π^- capture on the deuteron as a dibaryon search

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We consider radiative capture of a π^- from an atomic orbital around deuterium into a bound $NN\pi$ system with two units of isospin. It is found that it is a feasible method to search for such hypothetical states.

I. INTRODUCTION

The pion-nucleon p-wave interaction in the isospin $\frac{3}{2}$ state is known to be strongly attractive. This fact has led to suggestions that the pion-nucleus system might have bound states.[1] To maximize the number of isospin $\frac{3}{2}$ interactions, neutron-rich π^- systems are preferred. Thus, one has been led to consider the system $nn\pi^-$ (or its charge conjugate).

If this hadronic collection is indeed bound it is a dibaryon with $I=2$ and hence cannot be formed as an s-channel resonance in the nucleon-nucleon system. There is a long history of the study of this system.[2] Recently an investigation[3] using relativistic three-body equations has shown that the existence of such a bound $NN\pi$ state is strongly linked to the interaction range of the pion-nucleon system and hence to the basic hadronic size. The importance of these results is that, while the observance of such a state would be spectacular in itself, the failure of such a state to exist gives bounds on the range of the pion-nucleon interaction. These investigations focused on a pair of neutrons with zero relative angular momentum, with a pion of unit angular momentum, to give a 1^+ state.

The latest results of these three-body calculations,[4] based on the two nucleons being in a relative p state, do not have such a simple interpretation since they involve two interaction ranges. They indicate, however, that the best representation of the p-wave scattering volume coincides with those situations in which the bound state exists. This latest study finds that the most promising case is for the nucleon pair to have relative angular momentum unity and for the pion also to have unit angular momentum with respect to the center of mass of the two nucleons, to give a total 2^- state. The resolution of this question is far from complete, and it is important to know if such bound systems are to be found in nature.

There exist (to our knowledge) no predictions of $I=2$ dibaryons in this energy region based on quark dynamics (although there do exist predictions[5] of particles in this range based on phenomenological extrapolations from observed structures.[6]) Hence we will assume, as suggested, that the particle has an internal composition of three hadrons (two nucleons and one pion) and that the size of the system is of the order of 1 fm.

One possible way to search for this bound system is with the reaction

$$\pi^- + {}^2\mathrm{H} \rightarrow \pi^+ + X^- \tag{1}$$

(or its charge conjugate). This reaction has the advantage that the product (X^-) must be an $I=2$ object, hence the background from the formation of other states is low. It also has the merit that, if the object exists, beams can be produced, lifetimes measured, etc. It has the disadvantage that a pion must be produced, and the calculation of pion production by pions in ordinary nuclear systems is difficult enough.

A recent measurement of this process[7] has placed an upper limit on the cross section for producing this state with this reaction of 9–15 nb/sr for the region 0–30 MeV binding energy. This same paper also reports the preliminary results of a second experiment which gives yet lower bounds by at least an order of magnitude. These experiments provide a significant constraint on the existence of such a state. While the results of this experiment are 2 orders of magnitude smaller than a recent prediction,[8] we note that the related cross sections for (π, 2π) reactions on nuclei are difficult to calculate and that the radiochemically measured[9] cross sections in this energy region are very small (around 10 μb at 350 MeV). Hence, it is desirable to have a method for searching for this bound state which does not depend on pion production.

To this end we first note that if the $I=2$ state exists, it will be found in all of the possible charge states. While the $T_z=\pm 2$ state can only decay weakly (because of charge conservation) the $T_z=0,\pm 1$ states can decay by the emission of a photon. The width of the state is still small enough to provide a well-defined particle. (See Sec. III).

In this paper we consider the X^0 as formed in the reaction[10]

$$\pi^- + {}^2\mathrm{H} \rightarrow X^0 \gamma \tag{2}$$

from an atomic orbital. The advantage of this transformation is that it can proceed by a purely electromagnetic transition with the only uncertainties in the rate arising from the lack of knowledge of the detailed structure of the X^0 itself. While the Stark mixing assures that the capture takes place from an atomic s state, it is not necessarily from the $1s$ state. We make estimates here of the rate based on capture from the $1s$ state, which is consistent with the usual practice of quoting absolute rates assuming capture from the lowest s state. We do not consider the possibility of capture from the $2p$ state, which might be an experiment to be performed in a low-density gaseous target.

Indeed, if the X particle is well described as being a bound state of two nucleons and a pion, then (2) might have some advantages over (1) since all the particles needed to produce the X^0 are already present in the initial state, and it is only necessary to rearrange the wave function. Reaction (2) can produce particles with $I=0$, 1, or 2, but we shall focus only on the lattermost possibility.

The questions we wish to address are (i) what is the transition rate for (2), and (ii) what is the decay width of the X^0. If the branching ratio (B.R.)

$$\text{B.R.} = \frac{\Gamma(\pi^- + {}^2\text{H} \rightarrow \gamma + X^0)}{\Gamma(\pi^- + {}^2\text{H} \rightarrow \text{anything})} \tag{3}$$

is too small, or if the width of the X^0 is too large, then it will be difficult to detect using (2).

To calculate either one of these quantities requires a knowledge of the three-body wave function of X^0. In this paper we shall use simple wave functions and make crude estimates of the matrix elements that are needed.

II. CAPTURE RATE

To describe the possible quantum numbers of X, we shall use the set of Jacobi coordinates shown in Fig. 1. The internal orbital angular momentum is decomposed into

$$l = l_\rho + l_\lambda , \tag{4}$$

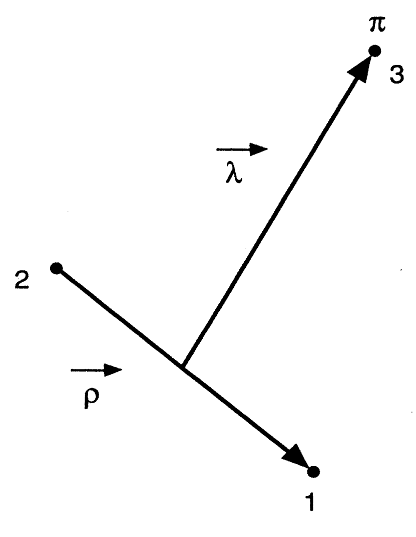

FIG. 1. Definition of the Jacobi coordinates used in the text.

where l_ρ is associated with the relative motion of the two nucleons and l_λ involve the motion of the pion with respect to the center of mass of the two nucleons. Since the isospin of the two nucleons must be unity in order that the total $I=2$, the allowed states of the two nucleons are ${}^1S_0$ with $l_\rho=0$, ${}^3P_J$ with $l_\rho=1$ $(J=0,1,2)$, etc. Considering only $l_\lambda=0$ or 1, the possible quantum numbers of the X are shown in Table I.

In the initial state the π^- is in a Bohr orbit $(l_\lambda=0)$ and the neutron-proton pair is in a ${}^3S_1$-${}^3D_1$ state, i.e., $l_\rho=0$ with a small admixture of $l_\rho=2$. The initial quantum members, therefore, are 1^-. To obtain an estimate of the transition rate, we assume that only one-body electromagnetic current operators need be considered. Since the isospin of the two nucleons must change from 0 to 1, only radiation from the nucleons makes a contribution to the matrix element.

The simple wave functions that we use factorize into three parts: nucleon spin $\times$ nucleon space $\times$ pion space. Consequently, there are only four independent matrix element that need be calculated corresponding to the four cases in Table I: $l_\lambda=0$ or 1 and $l_\rho=0$ or 1. Insofar as isospin is concerned, the X^0 is a 50-50 mixture of $\pi^- np$ and $\pi^0 NN$ (with the two nucleons having $I=1$). We note, finally, that since the $I=1$ nucleon-nucleon system does not have a bound state, the factorized spatial wave function cannot be correct asymptotically. Our hope is that it is the sizes of the spatial wave functions rather than any detail of their shapes that governs the calculation of the transition rate.

Since each one of the four possible transitions is a dipole, the general expression for the transition probability per unit time is

$$\Gamma = \frac{16\pi}{9} k^3 \sum_m |T_{1m}|^2 . \tag{5}$$

We now examine T for the four different cases.

A. $l_\lambda=0, l_\rho=0$ $(J^\pi=0^-)$

Neglecting the D state of the deuteron (which we shall do henceforth), it is only the spin (and isospin) of the two nucleons that changes. The transition is a magnetic dipole and $T_{1m}=M'_{1m}$, where

TABLE I. Quantum numbers J^π of an $I=2$ dibaryon for the lowest angular momentum states of the two nucleons and the pion. The 0^- and 2^- states having $l_\lambda=1$ are the most likely ones to exist as a bound $NN\pi$ state according to Ref. 4.

Pion angular momentum l_λ	Nucleon-nucleon state ${}^1S_0$	${}^3P_J$
0	0^-	0^+, 1^+, 2^+
1	1^+	0^-, 1^-, 2^-

$$M'_{1m}=\frac{1}{\sqrt{2}}\frac{e}{2m_N}2\mu_v\left[\frac{3}{4\pi}\right]^{1/2}\langle S'=0,m_s'=0|(-1)^m\sigma_{-m}|S=1,m_s\rangle$$
$$\times\langle I'=1,m_I'=0|\tau_z|I=0,m_I=0\rangle\int d^3\rho d^3\lambda\psi^*_{X^0}(\rho,\lambda)\psi_{\pi^-D}(\rho,\lambda)\ . \tag{6}$$

The significance of the various quantities is as follows. The factor $1/\sqrt{2}$ is present because only the $\pi^- np$ component of the X^0 wave function takes part in the transition. $e/2m_N$ is the nucleon magneton. $\mu_v=\frac{1}{2}(\mu_p-\mu_n)=2.35$ is the isovector magnetic moment in units of the nucleon magneton, and the factor 2 arises since the spin and isospin matrix elements for nucleon 2 are each the negative of those for nucleon 1; it is the latter that are displayed in Eq. (6) where σ_m is the spherical component of the spin and m_s is the initial z component of the deuteron's spin. The isospin matrix element in Eq. (6) has the value unity, and the sum over m that is called for in Eq. (5) together with the average over m_s yields

$$\frac{1}{3}\sum_{m_s}\sum_{m}|\langle S'=0,m_s'=0|\sigma_m|S=1,m_s\rangle|^2=1\ . \tag{7}$$

We take the spatial wave functions to factorize as follows:

$$\psi_{\pi^-D}(\rho,\lambda)=\psi_D(\rho)\psi_B(\lambda)\ , \tag{8}$$

where

$$\psi_B(\lambda)=N_Be^{-\kappa_B\lambda}Y_{00}(\theta_\lambda,\phi_\lambda) \tag{9}$$

is the wave function of the lowest s-state Bohr orbital, with

$$\kappa_B=2m_Nm_\pi\alpha/(2m_N+m_\pi)\simeq m_\pi\alpha$$

being the inverse of the Bohr radius and $N_B^2=4\kappa_B^3$. $\Psi_D(\rho)$ is the wave function of the deuteron.

We also assume a factorized form for the X^0 wave function

$$\psi_{X^0}(\rho,\lambda)=\psi_{NN}(\rho)\psi_\pi(\lambda)\ ,$$

and take $\psi_\pi(\lambda)$ to be a Yukawa with range κ_π^{-1},

$$\psi_\pi(\lambda)=N_\pi\frac{e^{-\kappa_\pi\lambda}}{\lambda}Y_{00}(\theta_\lambda,\phi_\lambda)\ , \tag{10}$$

with $N_\pi^2=2\kappa_\pi$. The spatial matrix element in Eq. (6) now factorizes into a nucleon part and a pion part,

$$\int d^3\rho\, d^3\lambda\psi^*_{X^0}(\rho,\lambda)\psi_{\pi^-D}(\rho,\lambda)=\langle I\rangle_{NN}\langle I\rangle_\pi\ , \tag{11}$$

with

$$\langle I\rangle_{NN}\equiv\int d^3\rho\psi^*_{NN}(\rho)\psi_D(\rho)\ , \tag{12}$$

and

$$\langle I\rangle_\pi\equiv\int d^3\lambda\psi^*_\pi(\lambda)\psi_B(\lambda)$$
$$=N_\pi N_B\int_0^\infty d\lambda\lambda e^{-(\kappa_\pi+\kappa_B)\lambda}$$
$$\simeq\sqrt{8}\left[\frac{\kappa_B}{\kappa_\pi}\right]^{3/2}\ , \tag{13}$$

where we have used the fact that $\kappa_B\ll\kappa_\pi$.

Since none of the two-body subsystems of the X^0 are bound, existence of a three-body bound state would require that there be an appreciable probability for finding all three particles simultaneously within the range of the forces. This leads us to expect that the size of the wave function in the NN separation variable is comparable with κ_π^{-1}, the size in the pion separation variable. This also means that the pion carries most of the kinetic energy, and leads to an estimate of $\kappa_\pi\simeq(2m_\pi B)^{1/2}$, where B is the binding energy of the X^0.

With B in the range of 10–40 MeV, κ_π^{-1} varies from ~ 4 to ~ 2 fm. Even if the size in the NN separation variable were only 1 fm, which would be surprising given the preceding discussion, the overlap with the deuteron wave function would be appreciable, $\gtrsim 0.5$, based on simple model calculations. This is the basis for the "standard" values of κ_π, $\langle I\rangle_{NN}$, $\langle\rho\rangle_{NN}$, and $\langle p_\rho\rangle$ that are used in the following numerical estimates; but the formulas are all written in a convenient manner for changing these values if it is desired to do so.

Putting all these results into Eq. (5) gives

$$\Gamma\simeq\frac{2}{3}k^3\frac{\alpha\mu_v^2}{m_N^2}\left[\frac{m_\pi\alpha}{\kappa_\pi}\right]^3|\langle I_{NN}\rangle|^2\ ,$$
$$\simeq\left[\frac{k}{35\text{ MeV}}\right]^3\left[\frac{0.5\text{ fm}^{-1}}{\kappa_\pi}\right]^3\left|\frac{\langle I_{NN}\rangle}{0.7}\right|^2$$
$$\times 8.7\times10^{12}/\text{sec}\ , \tag{14}$$

where the final form of Eq. (14) has been referred to the "standard" values $k=35$ MeV, $\kappa_\pi=0.5$ fm^{-1}, and $\langle I_{NN}\rangle=0.7$.

B. $l_\lambda=0, l_\rho=1$ ($J^\pi=0^+,1^+,2^+$)

Since the orbital angular momentum of the two nucleons changes from 0 to 1, but their spin state remains a triplet, an electric dipole transition occurs with $T_{1m}=Q_{1m}$, where

$$Q_{1m}=\frac{1}{\sqrt{2}}\frac{e}{2}\langle S'=1,m_s'|I|S=1,m_s\rangle \times\langle I'=1,m_I'=0|\tau_z|I=0,m_I=0\rangle \times\int d^3\rho\, d^3\lambda\, \psi^*_{X^0}(\rho,\lambda)\rho Y^*_{1m}(\theta_\rho,\phi_\rho)\psi_{\pi^- D}(\rho,\lambda)\ . \tag{15}$$

The other difference from Eq. (6) arises from the fact that the electric dipole operator for the two nucleons is $\frac{1}{2}e\tau_{1z}(\mathbf{r}_1\text{-}\mathbf{r}_2)$. Neglecting any spin-orbit interaction and carrying out the angular integration over the nucleon coordinates and the spin sums gives

$$\Gamma=\frac{16\pi}{9}k^3\frac{1}{2}\frac{\alpha}{4}\frac{1}{4\pi}\frac{2J+1}{3}|\langle I_\pi\rangle\cdot\langle\rho\rangle_{NN}|^2\ , \tag{16}$$

where I_π is given in Eq. (13) and

$$\langle\rho\rangle_{NN}=\int_0^\infty d\rho\,\rho^2 R^*_{NN}(\rho)\rho R_D(\rho) \tag{17}$$

is the radial part of the nucleon-nucleon matrix element. J is the spin of the X^0. The result is

$$\Gamma=\left[\frac{k}{35\text{ MeV}}\right]^3\left[\frac{0.5\text{ fm}^{-1}}{\kappa_\pi}\right]^3\left[\frac{\langle\rho\rangle_{NN}}{1\text{ fm}}\right]^2 \times(2J+1)\times 2.0\times 10^{12}/\text{sec}\ . \tag{18}$$

C. $l_\lambda=1,\ l_\rho=0\ (J^\pi=1^+)$

For this case, both the orbital angular momentum of the pion and the spin state of the nucleons change. Such a transition can be produced by the electric dipole operator Q'_{1m}, since $\mathbf{r}_1\times\boldsymbol{\sigma}_1+\mathbf{r}_2\times\boldsymbol{\sigma}_2$ involves the coordinate $\boldsymbol{\lambda}$, but weighted by the small factor m_π/m_N. Compared with Q_{1m} (case B above), Q'_{1m} is down by a factor of the order $\mu_v(k/m_N)(m_\pi/m_N)$, so that even a small admixture of the 1^+ wave function of type B would dominate the transition rate.

D. $l_\lambda=1,\ l_\rho=1\ (J^\pi=0^-,1^-,2^-)$

A magnetic dipole transition produced by the orbital angular momentum of the nucleons can change the orbital state of the pion as well. $T_{1m}=M_{1m}$, where

$$M_{1m}=\frac{1}{\sqrt{2}}\frac{1}{l+1}\frac{1}{2}\frac{e}{m_N}\left[\frac{3}{4\pi}\right]^{1/2}\langle S'=1,m_s'|I|S=1,m_s\rangle \times\int d^3\rho\, d^3\lambda\,\psi^*_{X^0}(\rho,\lambda)(-1)^m \times(L_1-L_2)_{-m}\psi_{\pi^- D}(\rho,\lambda)\ , \tag{19}$$

and the isospin matrix element is unity as before. Expressing the individual orbital angular momenta in terms of relative coordinates gives

$$\mathbf{L}_1-\mathbf{L}_2=-2\left[\frac{m_\pi}{2m_N+m_\pi}\boldsymbol{\lambda}\times\mathbf{p}_\rho+\tfrac{1}{4}\boldsymbol{\rho}\times\mathbf{p}_\lambda\right]+\mathbf{L}'\ , \tag{20}$$

where the terms in $\mathbf{L}'$ do not involve both $\boldsymbol{\lambda}$ and $\boldsymbol{\rho}$, and hence are not able to produce the transition. The spherical components of this operator are

$$(L_1-L_2)_{-m}\simeq 2\sqrt{2}i\begin{bmatrix}1 & 1 & 1\\ m_\lambda & m_\rho & -m\end{bmatrix}\left[\frac{m_\pi}{2m_N}\lambda p_\rho-\frac{1}{4}\rho p_\lambda\right]Y_{1m_\rho}(\theta_\rho,\phi_\rho)Y_{1m_\lambda}(\theta_\lambda,\phi_\lambda)+L'_{-m}\ . \tag{21}$$

The wave function of the X^0 is written

$$\psi^{J,M}_{X^0}(\rho,\lambda)=R_{NN}(\rho)R_\pi(\lambda)\sum_{m_\rho',m_\lambda'}\sum_{m_s'}\sum_{L',m_L'}\begin{bmatrix}1 & 1 & L'\\ m_\rho' & m_\lambda' & m_L'\end{bmatrix}\begin{bmatrix}L' & 1 & J\\ m_L' & m_s' & M\end{bmatrix}Y_{1m_\rho'}(\theta_\rho,\phi_\rho)Y_{1m_\lambda'}(\theta_\lambda,\phi_\lambda)X_1^{m_s'}\ , \tag{22}$$

and the radial function for the pion is taken to be

$$R_\pi(\lambda)=N'_\pi\frac{e^{-\kappa_\pi\lambda}}{\lambda}\left[1+\frac{1}{\kappa_\pi\lambda}\right]F(\lambda,R_0)\ , \tag{23}$$

where R_0 is the range of the interaction and F is a function chosen to give the wave function the correct behavior at the origin and approaches unity for $\lambda>R_0$. We expect that $N'^2_\pi\approx\kappa_\pi^2R_0$, with the extra factor of $\kappa_\pi R_0$ arising from the p-wave nature of the wave function.

Of the two terms in the bracket in Eq. (21), we expect the first to dominate in spite of the presence of the small factor m_π/m_N because the derivative with respect to λ in the second term brings down the very small factor κ_B. The expected magnitude of the radial part of the pion matrix element is

$$\langle\lambda\rangle\sim N_BN'_\pi\frac{1}{\kappa_\pi^3}\approx\beta N_B\sqrt{R_0}/\kappa_\pi^2\ . \tag{24}$$

We considered two forms for F:

$$F(\lambda,R_0)=1-e^{-(\lambda/R_0)^3}\ , \tag{25a}$$

$$F(\lambda,R_0)=(1-e^{-\lambda/R_0})^3\ , \tag{25b}$$

both having the property that they tend to unity for $\lambda\gg R_0$ and they are proportional to λ^3 for small λ. For a series of κ_π and R_0 in the relevant range, the values of β varied from 3.2 to 9.0. We use the smallest value from this set. The radial part of the nucleon-nucleon matrix element, $\langle p_\rho \rangle$, is expected to be of order 1 fm^{-1}.

Inserting Eqs. (21)–(23) into Eq. (19) and carrying out the spin and angular part of the matrix elements gives

$$M_{1m} \simeq \frac{1}{\sqrt{2}} \frac{1}{2} \frac{e}{2m_N} \left[\frac{3}{4\pi} \right]^{1/2} 2\sqrt{2}i \frac{m_\pi}{2m_N} N_B \frac{\beta\sqrt{R_0}}{\kappa_\pi^2} \times \langle p_\rho \rangle \frac{1}{3} \begin{bmatrix} 1 & 1 & J \\ -m & m_s & M \end{bmatrix} . \quad (26)$$

Squaring, summing over m and M, and averaging over m_s yields

$$\Gamma \simeq \frac{\beta^2}{648} k^3 \left[\frac{m_\pi \alpha}{m_N} \right]^4 \frac{m_\pi R_0}{\kappa_\pi^4} |\langle p_\rho \rangle|^2 (2J+1) \quad (27)$$

$$\simeq \left[\frac{k}{35\ \text{MeV}} \right]^3 \left[\frac{0.5\ \text{fm}^{-1}}{\kappa_\pi} \right]^4 \left[\frac{R_0}{2\ \text{fm}} \right] \times \left| \frac{\langle p_\rho \rangle}{1\ \text{fm}^{-1}} \right|^2 (2J+1) \times 6.6 \times 10^9/\text{sec} \quad (28)$$

for $J=0$, 1, or 2, which are the only states that can be reached by this dipole transition. The suppression of this rate compared with that produced by the M'_{1m} operator (see Sec. II A) arises primarily from the lever arm factor $(m_\pi/2m_N)^2$. For the $J^\pi=0^-$ case any small admixture of the 0^- wave function of type A would dominate the transition rate.

Comparison of the preceding four cases shows that the first two cases with $l_\lambda=0$, and $l_\rho=0$ or 1, would produce the largest transition rates, but still rather small because of the small energy release. If the binding energy of X^0 is less than 20 MeV, then it would appear to be unlikely that the rate for $\pi^- D \to \gamma X^0$ could be significantly larger than 10^{12}/sec. For the fourth case where $l_\lambda=1$ and $l_\rho=1$, which might be a preferred candidate for a bound X^0, the rate is down from this value by more than 3 orders of magnitude. For such a case consideration of exchange currents might increase the rate by a significant amount.

The rate for $\pi^- D \to \gamma nn$ is not known experimentally, but there is a theoretical calculation[11] that assigns the value

$$\Gamma(\pi^- D \to \gamma nn) = (4.3 \pm 0.5) \times 10^{14}/\text{sec} ,$$

using the value of the capture rate on the proton, which is in turn inferred from pion charge exchange and the Panofsky ratio.[12] We note that these are all "1s" rates. Using the measured ratio for other processes[13] one obtains a total rate of 2×10^{15}/sec.

In Fig. 2 we show some of the branching ratios as a function of photon energy. For this plot we have assumed "standard" values for all variables except k and κ_π. As previously described, the quantity κ_π was calculated by $\kappa_\pi = \sqrt{2m_\pi E_\gamma}$ in an attempt to correct for the change in size of the system due to the change in binding.

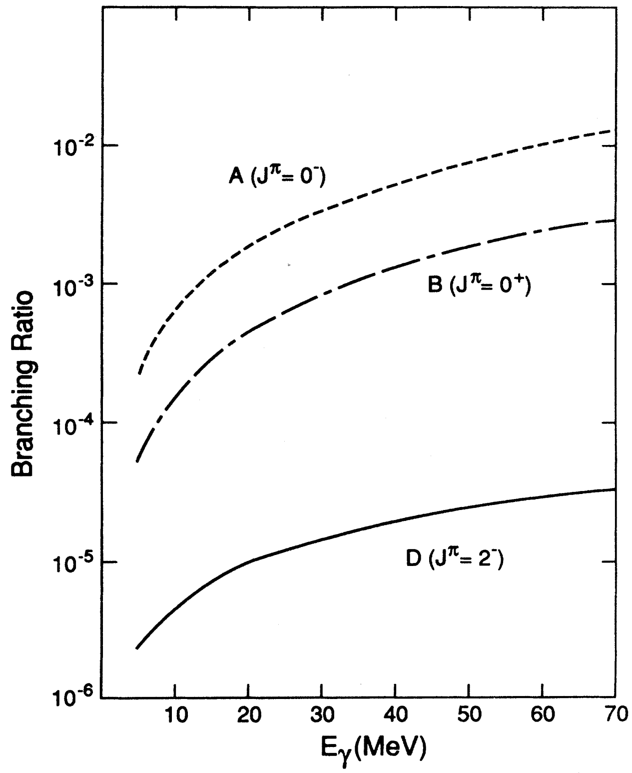

FIG. 2. The branching ratio from Eq. (3) for three of the cases discussed in the text. They are case A ($l_\lambda=0$, $l_\rho=0$), dotted curve; case B ($l_\lambda=0$, $l_\rho=1$), dotted-dashed curve; and case D ($l_\lambda=1$, $l_\rho=1$), solid curve. These values should only be considered as indicative since we have assumed the relationship $\kappa_\pi = \sqrt{2m_\pi E_\gamma}$ in their calculation. Only the least favorable case for B ($J=0$) and the most favorable for D ($J=2$) are shown. The other values scale as $2J+1$.

III. DECAY RATE

Another quantity of interest for a hypothetical X^0 is its decay rate. An electromagnetic interaction could take an $I=2$ particle to a γnn final state; and isospin mixing in the X^0 could lead to an NN final state. Since $X^0 \to \gamma nn$ goes via the same basic process, $\pi^- p \to \gamma n$, that also governs the rate for $\pi^- D \to \gamma nn$ (with the extra neutron in both cases acting as a spectator), we expect that the ratio of the two rates is primarily determined by the relative probabilities of finding the pion at the proton. For $\pi^- D$ capture, this is essentially the probability of finding the pion at the origin ($\lambda=0$). In the X^0 wave function, if the ranges in the ρ and λ variable are comparable, i.e., κ_π^{-1}, then we expect that

$$\frac{\Gamma(X^0 \to \gamma nn)}{\Gamma(\pi^- D \to \gamma nn)} \lesssim \left[\frac{\kappa_\pi}{\kappa_B} \right]^3 = \left[\frac{\kappa_\pi}{0.5\ \text{fm}^{-1}} \right]^3 9.0 \times 10^5 , \quad (29)$$

and together with the rate for $\pi^- D \to \gamma nn$ given above, the width of the X^0 would be expected to be less than 0.2 MeV.

We now want to show that Coulomb mixing of an $I=1$ component into the X^0 wave function followed by a nonradiative decay of this component to two neutrons is expected to contribute less to the width of the X^0 than the radiative decay of the dominant $I=2$ component. The matrix element for nn decay is given by

$$M=\sum_j \frac{\langle nn|D|I=1,j\rangle\langle I=1,j|V_c|X_0\rangle}{E_0-E_j}, \quad (30)$$

where V_c is the Coulomb potential and D is the strong interaction operator responsible for the decay. Even though the $NN\pi$ states with $I=1$ are all assumed to be in the continuum, a simple model leads to the expectation that

$$|M|\lesssim\frac{\langle V_c\rangle}{|E_0|}\langle D\rangle, \quad (31)$$

where the matrix elements $\langle D\rangle$ and $\langle V_c\rangle$ are to a discrete $I=1$ state. Since $\langle V_c\rangle\lesssim 1$ MeV and assuming the binding energy $|E_0|>10$ MeV, the effective amount of admixture of $I=1$ probability into the X^0 wave function is quite small, $P_{I=1}\lesssim 0.01$. Furthermore, one expects $\langle D\rangle$ to be such that

$$\frac{\Gamma(I=1\to nn)}{\Gamma(I=2\to nn\gamma)}\sim\frac{\Gamma(\pi^-D\to nn)}{\Gamma(\pi^-D\to nn\gamma)}\sim 3, \quad (32)$$

since the same basic mechanism is operating in X^0 decay as in π^-D capture. Therefore,

$$\frac{\Gamma(X^0\to nn)}{\Gamma(X^0\to nn\gamma)}=\frac{P_{I=1}\cdot\Gamma(I=1\to nn)}{\Gamma(I=2\to nn\gamma)}\ll 1, \quad (33)$$

and the $nn\gamma$ decay dominates the width of the X^0.

We note that a recent experiment[14] has been performed to search for the dibaryon based on the suggestion of Ref. 10. They[14] give an upper limit on the branching ratio from Eq. (3) that varies between $(1\text{–}6)\times 10^{-4}$ over the γ-ray energy interval 10–30 MeV. A comparison with Fig. 2 would appear to rule out the $J^\pi=0^-$ configuration in which all orbital angular momenta are zero (case A of Fig. 2), if the photon energy (which is approximately equal to the binding energy) is between 15 and 30 MeV.

We also note a recent variational calculation of the $I=2$ πNN system that does not find a bound state.[15]

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