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Weak nonmesonic decay spectra of hypernuclei

Eduardo Bauer[a], Alfredo P. Galeão[b], Mahir S. Hussein[c], and Franjo Krmpotić[d]*

[a]Facultad de Ciencias Exactas, Departamento de Física, Universidad Nacional de La Plata, 1900 La Plata, Argentina

[b]Instituto de Física Teórica, UNESP - Universidade Estadual Paulista, Caixa Postal 70532-2, 01156-970 São Paulo, SP, Brazil

[c]Departamento de Física Matemática, Instituto de Física, Universidade de São Paulo, Caixa Postal 66318, 05315-970 São Paulo, SP, Brazil

[d]Instituto de Física La Plata, CONICET, Universidad Nacional de La Plata, 1900 La Plata, Argentina

We compute one- and two-nucleon kinetic-energy spectra and opening-angle distributions for the nonmesonic weak decay of several hypernuclei, and compare our results with some recent data. The decay dynamics is described by transition potentials of the one-meson-exchange type, and the nuclear structure aspects by two versions of the independent-particle shell model (IPSM). In version IPSM-a, the hole states are treated as stationary, while in version IPSM-b the deep-hole ones are considered to be quasi-stationary and are described by Breit-Wigner distributions.

The mesonic mode, $\Lambda \to N\pi$, with a rather small Q-value, $Q_M = M_\Lambda - M_N - m_\pi \approx 37$ MeV, is heavily inhibited for Λ-hypernuclei, except for the very lightest, due to Pauli blocking. With increasing mass number, A, a new mode quickly becomes dominant, namely the nonmesonic weak decay (NMWD), $\Lambda N \to nN$, whose Q-value, $Q_{NM} = M_\Lambda - M_N + \varepsilon_\Lambda + \varepsilon_N \approx 120 - 135$ MeV, is sufficiently large to render this Pauli blocking less and less effective. NMWD can be seen as one of the most radical transmutations of an elementary particle inside the nuclear medium: the strangeness is changed by $\Delta S = -1$ and the mass by $\Delta M = M_\Lambda - M_N = 176$ Mev. From a practical point of view, the main interest in NMWD is that it is, at present, the only way available to probe the strangeness-changing interaction between baryons.

Lately, the quality of experimental data on NMWD has improved considerably, and today one has available, not only one-nucleon kinetic energy spectra, but also two-nucleon coincidence spectra and opening-angle distributions obtained in several laboratories around the world, such as, KEK, FINUDA, and BNL. Here we briefly discuss a simple but fully quantum-mechanical formalism for the theoretical investigation of these observables. For more details see Refs. [1] and [2].

*E-mail address: krmpotic@fisica.unlp.edu.ar

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We start from Fermi's golden rule for the NMWD transition rate, $\Gamma_N \equiv \Gamma(\Lambda N \to nN)$, with the states of the emitted nucleons approximated by plane waves and initial and final short-range correlations implemented at a simple Jastrow-like level. The initial hypernuclear state is taken as a Λ in single-particle state $j_\Lambda = 1s_{1/2}$ weakly coupled to an $(A-1)$ nuclear core of spin J_C, i.e., $|J_I\rangle = |(J_C j_\Lambda)J_I\rangle$. For the description of nuclear states we adopt the independent particle shell model (IPSM).

Let us consider first the simplest version of this model, IPSM-a, in which all the relevant particle and hole states are assumed to be stationary. Thus, if the nucleon inducing the decay is in state j_N, then the possible states of the residual nucleus are $|J_F\rangle = |(J_C j_N^{-1})J_F\rangle$ and the liberated energy is $\Delta_{j_N} = M_\Lambda - M + \varepsilon_{j_\Lambda} + \varepsilon_{j_N}$, where the ε's are single-particle energies. Within this scheme, we get

$$\Gamma_N = 2\pi \sum_{S j_N J_F} \int\int |\langle \mathbf{p}_n \mathbf{p}_N S; J_F|V|J_I\rangle|^2 \delta(E_n + E_N + E_r - \Delta_{j_N}) \frac{d\mathbf{p}_n}{(2\pi)^3}\frac{d\mathbf{p}_N}{(2\pi)^3} \tag{1}$$

$$= \frac{4M^3(A-2)}{\pi} \sum_{j_N} \int dE_N \int dE_n \, \mathcal{F}_{j_N}(p,P), \tag{2}$$

where $E_{n,N} = \mathbf{p}_{n,N}^2/(2M)$ are the kinetic energies of the emitted nucleons, $E_r = |\mathbf{p}_n + \mathbf{p}_N|^2/[2(A-2)M]$ accounts for the recoil, and V is the transition potential. It is understood that all integrations on kinematical variables run over the allowed phase space. In Eq. (2),

$$\mathcal{F}_{j_N}(p,P) = \sum_{J=|j_N-1/2|}^{j_N+1/2} F_{j_N}^J \sum_{SlL\lambda J} |\langle plPL\lambda SJ|V|j_\Lambda j_N J\rangle|^2, \tag{3}$$

where $F_{j_N}^J$ are spectroscopic factors, p and P are the relative and total momenta of the emitted nucleons, l and L are the corresponding orbital angular momenta, and the couplings $\boldsymbol{l} + \boldsymbol{L} = \boldsymbol{\lambda}$, $\boldsymbol{\lambda} + \boldsymbol{S} = \boldsymbol{J}$ and $\boldsymbol{j}_\Lambda + \boldsymbol{j}_N = \boldsymbol{J}$ are performed. The one-nucleon transition probability density $S_N(E_N)$ is obtained by taking the derivative of Γ_N, i. e.,

$$S_N(E_N) = \frac{d\Gamma_N}{dE_N} = \frac{4M^3(A-2)}{\pi} \sum_{j_N} \int dE_n \, \mathcal{F}_{j_N}(p,P). \tag{4}$$

The proton and neutron spectra are then given by $\Delta N_p(E) \propto S_p(E)$ and $\Delta N_n(E) \propto S_p(E) + 2S_n(E)$. Similar developments, with the appropriate choice of kinematical integration variables, yield the coincidence energy spectra $\Delta N_{nN}(E) \propto S_{nN}(E)$ and the opening-angle distributions $\Delta N_{nN}(\cos\theta) \propto S_{nN}(\cos\theta)$, with $N = n,p$. More details, including the way to fix the normalization of these spectra and distributions, are given in Ref. [2].

For p-shell and heavier hypernuclei, some of the $|j_N^{-1}\rangle$ are deep-hole states, having considerable spreading widths, γ_{j_N}, as revealed for instance in quasifree $(p,2p)$ reactions. It is, therefore, unreasonable to treat such cases as stationary, zero-width, states. Rather, they are better approximated as Breit-Wigner distributions in the liberated energy ε, $\mathrm{P}_{j_N}(\varepsilon) = \frac{2\gamma_{j_N}}{\pi} \frac{1}{\gamma_{j_N}^2 + 4(\varepsilon - \Delta_{j_N})^2}$. This leads to a slightly more sophisticated version of the IPSM, which we call IPSM-b. It turns out that the final expressions we need can be obtained from those of IPSM-a through the replacements: $\Delta_{j_N} \mapsto \varepsilon$, and $\sum_{j_N} \cdots \mapsto \sum_{j_N} \int_{-\infty}^{+\infty} d\varepsilon \mathrm{P}_{j_N}(\varepsilon) \cdots$. This is explained in more detail, for the case of $S_{nN}(E)$, in Ref. [1].

The full one-meson exchange potential (OMEP), that comprises the $(\pi, \eta, K, \rho, \omega, K^*)$ mesons, has been employed for V in numerical evaluations. In Fig. 1 we confront the two IPSM approaches for the spectra $S_{np}(E)$. One sees that, except for the ground states, the narrow peaks engendered by the recoil effect within the IPSM-a become pretty wide bumps within the IPSM-b.

On the other hand, preliminary calculations indicate that the two IPSM versions yield similar results for $S_N(E)$ and $S_{nN}(\cos\theta)$. In Fig. 2 are compared the experimental [3] and theoretical kinetic energy spectra $\Delta N_p(E)$ for ${}^{12}_{\Lambda}$C, using the just mentioned OMEP. The theoretical spectrum is peaked around 85 MeV, and reproduces quite well the data for energies larger than 50 MeV. Yet, it differs quite a lot at smaller energies, where the effect of final state interactions (FSI) is likely to be quite important.

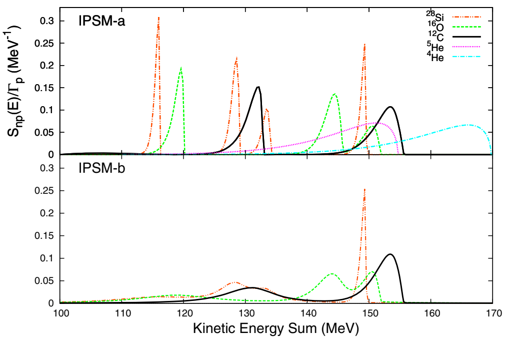

Figure 1. Normalized energy spectra $S_{np}(E)/\Gamma_p$ for ${}^{4}_{\Lambda}$He, ${}^{5}_{\Lambda}$He, ${}^{12}_{\Lambda}$C, ${}^{16}_{\Lambda}$O, and ${}^{28}_{\Lambda}$Si hypernuclei, taken from Ref. [1].

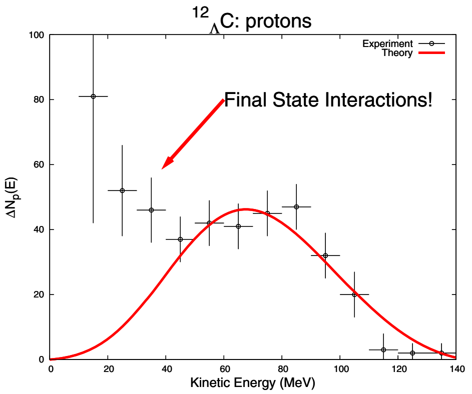

Figure 2. Comparison between the experimental [3] and theoretical kinetic energy spectra for protons from ${}^{12}_{\Lambda}$C decay.

The IPSM reproduces well [2] the BNL experiment for ${}^{4}_{\Lambda}$He [4], but it does not reproduce well the FINUDA experiment for the $S_N(E)$ spectra in ${}^{5}_{\Lambda}$He, ${}^{7}_{\Lambda}$Li, and ${}^{12}_{\Lambda}$C [3]. Once normalized to the transition rate, all the spectra are tailored basically by the kinematics of the corresponding phase space, depending very weakly on the dynamics governing the $\Lambda N \to nN$ transition proper. The IPSM is the appropriate lowest-order approximation for the theoretical description of the NMWD of hypernuclei. It is in comparison to this picture that one should appraise the effects of the FSI and of the two-nucleon-induced decay mode.

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