.032$    |
| 3 0                 | 0.762                       | 0.081               | 0.057                       | 0.705       | $-0.212$    | $-0.020$    |
| 3 3                 | 0.796                       | 0.136               | 0.095                       | 0.701       | $-0.072$    | $-0.007$    |

EXP: $1.14\pm0.2$ \cite{10} $1.33^{+1.12}_{-0.81}$ \cite{10} $0.31^{+0.18}_{-0.11}$ \cite{11} $-0.01 \pm 0.10$ \cite{12}

$0.89\pm0.15\pm0.03$ \cite{11} $1.87\pm0.59^{+0.32}_{-1.00}$ \cite{11}

$0.70\pm0.3$ \cite{13}

$0.52\pm0.16$ \cite{13}