Nonmesonic weak decay of the hypertriton

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The nonmesonic weak decay of ${}^{3}_{\Lambda}$H is evaluated microscopically in the pion exchange model. The correlated three-body wave function of the hypertriton is approximated by a bound Λ-deuteron system obtained by averaging the YN interaction over the deuteron wave function. The relevant matrix elements are calculated in momentum space. The resulting decay rate is 4.9% of the free Λ decay rate.

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I. INTRODUCTION

The hypertriton ${}^{3}_{\Lambda}$H, consisting of a $pn\Lambda$ bound state, is the lightest hypernucleus since the ΛN interaction has no one-pion-exchange tail [1] due to the isospin-zero nature of the Λ. The role of ${}^{3}_{\Lambda}$H in hypernuclear physics [2] is therefore similar to the deuteron in conventional nuclear physics in providing an important testing ground for the underlying interactions.

All experimental information obtained for the hypertriton comes from early bubble chamber measurements [3] and emulsion works [4] more than 20 years ago. It has not been studied in the series of hypernuclear experiments over the last decade [5] since its production via the (π,K) or (K,π) mechanism requires either using a tritium target or detecting a neutral meson π^0 or K^0 in the final state. However, with the upcoming of CEBAF it will be possible to measure kaon photoproduction [6] on ^{3}He, namely, ${}^{3}\mathrm{He}(\gamma,K^+){}^{3}_{\Lambda}\mathrm{H}$, thus producing the hypertriton with a more convenient target and minimally interacting probes.

In this study we present a calculation for total and partial nonmesonic decay rates of the hypertriton, ${}^{3}_{\Lambda}\mathrm{H}\rightarrow p+n+n$. The mesonic decay modes like ${}^{3}_{\Lambda}\mathrm{H}\rightarrow {}^{3}\mathrm{He}+\pi^-$, etc. dominate the weak decay of ${}^{3}_{\Lambda}$H in contrast to heavier hypernuclei where the pionic channel is Pauli suppressed. However, the hypertriton is an ideal testing ground for the nonmesonic decay channel since the wave function in the initial state can in principle be calculated exactly using the Faddeev formalism or variational techniques. On the other hand, computations for the nonmesonic decay rates of heavier shell-model hypernuclei [7] require single-particle potentials, spectroscopic nuclear structure information, and a ΛN correlation function as input. All three ingredients are normally taken from independent sources and introduce model dependences that can affect the total rate by a factor of 2 or more [7]. These ingredients are naturally included in a correlated three-body wave function of ${}^{3}_{\Lambda}$H, thus allowing a less model-dependent study of the basic transition amplitude $\Lambda N\rightarrow NN$.

The first theoretical estimates of the ${}^{3}_{\Lambda}$H nonmesonic decay rates were made by Dalitz and collaborators [8] more than 30 years ago. They extracted partial rates from existing data for light Λ-hypernuclear weak decays, treating the Λ decay induced by protons or neutrons as incoherent. Along with an average density for the hypertriton they obtained total nonmesonic rates of less than 1% of the free Λ decay rate.

In Sec. II we briefly present the hypertriton wave function [9] employed in our calculation. Though not a solution of the Faddeev equations it should be sufficient for a first-order calculation of the process ${}^{3}_{\Lambda}\mathrm{H}\rightarrow n+n+p$. Section III outlines the formalism in momentum space. The elementary $\Lambda N\rightarrow NN$ transition amplitude is restricted to pion exchange only; however, heavier mesons can easily be included at a later stage. The discussion and summary of our findings is given in Sec. IV.

II. THE HYPERTRITON WAVE FUNCTION

A simple model for the hypertriton has been developed in Ref. [9] which can easily be applied to processes involving the production and decay of ${}^{3}_{\Lambda}$H. In this model the hypertriton is taken to be a deuteron and a lambda particle moving in an effective Λd potential. The influence of the Λ on the two nucleons is neglected, thus the nucleon part of the ${}^{3}_{\Lambda}$H wave function is that of the free deuteron. Even though the Λ separation energy of $B_\Lambda=0.13\pm0.05$ MeV is less than 6% of the deuteron binding, $B_d=2.22$ MeV, this is a rough approximation since it excludes the coupling to the $\Lambda N\rightarrow\Sigma N$ conversion channel [2]. While this is unrealistic for describing the full ΛN interaction, it may, on the other hand, be a reasonable approach when the ΣNN components in the hypertriton wave function are small. Reference [10] found a probability of 0.36% for the ΣNN component in ${}^{3}_{\Lambda}$H using the phenomenological YN coupled channel potential of Wyech [11].

A simple Λd potential is constructed by first performing a separable fit to the ΛN s-wave interaction which is given by the most recent version of the Nijmegen YN soft-core potential [12]. The potential is then spin averaged over the ΛN configurations found in the hypertriton. In the next step the ΛN potential is summed over the two nucleons and averaged over the deuteron wave function. Finally, the resulting Λd potential is again fitted by a separable form retaining only the s-wave part.

The experimental Λ separation energy can be reproduced by a very small fine tuning of the Λd potential parameters. We note that this simple wave function successfully reproduces [9] the ratio of ${}^3_\Lambda\mathrm{H}\to\pi^- + {}^3\mathrm{He}$ to ${}^3_\Lambda\mathrm{H}\to\pi^- +$ all.

The ${}^3_\Lambda$H wave function in its center-of-mass (c.m.) system is given by

$$\Psi_M(\mathbf{q},\mathbf{p})=\varphi_\Lambda(q)\sum_{(l,S)}\Psi^l_{\mathrm{deut}}(p)[\mathcal{Y}^l_{l0}(\hat{\mathbf{p}},\hat{\mathbf{q}})\chi^S_{1\frac{1}{2}}]^{1/2}_M\ , \quad (1)$$

where we have introduced the Jacobi coordinates for the three-body system in order to properly separate out the center-of-mass motion. Labeling the Λ as particle 1 we have in the c.m. system

$$\mathbf{K}=0\ ,\quad \mathbf{p}=\tfrac{1}{2}(\mathbf{k}_2-\mathbf{k}_3)\ ,\quad \mathbf{q}=\mathbf{k}_\Lambda\ . \quad (2)$$

In Eq. (1) $\Psi^l_{\mathrm{deut}}(p)$ denotes the radial part of the deuteron wave function, taken from the Paris potential [13]. The spin part

$$\chi^{SM_s}_{1\frac{1}{2}}=\sum_{m_d m_\Lambda} C^{1\ 1/2\ S}_{m_d\ m_\Lambda\ M_s}\chi^1_{m_d}\chi^{1/2}_{m_\Lambda} \quad (3)$$

shows the coupling of the spin-1 deuteron to the spin-$\frac{1}{2}$ Λ to give a total spin of $\frac{1}{2}$ or $\frac{3}{2}$. Thus, the sum over (l,S) in Eq. (1) includes $(0,\frac{1}{2})$ and $(2,\frac{3}{2})$ since the total spin of ${}^3_\Lambda$H is $\frac{1}{2}$. The angular momentum part is defined via

$$\mathcal{Y}^{lm}_{l_1l_2}(\hat{\mathbf{p}},\hat{\mathbf{q}})=\sum_{m_1m_2} C^{l_1\ l_2\ l}_{m_1\ m_2\ m} Y^{m_1}_{l_1}(\hat{\mathbf{p}})Y^{m_2}_{l_2}(\hat{\mathbf{q}})\ . \quad (4)$$

The isospin wave function [not shown in Eq. (1)] is defined by the $T=0$ nature of the deuteron. Finally, the Λ part of the wave function is given by

$$\varphi_\Lambda(q)=N\frac{\exp[-(q/Q_\Lambda)^2]}{q^2+\alpha^2}\ , \quad (5)$$

where $Q_\Lambda=1.17\ \mathrm{fm}^{-1}$, $\alpha=0.068\ \mathrm{fm}^{-1}$, and $N=0.322$.

III. MATRIX ELEMENTS IN MOMENTUM SPACE

In the pion-exchange description of the $\Lambda N\to NN$ amplitude the weak vertex Hamiltonian is given by

$$H_\pi=ig_w\overline{\Psi}_N(1+\lambda\gamma_5)\boldsymbol{\tau}\cdot\boldsymbol{\Phi}_\pi\Psi_\Lambda\ , \quad (6)$$

where λ and g_w are empirical constants which determine the strengths of the parity-conserving (PC) and parity-violating (PV) interactions and are adjusted to the free Λ decay. The nucleon and pion fields are given by Ψ_N and $\boldsymbol{\Phi}_\pi$, respectively, while $\Psi_\Lambda=\binom{0}{\Psi_\Lambda}$ is the Λ spurion field which is used to enforce the empirical $\Delta I=\frac{1}{2}$ sum rule. After performing a nonrelativistic reduction the $\Lambda N\to NN$ transition operator reads [14]

$$t(\mathbf{q}_\pi)=g_w\frac{g_{\pi NN}}{2M_N}\left[1+\frac{\lambda}{M_\Lambda+M_N}\boldsymbol{\sigma}_1\cdot\mathbf{q}_\pi\right]\times\frac{\mathcal{F}^2(\mathbf{q}^2_\pi)}{\mathbf{q}^2_\pi+m^2_\pi}\boldsymbol{\sigma}_2\cdot\mathbf{q}_\pi\boldsymbol{\tau}_1\cdot\boldsymbol{\tau}_2\ . \quad (7)$$

The form factor is denoted by $\mathcal{F}(\mathbf{q}^2_\pi)$ which depends on the momentum transfer $\mathbf{q}_\pi$. We choose a form factor of the type $\mathcal{F}(\mathbf{q}^2)=(\Lambda^2_\pi-m^2_\pi)/(\Lambda^2_\pi+\mathbf{q}^2)$ at each vertex. The cutoff mass is $\Lambda_\pi=1.3$ GeV, which has been shown to be appropriate in the study of the NN interaction. For our momentum-space calculation it is useful to write the amplitude in spherical notation (suppressing isospin)

$$\begin{aligned}t(\mathbf{q}_\pi)&=(L+\boldsymbol{\sigma}_1\cdot\mathbf{K})\frac{\mathcal{F}^2(\mathbf{q}^2_\pi)}{\mathbf{q}^2_\pi+m^2_\pi}\boldsymbol{\sigma}_2\cdot\mathbf{q}_\pi\\&=\frac{\mathcal{F}^2(\mathbf{q}^2_\pi)}{\mathbf{q}^2_\pi+m^2_\pi}\sum_{\substack{S_1M_{S_1}\\M_{S_2}}}(-1)^{M_{S_1}+M_{S_2}}(\sigma_1)^{S_1}_{M_{S_1}}K^{S_1}_{-M_{S_1}}\\&\quad\times(\sigma_2)^1_{M_{S_2}}(q_\pi)^1_{-M_{S_2}}\ ,\end{aligned} \quad (8)$$

where

$$K^0_0=L=g_w\frac{g_\pi NN}{2M_N}\ ,\quad \sigma^0_0=1\ ,$$
$$\mathbf{K}=g_w\frac{g_{\pi NN}}{2M_N}\frac{\lambda}{M_\Lambda+M_N}\mathbf{q}_\pi\ . \quad (9)$$

In terms of the Jacobi coordinates introduced in the preceding section the full matrix element in momentum space is given by

$$\langle pnn|t_{\Lambda N\to NN}|{}^3_\Lambda\mathrm{H}\rangle=\frac{1}{(2\pi)^3}\int d^3q\,d^3p'\,d^3q'\Psi^*_f(\mathbf{q}',\mathbf{p}')t(\mathbf{q}_\pi)\Psi_{M_i}(\mathbf{q},\mathbf{p})\ , \quad (10)$$

where the deuteron pair momentum in the hypertriton is fixed by momentum conservation, $\mathbf{p}=\mathbf{p}'+\mathbf{q}_\pi/2$, and $\mathbf{q}_\pi=\mathbf{q}'-\mathbf{q}$. A diagrammatic illustration of the process ${}^3_\Lambda\mathrm{H}\to pnn$ is shown in Fig. 1. The initial ${}^3_\Lambda$H wave function depends on its spin projection M_i while the final continuum wave function depends on the asymptotic momenta of the three nucleons, $\mathbf{k}_1$, $\mathbf{k}_2$, and $\mathbf{k}_3$ with $\mathbf{k}_1+\mathbf{k}_2+\mathbf{k}_3=0$ and their spin projections m_1, m_2, and m_3. Since the full three-body scattering state is very difficult to handle, we neglect final-state interactions by writing

$$\Psi_f(\mathbf{q}',\mathbf{p}')=(2\pi)^3\delta(\mathbf{q}_a-\mathbf{q}')\delta(\mathbf{p}_a-\mathbf{p}')\chi_{m_1}\chi_{m_2}\chi_{m_3}\ , \quad (11)$$

where $\mathbf{q}_a$ and $\mathbf{p}_a$ are the asymptotic spectator and pair momenta of the three nucleons in the final state. Since we are interested only in inclusive observables and, therefore, integrate over all possible final states, this approximation should be adequate. We thus obtain for the matrix elements of Eq. (10)

$$t_{fi}=\int d^3q\ t(\mathbf{q}_a-\mathbf{q})\Psi_{M_i}(\mathbf{q},\mathbf{p}_a+(\mathbf{q}_a-\mathbf{q})/2)\ . \quad (12)$$

Using the ${}^3_\Lambda$H wave function as given in Eq. (1) and the $\Lambda N\to NN$ operator from Eq. (8)

$$t_{fi}=2\sqrt{18}(-1)^{3/2-M_i}\sum_{\substack{(l,S)S_1\\ m_l,m_\Lambda}}(-1)^{S-M_S-m_\Lambda-m_N}\hat{S}\hat{S}_1\begin{pmatrix} l & S & \frac{1}{2}\\ m_l & M_S & -M_i\end{pmatrix}$$
$$\times\begin{pmatrix} \frac{1}{2} & 1 & S\\ m_\Lambda & m_d & -M_S\end{pmatrix}\begin{pmatrix} \frac{1}{2} & \frac{1}{2} & 1\\ m_N & m_3 & -m_d\end{pmatrix}\begin{pmatrix} \frac{1}{2} & S_1 & \frac{1}{2}\\ -m_1 & M_{S_1} & m_\Lambda\end{pmatrix}\begin{pmatrix} \frac{1}{2} & 1 & \frac{1}{2}\\ -m_2 & M_{S_2} & m_N\end{pmatrix} I^{l\;S_1}_{m\,M_{S_1}M_{S_2}} \qquad (13)$$

where $\hat{S}=\sqrt{(2S+1)}$. The integrals

$$I^{l\;S_1}_{m_l\,M_{S_1}M_{S_2}}=\int d^3q\,\varphi_\Lambda(q)\Psi^{l}_{\rm deut}(p)Y_l^{m_l}(\hat{\mathbf{p}})K^{S_1}_{-M_{S_1}}(q_\pi)^1_{-M_{S_2}} \qquad (14)$$

are evaluated numerically. Note that the values of M_S, m_d, m_N, M_{S_1}, and M_{S_2}, in Eq. (13) are fixed by the $3j$ symbols while the spin projections m_1, m_2, and m_3 along with the respective isospin projections are summed incoherently to obtain the total decay rate

$$\Gamma_{\rm nm}=\frac{1}{(2\pi)^5}\int d^3k_1 d^3k_2\,\delta(M-E_1-E_2-E_3)\times\frac{1}{2}\sum_{\substack{M_i m_1\\ m_2 m_3}}|t^{\mathcal{A}}_{fi}|^2 \qquad (15)$$

where M, E_1, E_2, and E_3 are the mass of ${}^3_\Lambda$H and the total asymptotic energies of the three final nucleons, respectively. In the matrix element $t^{\mathcal{A}}_{fi}$ antisymmetry in the final state has properly been included. We have averaged over M_i, the projection of ${}^3_\Lambda$H in the initial state. Exploiting the energy conserving δ function as well as choosing $\mathbf{k}_1$ to be along the z direction reduces the six integrals in Eq. (15) to two which are also performed numerically. The largest contribution of the matrix element comes from the kinematical situation $\mathbf{k}_1=-\mathbf{k}_2$ since this allows maximal overlap of the wave functions in the integrals of Eq. (14).

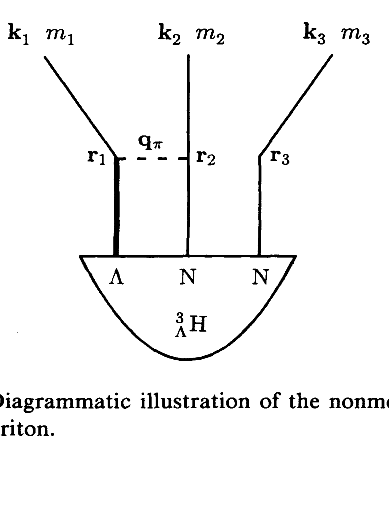

FIG. 1. Diagrammatic illustration of the nonmesonic decay of the hypertriton.

IV. RESULTS AND DISCUSSION

The results of our calculation for the total nonmesonic decay rate of ${}^3_\Lambda$H are presented in Table I in units of the free lambda decay, $\Gamma_\Lambda=3.8\times10^9\ {\rm s}^{-1}$. Without form factor we obtain a nonmesonic rate of $\Gamma_{\rm nm}=0.076$ for the hypertriton. Including the form factor with a cutoff mass [15] of $\Lambda_\pi=1300$ MeV reduces the decay rate by about 35% which is similar to the reduction found for the rate of ${}^{12}_{\Lambda}$C [7]. Our result, $\Gamma_{\rm nm}({}^3_\Lambda{\rm H})=0.049$, is substantially smaller than nonmesonic decay rates for heavier nuclei [7,14], such as ${}^{12}_{\Lambda}$C where $\Gamma_{\rm nm}({}^{12}_{\Lambda}{\rm C})=0.35$. This is to be expected since not only is the number of nucleons in ${}^3_\Lambda$H much smaller but the hypertriton is also a very dilute system due to the small Λ separation energy. Therefore the lambda wave function has fewer high momentum components which are required for a large momentum transfer process such as $\Lambda N\to NN$. This behavior is illustrated in Fig. 2 where $\Gamma_{\rm nm}({}^3_\Lambda{\rm H})$ is plotted against the Λ separation energy B_Λ. The total rate increases from 0.049 to 0.43 when B_Λ rises from 0.13 to 10 MeV. Changing B_Λ yields new parameters α and Q_Λ in the wave function $\phi_\Lambda(q)$ in Eq. (5); these quantities are connected through a simple transcendental equation. Our value of $\Gamma_{\rm nm}({}^3_\Lambda{\rm H})=0.049$ is, on the other hand, substantially larger than previous estimates in the early works of Dalitz and collaborators [8] who obtained total rates for the decay of ${}^3_\Lambda$H of less than 1%. However, no microscopic calculation was performed, instead Ref. [8] used a phenomenological model combined with some experimental information that had very large error bars [16].

Due to the approximations discussed in Sec. II the ${}^3_\Lambda$H wave function contains only two partial waves compared to 11 or more in the trinucleon system [17]. The two partial waves present in our hypertriton wave function are just the s- and d-wave part of the deuteron wave function since the Λ is restricted to be in a relative s state. The in-

TABLE I. Nonmesonic weak decay rates for ${}^3_\Lambda$H.

	$\Gamma_{\rm nm}$	Γ_n/Γ_p	$\Gamma_{\rm PV}/\Gamma_{\rm PC}$
S wave	0.072	0.70	0.44
S and D wave	0.076	0.68	0.45
S and D wave including form factor	0.049	0.68	0.46

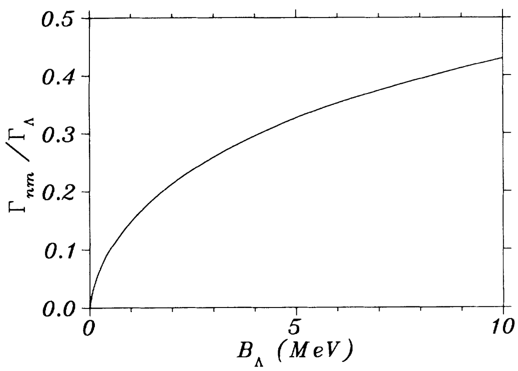

FIG. 2. The total decay rate $\Gamma_{\rm nm}$ as a function of the Λ separation energy.

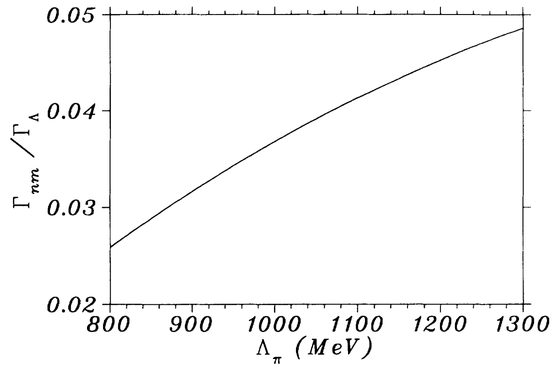

FIG. 3. The total decay rate $\Gamma_{\rm nm}$ as a function of the cutoff Λ_π in the form factor.

clusion of the d wave enhances the total rate by about 5%, as shown in Table I. Additional partial waves would include states where the p-n system is in a $L=1$, $S=0$, $T=0$ state and the Λ in a relative $L=1$ as well. Such contributions are known to be small [17] in the ^{3}He and ^{3}H wave functions, and should be even smaller in $^3_\Lambda$H.

Short-range correlations to which $\Gamma_{\rm nm}$ is known to be sensitive are in principle fully included in a correlated $^3_\Lambda$H three-body wave function. This is in contrast to heavier hypernuclei where shell-model wave functions are used and a correlation function has to be added by hand [7]. However, due to the approximations inherent in our $^3_\Lambda$H wave function ΛN short-range correlations are not properly included. Since Ref. [18] has shown that the effect of correlations on the total rate can be simulated by a softer form factor, we present in Fig. 3 the rate $\Gamma_{\rm nm}$ as a function of the cutoff mass Λ_π. A more sophisticated calculation with an exact hypertriton wave function should be able to remove this uncertainty.

Finally, we present in Table I the ratios of the neutron-to-proton–induced rates as well as the parity-violating to parity-conserving rates. The ratio $\Gamma_n/\Gamma_p=0.68$ is significantly larger than the results coming from calculations for heavier nuclei and nuclear matter [14]. Note that even though the final state consists of two neutrons and one proton for both neutron- and proton-induced reactions, the two processes should be distinguishable by their distinct kinematical patterns. As discussed above, the largest contribution to the decay rate corresponds to the two nucleons involved in the elementary process $\Lambda N \to NN$ going out back to back with a kinetic energy of about 80 MeV while the third "spectator" remains basically at rest. Thus, with appropriate kinematic cuts applied in an experiment it should be possible to identify the nucleons involved in the basic reaction.

Our result for the ratio $\Gamma_{\rm PV}/\Gamma_{\rm PC}=0.46$ is comparable to the ratios obtained in heavier nuclei but it should be regarded as preliminary since the effects of final-state interactions have been neglected in this study. Reference [7] found that for $^{12}_\Lambda$C the ratio $\Gamma_{\rm PV}/\Gamma_{\rm PC}$ changed by more than a factor of 2 when final-state distortions of the outgoing nucleons were included in the calculation. Since exact wave functions for three-nucleon scattering states are not yet available, one might consider including only the interaction between the nucleons involved in the reaction $\Lambda N \to NN$, treating the third particle as a spectator. In view of the kinematical argument given above this should be a reasonable approximation.

In conclusion, we have performed for the first time a microscopic calculation for the nonmesonic weak decay of the hypertriton. The obtained decay rate is 4.9% of that of the free Λ. More sophisticated calculations should employ better $^3_\Lambda$H wave functions and include final-state interactions for the two nucleons involved in the process $\Lambda N \to NN$. Furthermore, since the pion exchange model has been demonstrated to be insufficient to explain the experimental data for $^{12}_\Lambda$C total and partial rates heavier meson exchanges need to be considered. Since the three-nucleon problem can in principle be solved exactly, nuclear structure uncertainties can be minimized and, therefore, the decaying hypertriton provides an ideal laboratory to study weak interaction processes.

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