Nonmesonic Weak Decay of $\Lambda$ Hypernuclei within a Nuclear Matter Formalism

E. Bauer and F. Krmpotić

Departamento de Física, C.C. 67, Facultad de Ciencias Exactas,
Universidad Nacional de La Plata,
La Plata, 1900, Argentina

Abstract

The nonmesonic weak decay of $\Lambda$ hypernuclei using nonrelativistic nuclear matter is studied. As the basic building block we use the Polarization Propagator Method developed by Oset and Salcedo. It is shown that the exact calculation of exchange terms is required. Using the Local Density Approximation we evaluate the nonmesonic decay width for $^{12}_{\Lambda}C$ and compare the result with a finite nucleus calculation, obtaining a qualitative agreement.

PACS number: 21.80.+a, 25.80.Pw.

Keywords: $\Lambda$-hypernuclei, Non-mesonic decay of hypernuclei, $\Gamma_n/\Gamma_p$ ratio.
I. INTRODUCTION

The Λ decays within the hypernuclei via two mechanisms: the first one is called mesonic decay, where the product of the disintegration is a nucleon plus a pion. The mesonic decay rate $\Gamma_M \equiv \Gamma(\Lambda \to N\pi)$ is also the only process in the Λ free decay, $\Gamma^0$. The second mechanism is the nonmesonic decay, where the meson is absorbed by a nucleon. In this case the final product are two nucleons, $\Gamma_{NM} \equiv \Gamma(\Lambda N \to NN)$. Due to the Pauli principle the mesonic decay is strongly blocked for $A \geq 4$. At variance, this effect is not present in the nonmesonic channel. In principle, more than two nucleons can emerge in the nonmesonic decay. In this work we concentrate ourselves only on two nucleon emission. In this case, the corresponding transition rates are stimulated either by protons, $\Gamma_p \equiv \Gamma(\Lambda p \to np)$, or by neutrons, $\Gamma_n \equiv \Gamma(\Lambda n \to nn)$. The total nonmesonic decay rate is, $\Gamma_{NM} = \Gamma_n + \Gamma_p$. While theory fairly accounts for the experimental values for the total rate, the same is not true for the ratio $\Gamma_{n/p} \equiv \Gamma_n/\Gamma_p$, where theory underestimates data ($0.5 \leq \Gamma_{n/p}^{\text{exp}} \leq 2$).

Several models have been proposed to explain $\Gamma_{n/p}^{\text{exp}}$. A good review of the present status of the art can be found in Refs. [1] and [2]. Historically, Block and Dalitz (Refs. [3], [4]) developed a phenomenological model (see also Refs. [5]–[7]). From that point, microscopic models have been explored. The first one is due to Adams [8], who uses nuclear matter, one pion exchange model (OPE), $\Delta I = 1/2-\Lambda N\pi$ couplings and short range correlations (SRC). In the work of Adams, the $\Lambda N\pi$ coupling was too small to reproduce the Λ free lifetime. This mistake was corrected by McKellar and Gibson [9], whom also included the $\rho$-meson. Also using nuclear matter, Dubach et al. [10] introduced a one meson exchange model (OME) with $\pi$, $\eta$, $K$, $\rho$, $\omega$ and $K^*$-mesons.

Oset and Salcedo [11] developed the polarization propagator method (PPM). This scheme allows an unified treatment of mesonic and nonmesonic channels. It is performed in nuclear matter with the addition of the local density approximation (LDA). In the PPM one writes the expression for Λ self-energy where, in principle, propagators for nucleons and mesons can include correlations of all kinds. In Ref. [11] propagators are dressed with correlations.
of the Random Phase Approximation (RPA) type. The PPM is further developed in Refs. [12]-[16]. In addition of the above-mentioned mesons in Ref. [17] two-pion correlations are also considered. In a similar spirit of the PPM, Alberico and Garbarino [2] employed the bosonic loop expansion (BLE). These formalisms are particularly suitable when more than two nucleons emerge from the desintegration process.

Alternatively, some authors have employed finite nucleus wave functions instead of plane waves. This method is usually called Wave Function Method (WFM) (see Refs. [18]-[27]). The method makes use of shell model nuclear and hypernuclear wave function, as well as pion wave functions generated by pion-nucleus optical potential. Finally, let us mention that it was also included the quark degree of freedom to study this problem [26], [28]-[32].

This list of works does not pretend to be complete. We have tried to present the main physical ingredients considered in the literature to deal with the Λ-weak decay. Even though, the results for the ratio $\Gamma_{n/p}$ remains unsatisfactory. Some of the above mentioned issues are still in its preliminary stages and need further developments. Among which, we can mentioned the two-nucleon stimulated process (Refs. [12]-[15]), the inclusion of interaction terms that violates the isospin $\Delta I = 1/2$ rule (Refs. [21], [23], [32]), and the quark degree of freedom.

To deal with many body correlations like the RPA ones, a nuclear matter formalism is preferred over the WFM due to the difficulties in dealing with continuous wave functions within a WFM and also due to the big number of configurations involve. For this reason we have worked out a nuclear matter formalism. The first striking point is that most of the nuclear matter works usually neglect exchange terms or evaluated them in an approximate way. Due to this, we have explored a nuclear matter scheme which contains exchange terms and leads to results compatible with those of finite nucleus.

The present work is organized as follows. In Sect. II, an overview of the PPM is done and we present a nuclear matter formalism which includes exchange terms. In Sect. III we show results for the nonmesonic Λ decay together with a comparison with others authors. Finally, in Sect. IV some conclusions are drawn.
II. FORMALISM

We begin this section by briefly summarizing the main points of the PPM formalism introduced by Oset and Salcedo [11]. We first point out the advantages of the scheme, as well as its limitations. Afterwards we describe the formalism used here.

The PPM gives an unified description of both the mesonic and nonmesonic decay rates. Nonrelativistic nuclear matter is employed. The basic idea of the PPM is to evaluate the total decay rate $\Gamma_\Lambda = \Gamma_M + \Gamma_{NM}$, using the imaginary part of the $\Lambda$-self energy diagram of Fig. 1, as

$$\Gamma_\Lambda(k, k_F) = -2Im \Sigma_\Lambda.$$  \hspace{1cm} (1)

where $k$ and $k_F$ are the $\Lambda$ and Fermi momentum, respectively. The connection with $\Gamma_\Lambda$ is given by

$$\Gamma_\Lambda(k_F) = \int dk \Gamma_\Lambda(k, k_F) |\psi_\Lambda(k)|^2$$ \hspace{1cm} (2)

where for the $\Lambda$ wave function $\psi_\Lambda(k)$, we take the $1s_1/2$ wave function of a harmonic oscillator.

To evaluate $\Gamma_\Lambda$ for a particular nuclei one uses either an effective Fermi momentum or the Local Density Approximation (LDA) [11]. In the last case $k_F$ is spatially dependent and the transition rate reads

$$\Gamma_\Lambda = \int dr \Gamma_\Lambda(k_F(r)) |\tilde{\psi}_\Lambda(r)|^2$$ \hspace{1cm} (3)

where $\tilde{\psi}_\Lambda(r)$ is the Fourier transform of $\psi_\Lambda(k)$.

The $\Lambda N\pi$ vertex in Fig. 1 is described by the weak Hamiltonian

$$\mathcal{H}_{\Lambda N\pi} = iG_F m_\pi^2 \bar{\psi}_N (A_\pi + B_\pi \gamma_5) \tau \cdot \phi_\pi \psi_\Lambda + h.c.$$ \hspace{1cm} (4)

where $G_F m_\pi^2 = 2.21 \times 10^{-7}$ and the constants $A_\pi = 1.05$ and $B_\pi = -7.15$ are the parity violating and parity conserving couplings constants [21], respectively. In Eq. (4) we assume the $\Delta I = 1/2$ rule by taking the hyperon as an isospin spurion with $I_3 = -1/2$. The Hamiltonian for the strong $NN\pi$ vertex, which will be used latter, is given by
\[ \mathcal{H}_{\pi NN} = ig_{\pi NN} \bar{\psi}_N \gamma_5 \tau \cdot \phi_\pi \psi_N \] (5)

where the value of the strong-coupling constant is \( g_{\pi NN} = 13.3 \).

Using the standard Feynman rules, one gets in the nonrelativistic limit,

\[ \Sigma_\Lambda(k, k_F) = 3i(G_Fm_\pi^2)^2 \int \frac{d^4q}{(2\pi)^4} \left( A_2^2 + \frac{B_\pi^2}{4\bar{M}^2} q^2 \right) F_\pi^2(q) G_N(k - q) G_\pi(q), \] (6)

where the nucleon and pion propagators in the nuclear medium are, respectively,

\[ G_N(p) = \frac{\theta(|p| - k_F)}{p_0 - E_N(p) - V_N + i\varepsilon} + \frac{\theta(k_F - |p|)}{p_0 - E_N(p) - V_N - i\varepsilon}, \] (7)

and

\[ G_\pi(q) = \frac{1}{q_0^2 - q^2 - m_\pi^2 - \Sigma_\pi(q)}. \] (8)

Here, \( p = (p_0, \mathbf{p}) \) and \( q = (q_0, \mathbf{q}) \) denote the energy-momentum four-vector, \( E_N \) is the nucleon total free energy, \( V_N \) is the nucleon binding energy and \( \Sigma_\pi \) is the pion self-energy in nuclear matter. The constant \( \bar{M} \) is the average between the nucleon and \( \Lambda \) masses and \( k_F \) is the Fermi momentum. Plane waves for nucleons were employed in the derivation of Eq. (6), together with a step function which tells us if the nucleon is a particle or a hole. The finite nucleon size is shaped by the monopole form factor

\[ F_\pi(q) = \frac{\Lambda_\pi^2 - m_\pi^2}{\Lambda_\pi^2 - q_0^2 + q^2}, \] (9)

where \( m_\pi \) is the pion mass and \( \Lambda_\pi = 1.3 \text{ GeV} \).

The PPM is based on the behavior of the pion in the nuclear medium. The Feynman diagram displayed in Fig. 1 can be expanded in terms of Goldstone diagrams. In fact, this is an infinite series expansion, which contains RPA-diagrams, self-energy ones, etc. In Fig. 2, we show some of these terms, where the pion decays into a particle-hole pair (ph), a \( \Delta \)-hole pair (\( \Delta h \)), etc. Eventually, the ph or \( \Delta h \) can propagate within the nuclear medium. Each term has its own poles, and it is possible to evaluate the mesonic and nonmesonic contributions to the \( \Lambda \) decay separately. For example, when the pion self energy is neglected in Eq. (8), a pole in the meson propagator (which means that the meson is on the mass
shell) gives a contribution to the mesonic rate. On the other hand, in the nonmesonic decay rate the mesons are off the mass shell.

So far we have pointed out that the PPM is the sum of an infinite series of diagrams, containing both mesonic and nonmesonic contributions. Specific formulas and further details, can be found in Ref. [11]. From now on, we concentrate on the nonmesonic decay rate $\Gamma_{NM}$, considering only the contribution of the second diagrams displayed in Fig. 2, which is evaluated by means of the Goldstone rules. To obtain an analytical expression for $\Gamma_{NM}$ we should first specify the one pion exchange transition potential in momentum space

$$V_\pi(q) = G_F m_\pi^2 \frac{g_{\pi NN}}{2M} F_\pi^2(q) \left( \hat{A} + \frac{\hat{B}}{2M} \sigma_1 \cdot q \right) \frac{\sigma_2 \cdot q}{q_0^2 - q^2 - m_\pi^2},$$

which is the nonrelativistic reduction of Eqs. (4) and (5). Here $q$ is the momentum carried by the pion and $M$ is the nucleon mass. Note that we have added the form factor $F_\pi(q)$.

The operators $\hat{A}$ and $\hat{B}$, which contains the isospin dependence of the potential, are

$$\hat{A} = A_\pi \tau_1 \cdot \tau_2,$$

$$\hat{B} = B_\pi \tau_1 \cdot \tau_2.$$  

Let us consider analytical expression for nonmesonic decay rate from the above mentioned diagram. It is convenient to distinguish between the proton ($t_h = 1/2$) and neutron ($t_h = -1/2$) decay rates

$$\Gamma_{dir}^{th}(k, k_F) = -2 \text{Im} \int \frac{d^4 q}{(2\pi)^4} \int \frac{d^4 \kappa}{(2\pi)^4} G_N(\kappa + q/2) G_N(\kappa - q/2) G_N(k - q) \sum_{s_p t_p' s_{p'} t_{p'}} |s_p t_p s_{p'} t_{p'}| V_\pi(q) |s_{\Lambda t_{\Lambda s_{h} t_{h}}}|^2,$$

where the $s$'s and the $t$'s stand for the spin and isospin quantum numbers (see Fig. 2). When performing the energy integrations one keeps only the two particles-one hole (2p1h) cut, which gives the nonmesonic character to the transition rate. For the neutron decay we obtain,

$$\Gamma_{dir}^{n}(k, k_F) = \left( G_F m_\pi^2 \frac{g_{\pi NN}}{2M} \frac{1}{2\pi} \right)^2 \int dq \, \theta(q_0) \theta(|k - q| - k_F) \int dq_0 \, \theta(q_0) \theta(q_0 - q) F_\pi^2(q) \left[ A_\pi^2 + q^2 \left( \frac{B_\pi}{2M} \right)^2 \right] \frac{q^2}{(q_0^2 - q^2 - m_\pi^2)^2} \mathcal{U}(q_0, q),$$
while $\Gamma_{\nu}^{\text{dir}}(k, k_F) = 5\Gamma_{n}^{\text{dir}}(k, k_F)$. Here $q_0 = k_0 - E_N(k - q) - V_N$, with $k_0$ being the energy of the $\Lambda$ and

$$
\mathcal{U}(q_0, q) = \frac{2}{(2\pi)^2} \int d\kappa \ \theta(|\kappa + \frac{q}{2} - k_F|) \theta(k_F - |\kappa - \frac{q}{2}|) \\
\delta(q_0 - (E_N(\kappa + \frac{q}{2}) - E_N(\kappa - \frac{q}{2})))
$$

is the Lindhard function. The corresponding $\Gamma_{n,p}^{\text{dir}}$ are obtained from Eqs. (2) and (3).

At this point, we would like to discuss the limitations of the PPM. The first one refers to the incorporation of other mesons beyond the pion. By construction, in the PPM the only weak vertex is the $\Lambda N \pi$ one. The incorporation of other mesons is certainly important for the nonmesonic decay width, but it is not compatible with Eq. (6). Note that the naive attempt to solve this problem by replacing the pion propagator $G_\pi$ with the sum of other mesons propagators ($G_\pi \rightarrow G_\pi + G_{\rho} + G_{\eta} + ...$), fails to incorporate interference terms between mesons. The second point refers to the exchange terms, which are obviously not included in the PPM, as all diagrams are originated from the expansion of the dressed pion propagator of Eq. (8). In this work we develop a nuclear matter scheme, which overcome both just mentioned difficulties. We will limit our attention to the diagrams displayed in the Fig. 3, where the second graph of Fig. 2 is redrawn in the part a), while in the part b) we show it the corresponding exchange contribution.

We use the standard strangeness-changing weak $\Lambda N \rightarrow NN$ transition potential which involves the exchange of the complete pseudoscalar and vector meson octets ($\pi, \eta, K, \rho, \omega, K^*$). It was taken from Ref. [20] and the explicit expressions are listed in Appendix A. The incorporation of the short range correlations (SRC) is explained in Appendix B. For the sake of convenience, the total transition potential is written as

$$V_{\text{SRC}}(q) = \sum_{\tau=0,1} \mathcal{O}_\tau \mathcal{V}_\tau(q), \quad \mathcal{O}_\tau = \begin{cases} 1 \\ \tau_1 \cdot \tau_2 \end{cases}, \quad (16)$$

where

$$
\mathcal{V}_\tau(q) = (G_F m_\pi^2) \{ S_\tau(q) \sigma_1 \cdot \hat{q} + S'_\tau(q) \sigma_2 \cdot \hat{q} + P_{L,\tau}(q) \sigma_1 \cdot \hat{q} \sigma_2 \cdot \hat{q} + P_{C,\tau}(q) + P_{T,\tau}(q)(\sigma_1 \times \hat{q}) \cdot (\sigma_2 \times \hat{q}) + iS_{V,\tau}(q)(\sigma_1 \times \sigma_2) \cdot \hat{q}\}.
$$

(17)
The partial decay widths are defined as,

\[ \Gamma_{\tau} = \frac{\alpha^2}{4\pi} \int \frac{dq}{2\pi} \frac{1}{E_0^2 - |q| - k_F} \theta(|\kappa + q/2| - k_F) \delta(q_0 - (E_N(\kappa + q/2) - E_N(\kappa - q/2))) \]

for the direct contribution, and

\[ \Gamma^{exch}_{\tau} (\kappa, k_F) = \frac{\alpha^2}{4\pi} \int \frac{dq}{2\pi} \frac{1}{E_0^2 - |q| - k_F} \theta(\kappa - q/2) \delta(q_0 - (E_N(\kappa + q/2) - E_N(\kappa - q/2))) \]

for the exchange one, where \( Q = k - \kappa - q/2 \) and \( Q_0 = k_0 - E_N(\kappa + q/2) - V_N \). The integration over \( \kappa \) in Eq. (20) is factorized as the Lindhard function, which simplifies the
evaluation of $\tilde{\Gamma}_{\tau'\tau}^{\text{dir}}$. Summation over spin is already performed in $S_{\tau'\tau}^{\text{dir}}(q)$ and $S_{\tau'\tau}^{\text{exch}}(q,Q)$, which are defined as,

$$S_{\tau'\tau}^{\text{dir}}(q) = 4 \{ S_{\tau}(q) S_{\tau'}(q) + S_{\tau'}(q) S_{\tau}(q) + P_{L,\tau}(q) P_{L,\tau'}(q) + P_{C,\tau}(q) P_{C,\tau'}(q) +$$
$$+ 2 P_{T,\tau}(q) P_{T,\tau'}(q) + 2 S_{V,\tau}(q) S_{V,\tau'}(q) \} \quad (22)$$

and

$$S_{\tau'\tau}^{\text{exch}}(q,Q) = (\hat{q} \cdot \hat{Q}) S_{\tau}(q) S_{\tau'}(Q) + (2(\hat{q} \cdot \hat{Q})^2 - 1) P_{L,\tau}(q) P_{L,\tau'}(Q) +$$
$$+ P_{C,\tau}(q) P_{C,\tau'}(Q) + 2((\hat{q} \cdot \hat{Q})^2 - 1) P_{T,\tau}(q) P_{T,\tau'}(Q) -$$
$$- 2(\hat{q} \cdot \hat{Q})^2 (P_{L,\tau}(q) P_{T,\tau'}(Q) + P_{L,\tau}(Q) P_{T,\tau'}(q)) \quad (23)$$

where

$$S_{\tau}(q) S_{\tau'}(Q) = (S_{\tau}(q) + S_{\tau'}(q)) (S_{\tau'}(Q) + S_{\tau'}(q))$$
$$+ 2(S_{\tau}(q) S_{V,\tau'}(Q) + S_{V,\tau}(q) S_{\tau'}(Q))$$
$$- 2(S_{\tau'}(q) S_{V,\tau'}(Q) + S_{V,\tau}(q) S_{\tau'}(Q)). \quad (24)$$

The partial widths $\tilde{\Gamma}_{\tau'\tau}^{\text{dir,exch}}(k; k_F)$ depend on the momentum of $\Lambda$ and on the Fermi momentum, $k_F$. The $k_F$-dependence is eliminated by means of the LDA, as shown in Eq. (2), i.e.,

$$\tilde{\Gamma}_{\tau'\tau}^{\text{dir,exch}} = \int d\mathbf{k} |\psi_{\Lambda}(\mathbf{k})|^2 \tilde{\Gamma}_{\tau'\tau}^{\text{dir,exch}}(\mathbf{k}). \quad (25)$$

The final result from Eq. (18) is respectively,

$$\Gamma_n = \tilde{\Gamma}_{11}^{\text{dir}} - \tilde{\Gamma}_{11}^{\text{exch}} + \tilde{\Gamma}_{00}^{\text{dir}} - \tilde{\Gamma}_{00}^{\text{exch}} + \tilde{\Gamma}_{01}^{\text{dir}} - \tilde{\Gamma}_{01}^{\text{exch}} + \tilde{\Gamma}_{10}^{\text{dir}} - \tilde{\Gamma}_{10}^{\text{exch}}$$

$$\Gamma_p = 5 \tilde{\Gamma}_{11}^{\text{dir}} + 4 \tilde{\Gamma}_{11}^{\text{exch}} + \tilde{\Gamma}_{00}^{\text{dir}} + \tilde{\Gamma}_{00}^{\text{exch}} - (\tilde{\Gamma}_{01}^{\text{dir}} + \tilde{\Gamma}_{01}^{\text{exch}} + 2 \tilde{\Gamma}_{01}^{\text{dir}} + 2 \tilde{\Gamma}_{01}^{\text{exch}}). \quad (26)$$

Note that $\Gamma_{n/p} = 1/5$, when only the direct isovector contributions are considered. In the next section we give numerical results and also analyze the importance of different terms entering into our scheme.
In the evaluation of exchange term (diagram b of Fig. 3) Jido, Oset and Palomar [17] have approximated the momentum $Q$ by $-q$, which greatly simplifies the calculation. This implies that the exchange terms in Ref. [17] are approximated by direct ones, but with the spin-isospin factors corresponding to actual exchange diagrams. Simultaneously, they consider that the $\Lambda$ carries a non-vanishing $k$-momentum in both the direct and the exchange terms. This last point is somehow contradictory with $Q \approx -q$, as the later approximation is based on:  

1. the hyperon is considered to be at rest ($k = 0$), and 
2. the momentum of the hole is neglected ($\kappa - q/2 = 0$). 
To arrive to the same simplification from our scheme, one simply replaces $Q$ by $-q$ and $Q_0$ by $q_0$ in Eq. (23), which makes the quantity $S^{exch}_{\tau,\tau'}$ to depend only on $q$. As a further consequence, all factors $(\hat{q} \cdot \hat{Q})$ goes to -1, the term $P_{T,\tau}(q)P_{T,\tau'}(Q)$ disappears and $\kappa$-integral in Eq. (21) reduces to the Lindhard function.
III. RESULTS

In this section we give numerical values for the nonmesonic \( \Lambda \)-decay width. All calculations were done in nuclear matter with the transition potential presented in the last section and in Appendices A and B. The results for \( ^{12}\Lambda C \) comes from the LDA. The multiple integrations have been performed using a Monte Carlo technique. As mentioned in Sect. II, the hyperon is assumed to be in the \( 1s_{1/2} \) orbit of a harmonic oscillator well with frequency 
\[
\hbar \omega = (45A^{-1/3} - 25A^{-2/3}) \text{ MeV}.
\]

In order to analyze the importance of the exchange terms we show in Table I the numerical results for the neutron and proton decay widths. Two results are displayed for the total (direct plus exchange) decay rates,

- **Calculations I:** The simplification explained at the end of the last section has been implemented for the exchange term.
- **Calculations II:** The exchange contribution is evaluated in the exact way.

We start with the results for the OPE and then we add one by one the contributions of the remainder transition potentials. As expected the exchange terms are quite important. Furthermore, one sees that the total transition rates strongly depend on the way these terms are evaluated. The final result shows that the exchange terms increase the value of \( \Gamma_n \) while it has the opposite effect over \( \Gamma_p \), improving the ratio \( \Gamma_n/\Gamma_p \).

We have paid some attention to the role of the \( \rho \)-meson and the relative contributions of the parity violating (PV) and parity conserving (PC) decay widths. Since the work of McKeller and Gibson [9], there was a controversy referring to the importance of this meson. Yet in the work of Parreño et al. [20], it was established that when the \( \rho \)-meson is added to the pion the total rate is reduced by about 10-15 \%. In Table II, we compare our nuclear matter results with the finite nucleus calculation of Barbero et al. [27]. In the present calculation, the reduction of total rate is slightly bigger than in [20,27], although the overall agreement is rather good. It is worth noting that the present values for \( \Gamma_{n,p}^{PV} \) and \( \Gamma_{n,p}^{PC} \) differ from those of finite nucleus in the case of the pion, but for the \( \rho \)-meson the agreement is
satisfactory. While in finite nucleus calculations the pion PV contribution is about 40% of the total $\pi$-meson decay width [22,27], we get that it is only of about 23%. Note however, that the latter percentage is appreciable larger that in previous nuclear matter estimates: it is negligible in Ref. [9] and of the order of 15% in Ref. [10].

In Table III, we analyze the role of the $K$-meson, which improves the value of $\Gamma_n/\Gamma_p$. We compare our results with the nuclear matter calculations done in Refs. [26,17]. In the first work a partial wave expansion of the nuclear matter plane waves is done. The decay width is also evaluated in an approximate way: the summation over momentum of the two outgoing particles is performed with no restrictions (which means that they could take values below the Fermi momentum). The values for $\Gamma_n/\Gamma_p$ from Ref. [26] are in agreement with ours, but the individual transitions rates, $\Gamma_{n,p}$, are bigger. Regarding the second work, it should be stressed that the differences with our Calculation I are: 1) in Ref. [17] are also included the RPA correlations, and 2) the effective interaction is somewhat different. The second effect turns out to be the most relevant, as can be seen from the Calculation I', where $\Gamma_n$ and $\Gamma_p$ are evaluated by employing both the approximation and the interaction from Ref. [17], but without the RPA correlations. Finally, note that the inclusion of the kaon increases the ratio $\Gamma_n/\Gamma_p$ within the Calculation II as well.

In Table IV we compare our results for the full OME with those of Ref. [27]. They are quite similar, except for the vector mesons $\omega$ and $K^*$. One should keep in mind that the finite nucleus formalism of Ref. [27] has notable differences compared with the present nuclear matter model. Among the sources of difference, we can mention that we employ plane waves for both the incoming and outgoing wave functions, while in Ref. [27] harmonic oscillator wave functions for the incoming particles are used and the outgoing nucleons are expanded into partial waves.

Before ending this section we must call attention on the RPA-correlations. Many nuclear matter calculations dress the mesons propagators with RPA-type correlations. More precisely, very frequently only the direct RPA terms are considered, an estimation which is usually called ring approximation (RA). Within this framework a strong dependence of
the total nonmesonic decay width on the Landau-Migdal coupling $g'$ has been reported recently in Ref. [15] (see Fig. 3 of this work). Yet, it is well known that the RA leads to significantly different results for the electron scattering strength function than the full RPA [34,35]. Thus, we consider that it is encouraging to explore the consequences of the full RPA on the $\Lambda$-decay, which certainly is a complex issue and is beyond scope of the present work.

It is worthwhile to say a few words on the final state interactions (FSI), which is a very general denomination for all kind of interactions between the two outgoing particles. Parreño and Ramos [25] have treated them recently through the solution of a $T$-matrix using realistic NN interactions. Their results show that the FSI demand this kind of calculations over the phenomenological approach, which has been used in the present work.

As a final comment for this section, we wish to restate that there are several methods to relate the nuclear matter results to experimental data. Besides the LDA, we can mention the use of an effective Fermi momentum [36], and the employment of a diffused Fermi surface [37]. These last two approximations have been successfully employed in the context of the electron-nucleus quasi-elastic scattering. However, they lead to non-physical results for the mesonic decay, which is totally forbidden when the first method is employed and becomes artificially big when the second one is used [1]. These elements suggest that the LDA is a more adequate approximation for hypernuclei decays.
IV. CONCLUSIONS

A nuclear matter scheme for calculating the nonmesonic $\Lambda$-decay width has been presented, with explicit inclusion of exchange terms. It has been assumed that the transition is triggered by the full pseudoscalar-vector meson octet, with the corresponding form factors and short range correlations. To evaluate the decay rate of $^{12}_\Lambda C$ the LDA has been employed as well. Our numerical results were compared with finite nucleus ones and, except for the $\omega$ and $K^*$ mesons, good agreement was obtained.

At variance with finite nucleus calculations, the exchange terms are not always taken into account in the nuclear matter studies. In fact, the last ones can be classified in two groups, depending on whether the partial wave expansion of the nuclear matter plane waves is performed or not. In the first case, the Pauli principle is considered, but the $\Lambda$ is taken to be at rest and it is implicitly assumed that the exchange term carries the same momentum as the direct one (for details see Ref. [9]). Within the second group the most relevant formalism is, in our opinion, the PPM put forward by Oset and Salcedo [11]. The majority of works done within this model do not include the exchange term, which implies a separate and more complex calculation. An exception is Ref. [17], where they are incorporated in an approximate way (as stated at the end of Section II). Contrarily, we have evaluated them exactly, arriving to the conclusion that they are not only important but that they should also be calculated accurately.

Our numerical results agree fairly well with those obtained within the shell model framework [22,27]. Same as in these works, we are able to reproduce the data for the total nonmesonic decay width: $\Gamma_{NM}^{exp} \sim \Gamma_0$ [38], but not that for the $n/p$ ratio: $\Gamma_{n/p}^{exp} = 1.17^{+0.09+0.22}_{-0.08-0.18}$ [39]. This suggests that some others relevant physical ingredients are still missing. In this sense, our nuclear matter formalism is particularly suitable for: 1) analyzing the RPA correlations, and 2) the inclusion of the $\Lambda NN \rightarrow NNN$ decay [12–15].
ACKNOWLEDGMENTS

The authors are fellows of the CONICET Argentina, and acknowledge the support of ANPCyT (Argentina) under grant BID 1201/OC-AR (PICT 03-04296). One of us (E.B.) would like to thank A. Ramos for very helpful and illuminating discussions, and for a careful and critical reading of the manuscript.
APPENDIX A

In this appendix we show explicit expressions for the \((\eta, K, \rho, \omega \text{ and } K^*)-N\Lambda \rightarrow NN\) transition potential. The formulation was taken from [20], while the values of the different coupling constants and cutoff parameters appearing in the transition potential were taken from [33]. Weak couplings are in units of \(G_F m^2_\pi\).

For the pseudoscalar mesons we have expressions similar to Eq. (10) but making the following replacements,

\[
\begin{align*}
g_{N\pi} \rightarrow g_{N\eta}, \\
m_\pi \rightarrow m_\eta, \\
\hat{A} \rightarrow A_\eta, \\
\hat{B} \rightarrow B_\eta, \\
\end{align*}
\]

for the exchange of the isoscalar \(\eta\)-meson and

\[
\begin{align*}
g_{N\pi} \rightarrow g_{\Lambda K}, \\
m_\pi \rightarrow m_K, \\
\hat{A} \rightarrow \left(\frac{C_{PV}^K}{2} + D_{PV}^K + \frac{C_{PC}^K}{2}\tau_1 \cdot \tau_2\right) \frac{M}{M}, \\
\hat{B} \rightarrow -\left(\frac{C_{PC}^K}{2} + D_{PC}^K + \frac{C_{PV}^K}{2}\tau_1 \cdot \tau_2\right)
\end{align*}
\]

(27)

together with the exchange of index 1 and 2 in spin, for the isodoublet kaon. We employ, \(g_{NN\eta} = 6.4, A_\eta = 1.8, B_\eta = -14.3\) and \(\Lambda_\eta = 1.3\) GeV, for the \(\eta\)-meson. For the \(K\) meson, \(g_{\Lambda K} = -14.1, C_{PV}^K = 0.76, C_{PC}^K = -18.9, D_{PV}^K = 2.09, D_{PC}^K = 6.63\) and \(\Lambda_K = 1.2\) GeV.

In the case of vector mesons, we start with the \(\rho\)-meson,

\[
V_\rho(q) = G_F m^2_\pi (F_1 \hat{\alpha} - \frac{(\hat{\alpha} + \hat{\beta})(F_1 + F_2)}{4MM} (\sigma_1 \times q) \cdot (\sigma_2 \times q) - \frac{i\hat{\varepsilon} F_1 + F_2}{2M} (\sigma_1 \times \sigma_2) \cdot q) \frac{1}{q_0^2 - q^2 - m^2_\rho}
\]

(29)

where \(F_1 = g^V_{NN\rho}\) and \(F_1 = g^{T}_{NN\rho}\) and the operators \(\hat{\alpha}, \hat{\beta}\) and \(\hat{\varepsilon}\) are,
\[ \hat{\alpha} = \alpha_\rho \, \tau_1 \cdot \tau_2, \]
\[ \hat{\beta} = \beta_\rho \, \tau_1 \cdot \tau_2, \]
\[ \hat{\varepsilon} = \varepsilon_\rho \, \tau_1 \cdot \tau_2, \] (30)

with \( g_{N N \rho}^V = 3.16, g_{N N \rho}^T = 13.3, \alpha_\rho = -3.50, \beta_\rho = -6.11, \varepsilon_\rho = 1.09 \) and \( \Lambda_\rho = 1.4 \text{ GeV} \).

Finally, to obtain the \( \omega \) and \( K^* \) terms, one has to make the following substitutions in Eq. (29),

\[ m_\rho \to m_\omega, \]
\[ F_1 \to g_{N N \omega}^V, \]
\[ F_2 \to g_{N N \omega}^T, \]
\[ \hat{\alpha} \to \alpha_\omega, \]
\[ \hat{\beta} \to \beta_\omega, \]
\[ \hat{\varepsilon} \to \varepsilon_\omega \] (31)

and

\[ m_\rho \to m_{K^*}, \]
\[ F_1 \to g_{A N K^*}^V, \]
\[ F_2 \to g_{A N K^*}^T, \]
\[ \hat{\alpha} \to \frac{C_{K^*}^{PC,V}}{2} + D_{K^*}^{PC,V} + \frac{C_{K^*}^{PV,V}}{2} \tau_1 \cdot \tau_2 \]
\[ \hat{\beta} \to \frac{C_{K^*}^{PC,T}}{2} + D_{K^*}^{PC,T} + \frac{C_{K^*}^{PC,T}}{2} \tau_1 \cdot \tau_2 \]
\[ \hat{\varepsilon} \to \left( \frac{C_{K^*}^{PV}}{2} + D_{K^*}^{PV} + \frac{C_{K^*}^{PV}}{2} \tau_1 \cdot \tau_2 \right) \frac{M}{M^*}, \] (32)

with \( g_{N N \omega}^V = 10.5, g_{N N \omega}^T = 3.22, \alpha_\omega = -3.69, \beta_\omega = -8.04, \varepsilon_\omega = -1.33, \Lambda_\omega = 1.50 \) GeV, \( g_{A N K^*}^V = -5.47, g_{A N K^*}^T = -11.9 \) \( C_{K^*}^{PC,V} = -3.61, C_{K^*}^{PC,T} = -17.9, C_{K^*}^{PV} = -4.48, \)
\( D_{K^*}^{PC,V} = -4.89, D_{K^*}^{PC,T} = 9.30, D_{K^*}^{PV} = 0.60 \) and \( \Lambda_{K^*} = 2.20 \text{ GeV} \).
APPENDIX B

In momentum space the short range correlated (SRC) transition potential is obtained as,

\[ V_{SRC}(q) = V(q) - \int \frac{dp}{(2\pi)^3} \tilde{\xi}(|p + q|) V(p) \]  

(33)

where,

\[ \tilde{\xi}(p) = \frac{2\pi^2}{q_c^2} \delta(p - q_c) \]  

(34)

is the correlation function in momentum space. We have used \( q_c = 780 \). As an example let us show the result of Eq. (33) with the central part of the parity conserving one pion exchange potential, which we write in a simplify manner as,

\[ V^C_\pi(q) = C_\pi \frac{q^2}{q^2 + m^2_\pi} \sigma_1 \cdot \sigma_2 \tau_1 \cdot \tau_2 \]  

(35)

with \( C_\pi = -G_F m^2_\pi \frac{a_{NN\pi} \rho_{NN\pi}}{2M} \). Using this potential in Eq. (33) we obtain,

\[ V^{SRC,C}_\pi(q) = V^C(q) - C_\pi \frac{1}{2} (2 + \frac{m^2_\pi}{2q_c|q|} \ln \left| \frac{q_c^2 + m^2_\pi + q^2 - 2q_c|q|}{q_c^2 + m^2_\pi + q^2 + 2q_c|q|} \right|) \sigma_1 \cdot \sigma_2 \tau_1 \cdot \tau_2 \]  

(36)

if we call \( \kappa = 2q_c|q|/(q_c^2 + m^2_\pi + q^2) \) and now we use,

\[ \ln(1 + \kappa) \approx \kappa \]  

(37)

we finally obtain,

\[ V^{SRC,C}_\pi(q) = V^C(q) - C_\pi \frac{q_c^2 + q^2}{q_c^2 + m^2_\pi + q^2} \sigma_1 \cdot \sigma_2 \tau_1 \cdot \tau_2 \]  

(38)

which means that the contribution steaming from the second term of the r.h.s of Eq. (33) is simply \( V^C(q^2 \rightarrow q_c^2 + q^2) \). The procedure is analogous for rest of the interaction. We present now the final results of the short range correlated (\( \pi + \eta + K + \rho + \omega + K^* \))-transition potential. First, we define the following quantities,
\[
\begin{align*}
\mathcal{W}_\pi(q) &= \frac{g_{NN\pi}}{2M} \frac{B_\pi}{2M} F_\pi^2(q) G_\pi(q) \\
\mathcal{W}_\eta^S(q) &= \frac{g_{NN\pi}}{2M} \frac{A_\pi}{2M} F_\pi^2(q) G_\pi(q) \\
\mathcal{W}_\eta(q) &= \frac{g_{NN\eta}}{2M} \frac{B_\eta}{2M} F_\pi^2(q) G_\pi(q) \\
\mathcal{W}_\eta^S(q) &= \frac{g_{NN\eta}}{2M} \frac{A_\eta}{2M} F_\pi^2(q) G_\pi(q) \\
\mathcal{W}_K^0(q) &= -\frac{g_{ANK}}{2M} \frac{1}{2M} \left( \frac{C_{K^*}^{PC}}{2} + D_{K^*}^{PC} \right) \frac{M}{M} F_K^2(q) G_K(q) \\
\mathcal{W}_K^{S,0}(q) &= \frac{g_{ANK}}{2M} \left( \frac{C_{K^*}^{PV}}{2} + D_{K^*}^{PV} \right) F_K^2(q) G_K(q) \\
\mathcal{W}_K^1(q) &= -\frac{g_{ANK}}{2M} \frac{1}{2M} \frac{C_{K^*}^{PC}}{2} F_K^2(q) G_K(q) \\
\mathcal{W}_K^{S,1}(q) &= \frac{g_{ANK}}{2M} \frac{C_{K^*}^{PV}}{2} F_K^2(q) G_K(q) \\
\mathcal{W}_\rho^C(q) &= \alpha_\rho g_{NN\rho} F_\rho^2(q) G_\rho(q) \\
\mathcal{W}_\rho^T(q) &= -\frac{1}{4MM} (\alpha_\omega + \beta_\omega) (g_{NN\omega}^V + g_{NN\omega}^T) F_\rho^2(q) G_\rho(q) \\
\mathcal{W}_\rho^{PV}(q) &= \frac{\varepsilon_\rho (g_{NN\rho}^V + g_{NN\rho}^T)}{2M} F_\rho^2(q) G_\rho(q) \\
\mathcal{W}_\omega^C(q) &= \alpha_\omega g_{NN\omega}^V F_\omega^2(q) G_\omega(q) \\
\mathcal{W}_\omega^T(q) &= -\frac{1}{4MM} (\alpha_\omega + \beta_\omega) (g_{NN\omega}^V + g_{NN\omega}^T) F_\omega^2(q) G_\omega(q) \\
\mathcal{W}_\omega^{PV}(q) &= \frac{\varepsilon_\omega (g_{NN\omega}^V + g_{NN\omega}^T)}{2M} F_\omega^2(q) G_\omega(q) \\
\mathcal{W}_K^{C,0}(q) &= \left( \frac{C_{K^*}^{PC,V}}{2} + D_{K^*}^{PC,V} \right) g_{\Lambda K^*}^V F_K^2(q) G_K(q) \\
\mathcal{W}_K^{T,0}(q) &= -\frac{1}{4MM} \left( \frac{C_{K^*}^{PC,V}}{2} + D_{K^*}^{PC,V} + \frac{C_{K^*}^{PC,T}}{2} + D_{K^*}^{PC,T} \right) (g_{\Lambda K^*}^V + g_{\Lambda K^*}^T) F_K^2(q) G_K(q) \\
\mathcal{W}_K^{PV,0}(q) &= \frac{1}{2M} \left( \frac{C_{K^*}^{PC,V}}{2} + D_{K^*}^{PV} \right) (g_{\Lambda K^*}^V + g_{\Lambda K^*}^T) F_K^2(q) G_K(q) \\
\mathcal{W}_K^{C,1}(q) &= \frac{C_{K^*}^{PC,V}}{2} g_{\Lambda K^*}^V F_K(q) G_K(q) \\
\mathcal{W}_K^{T,1}(q) &= -\frac{1}{4MM} \left( \frac{C_{K^*}^{PC,V}}{2} + \frac{C_{K^*}^{PC,T}}{2} \right) (g_{\Lambda K^*}^V + g_{\Lambda K^*}^T) F_K^2(q) G_K(q) \\
\mathcal{W}_K^{PV,1}(q) &= \frac{1}{2M} \left( \frac{C_{K^*}^{PC,V}}{2} + D_{K^*}^{PV} \right) (g_{\Lambda K^*}^V + g_{\Lambda K^*}^T) F_K(q) G_K(q) \\
\end{align*}
\]  

(39)
we further introduce,

\[
G_i(q) = \frac{1}{q_0^2 - q^2 - m_i^2} \tag{40}
\]

\[
S_0 = (\mathcal{W}_\eta^S - \tilde{\mathcal{W}}_\eta^S)|q|
\]

\[
S_1 = (\mathcal{W}_\pi^S - \tilde{\mathcal{W}}_\pi^S)|q|
\]

\[
S_0' = (\mathcal{W}_{K,0}^S - \tilde{\mathcal{W}}_{K,0}^S)|q|
\]

\[
S_1' = (\mathcal{W}_{K,1}^S - \tilde{\mathcal{W}}_{K,1}^S)|q|
\]

\[
S_{V,0} = (\mathcal{W}_{\omega}^{PV} - \tilde{\mathcal{W}}_{\omega}^{PV} + \mathcal{W}_{K^*}^{PV,0} - \tilde{\mathcal{W}}_{K^*}^{PV,0})|q|
\]

\[
S_{V,1} = (\mathcal{W}_{\rho}^{PV} - \tilde{\mathcal{W}}_{\rho}^{PV} + \mathcal{W}_{K^*}^{PV,1} - \tilde{\mathcal{W}}_{K^*}^{PV,1})|q|
\]

\[
P_{L,0} = q^2(\mathcal{W}_\eta + \mathcal{W}_K^0) - (q^2 + \frac{1}{3}q_c^2)(\tilde{\mathcal{W}}_\eta + \tilde{\mathcal{W}}_K^0) - \frac{2}{3}q_c^2(\tilde{\mathcal{W}}_\omega + \tilde{\mathcal{W}}_{K^*}^{T,0})
\]

\[
P_{L,1} = q^2(\mathcal{W}_\pi + \mathcal{W}_K^1) - (q^2 + \frac{1}{3}q_c^2)(\tilde{\mathcal{W}}_\pi + \tilde{\mathcal{W}}_K^1) - \frac{2}{3}q_c^2(\tilde{\mathcal{W}}_\rho + \tilde{\mathcal{W}}_{K^*}^{T,1})
\]

\[
P_{T,0} = q^2(\mathcal{W}_\omega^T + \mathcal{W}_{K^*}^{T,0}) - (q^2 + \frac{2}{3}q_c^2)(\tilde{\mathcal{W}}_\omega^T + \tilde{\mathcal{W}}_{K^*}^{T,0}) - \frac{1}{3}q_c^2(\tilde{\mathcal{W}}_\eta + \tilde{\mathcal{W}}_K^0)
\]

\[
P_{T,1} = q^2(\mathcal{W}_\rho^T + \mathcal{W}_{K^*}^{T,1}) - (q^2 + \frac{2}{3}q_c^2)(\tilde{\mathcal{W}}_\rho^T + \tilde{\mathcal{W}}_{K^*}^{T,1}) - \frac{1}{3}q_c^2(\tilde{\mathcal{W}}_\pi + \tilde{\mathcal{W}}_K^1)
\]

\[
P_{C,0} = \mathcal{W}_\omega^C - \tilde{\mathcal{W}}_\omega^C + \mathcal{W}_{K^*}^{C,0} - \tilde{\mathcal{W}}_{K^*}^{C,0}
\]

\[
P_{C,1} = \mathcal{W}_\rho^C - \tilde{\mathcal{W}}_\rho^C + \mathcal{W}_{K^*}^{C,1} - \tilde{\mathcal{W}}_{K^*}^{C,1}
\]

\[
\tilde{\mathcal{W}}(q) = \mathcal{W}(q^2 \to q_c^2 + q^2) \tag{42}
\]

where the meaning of the tilde is,

the final expression for the interaction is given by Eq. (17)
REFERENCES

[1] E. Oset and A. Ramos, Prog. Part. Nucl. Phys. 41, 191 (1998).

[2] W. M. Alberico and G. Garbarino, nucl-th/0112036.

[3] R. H. Dalitz and G. Rajasekaran, Phys. Lett 1, 58 (1962).

[4] M. M. Block and R. H. Dalitz, Phys. Rev. Lett. 11, 96 (1963).

[5] C. B. Dover, Few-Body Systems Suppl. 2, 77 (1987)

[6] R. A. Schumacher, Nucl. Phys. A 547, 143c (1992).

[7] W. M. Alberico and G. Garbarino, Phys. Lett. B 486, 362 (2000).

[8] J. B. Adams, Phys. Rev. 156, 1611 (1967).

[9] B. H. J. McKellar and B. F. Gibson, Phys. Rev. C 30, 322 (1984).

[10] J. F. Dubach, G. B. Feldman, B. R. Holstein and L. de la Torre, Nucl. Phys. A 450 71c (1986).

J. F. Dubach, G. B. Feldman and B. R. Holstein, Ann. Phys. 249, 146 (1996)

[11] E. Oset and L. L. Salcedo, Nucl. Phys. A 443, 704 (1985).

[12] W. M. Alberico, A. De Pace, M. Ericson and A. Molinari, Phys. Lett. B 256, 134 (1991)

[13] A. Ramos, E. Oset and L. L. Salcedo, Phys. Rev. C 50, 2314 (1994).

[14] A. Ramos, M. J. Vicente-Vacas and E. Oset, Phys. Rev. C 55, 735 (1997).

[15] W. M. Alberico, A. De Pace, G. Garbarino and A. Ramos, Phys. Rev. C 61, 044314 (2000).

[16] W. M. Alberico, A. De Pace, G. Garbarino and A. Ramos, Nucl. Phys. A 668, 113 (2000).

[17] D. Jido, E. Oset and J. E. Palomar, Nucl. Phys. A 694, 525 (2001)
[18] J. Nieves and E. Oset, Phys. Rev. C 47, 1478 (1993).

[19] T. Motoba and K. Itonaga, Prog. Theor. Phys. Suppl. 117, 477 (1994).

[20] A. Parreño, A. Ramos and C. Bennhold, Phys. Rev. C 52, R1768 (1995) and ibid. 54, 1500 (1996).

[21] J. Golak, K. Miyagawa, H. Kamada, H. Witala, W. Glöckle, A. Parreño, A. Ramos and C. Bennhold, Phys. Rev. C 55, 2196 (1997).

[22] A. Parreño, A. Ramos and C. Bennhold, Phys. Rev. C 56, 339 (1997).

[23] A. Parreño, A. Ramos, C. Bennhold and Maltman, Phys. Lett. B 435, 1 (1998).

[24] A. Parreño, A. Ramos, N. G. Kelkar and C. Bennhold, Phys. Rev. C 59, 2122 (1999).

[25] A. Parreño and A. Ramos, Phys. Rev. C 65, 015204 (2002).

[26] K. Sasaki, T. Inoue and M. Oka, Nucl. Phys. A 669, 331 (2000); Erratum: ibid A 678, 455 (2000).

[27] C. Barbero, D. Horvat, F. Krmpotić, T. T. S. Kuo, Z. Narančić and D. Tadić, to be published.

[28] C. Y. Cheung, D. P. Heddle and L. S. Kisslinger, Phys. Rev. C 27, 335 (1983).

[29] D. P. Heddle and L. S. Kisslinger, Phys. Rev. C 33, 608 (1986).

[30] K. Maltman and M. Shmatikov, Phys. Lett B 331, 1 (1994); Nucl. Phys. A 585, 343c (1995).

[31] T. Inoue, S. Takeuki and M. Oka, Nucl. Phys. A 597, 563 (1996).

[32] T. Inoue, M. Oka, T. Motoba and K. Itonaga, Nucl. Phys. A 633, 312 (1998).

[33] M. N. Nagels, T. A. Rijiken and J. J. de Swart, Phys. Rev. D 15, 2547 (1977); P. M. M. Maessen, T. A. Rijiken and J. J. de Swart, Phys. Rev. C 40, 2226 (1989).
[34] E. Bauer, A. Ramos and A. Polls, Phys. Rev. C 54, 2959 (1996).

[35] E. Bauer and A. M. Lallena, Phys. Rev. C 59, 2603 (1999).

[36] J. E. Amaro, A. M. Lallena and G. Co’, Int. J. Mod. Phys. E 3, 735 (1994).

[37] E. Bauer and G. Co’, J. Phys. G 27, 1813 (2001).

[38] A. Montwill et al., Nucl. Phys. A 234, 413 (1974); J. J. Szymanski et al., Phys. Rev. C 43, 849 (1991); H. Noumi et al., Phys. Rev. C 52, 2936 (1995); H. Bhang et al., Phys. Rev. Lett. 81, 4321 (1998).

[39] O. Hashimoto et al., Phys. Rev. Lett. 88, 042503 (2002).
Table I: Proton and neutron decay widths for $^{12}\Lambda C$ in units of $\Gamma^0 = 2.52 \cdot 10^{-6}$ eV. The direct contributions are given in columns $\text{dir.}$, while the results for the total transition rates (direct plus exchange), obtained in the $\text{Calculations I}$ and $\text{II}$ (see text), are listed in columns $\text{Cal. I}$ and $\text{Cal. II}$, respectively.

| meson | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_{n/p}$ |
|-------|------------|------------|----------------|
|       | dir. Cal. I Cal. II | dir. Cal. I Cal. II | dir. Cal. I Cal. II |
| $\pi$ | 0.191 0.133 0.113 | 0.954 1.184 1.266 | 0.200 0.113 0.089 |
| $+\eta$ | 0.240 0.160 0.110 | 0.924 1.119 1.152 | 0.260 0.143 0.095 |
| $+K$ | 0.192 0.255 0.256 | 0.648 0.657 0.734 | 0.295 0.389 0.349 |
| $+\rho$ | 0.175 0.276 0.267 | 0.579 0.498 0.535 | 0.302 0.554 0.499 |
| $+\omega$ | 0.315 0.373 0.332 | 0.683 0.606 0.594 | 0.461 0.616 0.559 |
| $+K^*$ | 0.271 0.427 0.380 | 0.986 0.780 0.978 | 0.274 0.547 0.389 |
Table II: Contribution of the $\rho$-meson in the nonmesonic decay of $^{12}_\Lambda C$. $\Gamma^{PC}$ and $\Gamma^{PV}$ stand for the parity conserving and parity violating rates, respectively. Units are the same as in Table I.

|        | $\Gamma^{PC}_n$ | $\Gamma^{PV}_n$ | $\Gamma^{PC}_p$ | $\Gamma^{PV}_p$ | $\Gamma_\Lambda$ |
|--------|-----------------|-----------------|-----------------|-----------------|-----------------|
| Ref. [27] | 0.009           | 0.151           | 0.734           | 0.383           | 1.277           |
| Cal. II  | 0.005           | 0.108           | 1.004           | 0.262           | 1.379           |

|        | $\Gamma^{PC}_n$ | $\Gamma^{PV}_n$ | $\Gamma^{PC}_p$ | $\Gamma^{PV}_p$ | $\Gamma_\Lambda$ |
|--------|-----------------|-----------------|-----------------|-----------------|-----------------|
| Ref. [27] | 0.005           | 0.003           | 0.109           | 0.008           | 0.125           |
| Cal. II  | 0.007           | 0.003           | 0.100           | 0.012           | 0.122           |

|        | $\Gamma^{PC}_n$ | $\Gamma^{PV}_n$ | $\Gamma^{PC}_p$ | $\Gamma^{PV}_p$ | $\Gamma_\Lambda$ |
|--------|-----------------|-----------------|-----------------|-----------------|-----------------|
| Ref. [27] | 0.009           | 0.133           | 0.583           | 0.461           | 1.186           |
| Cal. II  | 0.004           | 0.128           | 0.727           | 0.204           | 1.063           |
Table III: Contribution of the $K$-meson. The cutoffs in Ref. [26] are: $\Lambda_\pi = \Lambda_K = 1.300$ MeV. Units are the same as in Table I.

|          | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_n/\Gamma_p$ |
|----------|------------|------------|----------------------|
| Ref. [26]| 0.221      | 2.354      | 0.094                |
| Ref. [17]| 0.119      | 0.956      | 0.124                |
| Cal. I   | 0.133      | 1.184      | 0.113                |
| Cal. I’  | 0.120      | 1.090      | 0.110                |
| Cal. II  | 0.113      | 1.266      | 0.089                |

|          | $\Gamma_n$ | $\Gamma_p$ | $\Gamma_n/\Gamma_p$ |
|----------|------------|------------|----------------------|
| Ref. [26]| 0.459      | 1.300      | 0.353                |
| Ref. [17]| 0.273      | 0.522      | 0.523                |
| Cal. I   | 0.217      | 0.697      | 0.311                |
| Cal. I’  | 0.223      | 0.485      | 0.460                |
| Cal. II  | 0.229      | 0.802      | 0.285                |
Table IV: Contributions of individual mesons to the decay width for $^{12}\Lambda C$. Units are the same as in Table I.

| meson  | $\Gamma_n$ (Ref. [27], Cal. II) | $\Gamma_p$ (Ref. [27], Cal. II) |
|--------|-------------------------------|-------------------------------|
| $\pi$  | 0.159 0.113 1.107 1.266       |                               |
| $\eta$ | 0.007 0.004 0.009 0.009        |                               |
| $K$    | 0.076 0.048 0.139 0.157        |                               |
| $\rho$ | 0.008 0.010 0.116 0.112        |                               |
| $\omega$ | 0.011 0.069 0.069 0.150       |                               |
| $K^*$  | 0.058 0.168 0.083 0.268        |                               |
| $\pi + \eta$ | 0.215 0.110 1.004 1.152 |                               |
| $\pi + K$ | 0.269 0.229 0.830 0.802  |                               |
| $\pi + \rho$ | 0.141 0.132 1.035 0.932 |                               |
| $\pi + \omega$ | 0.189 0.174 1.308 1.465 |                               |
| $\pi + K^*$ | 0.118 0.359 1.462 2.050 |                               |
| all mesons | 0.275 0.380 1.061 0.978 |                               |
FIGURES

FIG. 1. Λ self-energy in nuclear matter. Dot-dashed line represents a dressed-pion in nuclear matter. The continuous lines stand either for a nucleon or for the Λ (as indicated in the figure).

FIG. 2. A few lowest order terms for the Λ self-energy in nuclear matter. The dotted and wavy lines represent, respectively, the undressed pion and $NN$ strong interaction, while the $Δ$ excitation is denoted by the double continuous line.

FIG. 3. Direct (a) and exchange (b) contributions to the Λ decay width. The dashed-double dotted lines represent the full ($π + η + K + ρ + ω + K^*$)-transition potential.
Fig. 1
Fig. 2
Fig. 3