PHYSICAL REVIEW VOLUME 160, NUMBER 4 20 AUGUST 1967

$C^{14}(p,n)N^{14}$ Reaction and the Two-Body Force*

C. WONG, J. D. ANDERSON, J. MCCLURE, AND B. POHL
Lawrence Radiation Laboratory, University of California, Livermore, California

AND

V. A. MADSEN AND F. SCHMITTROTH
Oregon State University, Corvallis, Oregon

(Received 31 March 1967)

The $C^{14}(p,n)N^{14}$ reaction to the ground and first two levels of N^{14} has been investigated for proton energies between 6 and 14 MeV. The angular distributions at the highest energies have been analyzed by using a finite range distorted-wave Born-approximation formalism, assuming the two-body force to be of the form $\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i(b+a\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i)f(r_{0i})$, where $\boldsymbol{\tau}_0$ and $\boldsymbol{\sigma}_0$ refer to the isospin and spin of the incident nucleon, $\boldsymbol{\tau}_i$ and $\boldsymbol{\sigma}_i$ refer to isospin and spin of the extra-core target nucleons, and $f(r_{0i})$ is the Yukawa form factor with a range of 1.4 F. Calculations employing Visscher-Ferrell wave functions for C^{14} and N^{14} gave a good fit for the $0^+\rightarrow 0^+$ transition and a reasonable value for b of 9 MeV. However, the calculations for the $0^+\rightarrow 1^+$ ground-state and second-excited-state transitions yielded inconsistent values for a of 7 MeV for the excited transition and 21 MeV for the ground-state transition. This latter result is not unexpected, since $(\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i)(\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i)$ is related to the Gamow-Teller β-decay operator. The fact that experimentally the (p,n) reaction to the ground state is not inhibited (while the corresponding β decay is strongly suppressed) indicates that the spin and charge-exchange operator is not sufficient to account for the (p,n) ground-state transition. An explanation of the observed ground-state cross section in $C^{14}(p,n)$ requires additional spin-flip mechanisms such as a tensor interaction or particle exchange.

I. INTRODUCTION

THIS paper is the latest in a continuing series on (p,n) reactions and the two-body force deduced therefrom. Previous publications[1–4] dealt, respectively, with mirror transitions in light nuclei [i.e., $N^{15}(p,n)$, $C^{13}(p,n)$; $Be^9(p,n)$, $B^{11}(p,n)$]; and (p,n) reactions in s and d shell nuclei, [i.e., $Al^{27}(p,n)$, $O^{18}(p,n)$, $O^{17}(p,n)$; $Mg^{25}(p,n)$, $Mg^{26}(p,n)$]. The observation of $0^+\rightarrow 1^+$ transitions[5] in $O^{18}(p,n)$ has been interpreted as evidence for a $(\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i)(\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i)$ component in the effective two-body force used in reaction theory. The simultaneous measurement of the $0^+\rightarrow 0^+$ transition in $O^{18}(p,n)$ has yielded an estimate of the ratio of the pure-charge exchange $(\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i)$ to charge plus spin-exchange $(\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i)\times(\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i)$ strength for the effective two-body force.[3,5] The isotope O^{18} is a particularly unfavorable case for the observation of the $0^+\rightarrow 0^+(p,n)$ transition since the isobaric transition cannot be resolved experimentally from the transitions to three close-lying levels[6] separated by 50–100 keV from each other. In this respect, $C^{14}(p,n)N^{14}$ is particularly favorable since the 0^+ first-excited level in N^{14} at 2.3 MeV is well separated from the ground state and the next higher 1^+ level at 3.95 MeV.[6] In addition, it was recognized that the $(\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i)(\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i)$ operator is related to the Gamow-Teller β-decay operator. If the $C^{14}(p,n)0^+\rightarrow 1^+$ ground-state transition is indeed induced by this operator, then the (p,n) cross section to the ground state should be suppressed simply because the corresponding C^{14} β decay is inhibited.[7] Experimentally, the ground-state (p,n) transition is not strongly suppressed; its cross section is approximately one-half of that to the higher 1^+ level. The immediate conclusion is that additional spin-flip mechanisms, such as particle exchange and/or a tensor component in the charge-exchange part of the effective two-body force, are needed to account for the ground-state transition.

II. EXPERIMENTAL METHOD

The 6–14-MeV protons were accelerated by the Livermore variable-energy cyclotron. The experimental geometry and time-of-flight electronics have been fully described in a previous paper.[4] The C^{14} (isotopic purity 91.4%) was obtained in the form of $C^{14}O_2$. The $C^{14}O_2$ was contained in a 1-in.-long by 1-in.-diam gas cell with $\frac{1}{4}$-mil tantalum entrance and exit windows at a pressure slightly below atmospheric. The energy spread introduced by energy loss in the $C^{14}O_2$ gas was typically ±100 keV for 10-MeV protons.

III. EXPERIMENTAL RESULTS

Figure 1 shows the time-of-flight spectrum for $C^{14}(p,n)$ at a bombarding energy of 10.4 MeV and laboratory observation angle of 45°. Two target γ rays are present since double display was employed, one time-to-height converter stop pulse for every two beam pulses. The peaks between the two γ rays are neutron groups leaving N^{14} in its ground state (n_0) and various excited states (excitation energy in parentheses). The

* Work performed under the auspices of the U. S. Atomic Energy Commission.

[1] C. Wong, J. D. Anderson, S. D. Bloom, J. W. McClure, and B. D. Walker, Phys. Rev. **123**, 598 (1961).
[2] B. D. Walker, C. Wong, J. D. Anderson, and J. W. McClure, Phys. Rev. **137**, B1504 (1965).
[3] J. D. Anderson, S. D. Bloom, W. F. Hornyak, V. A. Madsen, and C. Wong (to be published).
[4] C. Wong, J. D. Anderson, J. W. McClure, and B. A. Pohl, Phys. Rev. **156**, 1266 (1967).
[5] S. D. Bloom, J. D. Anderson, W. F. Hornyak, and C. Wong, Phys. Rev. Letters **15**, 264 (1965).
[6] F. Ajzenberg-Selove and T. Lauritsen, Nucl. Phys. **11**, 1 (1959).
[7] W. M. Visscher and R. A. Ferrell, Phys. Rev. **107**, 781 (1957).

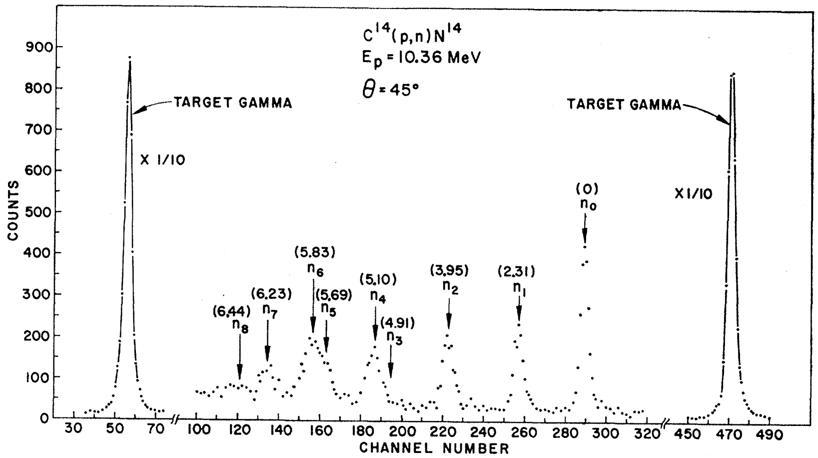

FIG. 1. Time-of-flight spectrum from $C^{14}(p,n)$ for a laboratory angle of 45° and bombarding energy of 10.36 MeV. Increasing time of flight is toward the left. The arrows are the calculated positions of the neutron groups leading to the ground (n_0) and various excited states in N^{14}. The excitation energies are given in parentheses.

arrows are the calculated positions of the various groups; the agreement between observation and calculation is excellent. With the exception of group $n_3(0^+ \rightarrow 0^-$ transition), all levels up to n_8 are excited in the (p,n) reaction. The spectrum of Fig. 1 was measured with a detector bias of 1.6-MeV neutron energy, and theoretically levels up to approximately 8.2 MeV in N^{14} could be observed. However, the region between 6.44 and 8.2 MeV excitation energy (below channel 100 in Fig. 1) was not plotted, since it is obscured by the tail of the intense target γ ray. Angular distributions for n_0, n_1, and n_2 as a function of bombarding energy are displayed in Fig. 2. The angular distributions have been fitted by a Legendre polynomial expansion. The zero-order coefficient in this expansion yields the integrated cross sections shown in Fig. 3. The rapid rise below 8 MeV for all three excitation functions is attributed to the presence of intermediate structure resonances. Since the direct reaction theory[8] has not been generalized to include the effects of such resonances, the angular distributions have

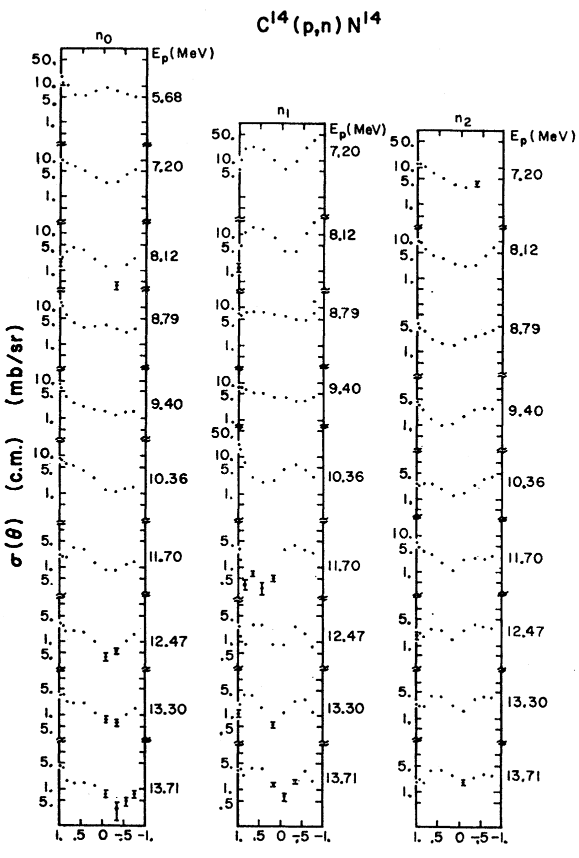

FIG. 2. Angular distributions as a function of bombarding energy for $C^{14}(p,n_0)$, $C^{14}(p,n_1)$, and $C^{14}(p,n_2)$. The n_1 is the neutron group leading to the 0^+ first-excited level at 2.31 MeV in N^{14}; n_0 and n_2 are the neutron groups leading to the 1^+ ground state and 1^+ second excited state at 3.95 MeV in N^{14}, respectively.

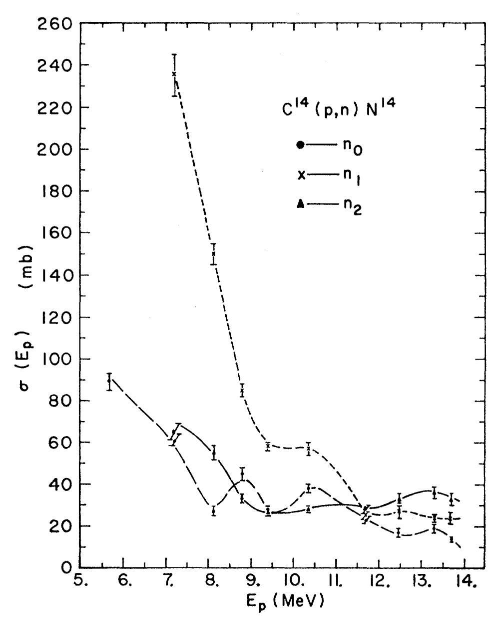

FIG. 3. Integrated cross sections as a function of bombarding energy for $C^{14}(p,n_0)$, $C^{14}(p,n_1)$, and $C^{14}(p,n_2)$.

[8] V. A. Madsen, Nucl. Phys. 80, 177 (1966).

been analyzed only at the highest bombarding energies where the excitation functions are slowly varying (see Fig. 3).

IV. THEORY

Recently, several treatments[8–10] of direct inelastic scattering have appeared. Here we summarize the main results of Ref. 8 applied to the (p,n) reaction. The differential cross section is given by the expression

$$\frac{d\sigma}{d\Omega}=\left(\frac{2m}{4\pi\hbar^2}\right)^2\frac{k_f}{k_i}\frac{1}{2(2J_i+1)}\times\sum_{II'LM}(2I+1)(2I'+1)\Big|\sum_{j_1j_2}D_{j_1j_2}(I,I',L)\times F_{LM}^{j_1j_2}(2L+1)^{-1/2}\Big|^2,\quad(1)$$

where[11]

$$F_{LM}^{j_1j_2\alpha_1\alpha_2}(\hat{k}_f)\equiv\langle\psi_{k_f}^{(-)}\psi(\mathbf{R})|Y_L^M(\hat{R})g_L^{j_1j_2\alpha_1\alpha_2}(R)|\psi_{k_i}^{(+)}\psi(\mathbf{R})\rangle,$$

$$g_L^{j_1j_2\alpha_1\alpha_2}(R)=\int R_{j_2l_2\alpha_2}(r_1)v_L(R,r_1)R_{j_1l_1\alpha_2}(r_1)r_1^2dr_1,$$

and

$$D_{j_1j_2}(II'L)=4(2j_1+1)^{1/2}(2j_2+1)^{1/2}\langle l_2\|Y_L\|l_1\rangle\times\begin{Bmatrix}j_1&\frac{1}{2}&l_1\\j_2&\frac{1}{2}&l_2\\I&I'&L\end{Bmatrix}S(J_iJ_fI;\,T_iT_f1;\,j_1j_2)\times C(T_iT_f1;P_i\,-P_f1)[\delta_{I'1}a+\delta_{I'0}b].\quad(2)$$

In Eq. (1), $F_{LM}^{j_1j_2}$ is the single-particle transition amplitude[8,11] from orbit j_1 to j_2, and $v_L(R,r_1)$ is the radial function appearing in the spherical harmonic expansion of the space part of the projectile-nucleon interaction. The terms a and b in Eq. (2) are the strengths of the effective nucleon-nucleon charge-exchange interaction

$$V(0,i)=V(r_{0i})\boldsymbol{\tau}_0\cdot\boldsymbol{\tau}_i[a\boldsymbol{\sigma}_0\cdot\boldsymbol{\sigma}_i+b],\quad(3)$$

and S is the spectroscopic amplitude,

$$S(J_iJ_fI;\,T_iT_f\tau;\,j_1j_2)=\frac{\langle\Phi_{J_fT_f}\|A_{I,\tau}\|\Phi_{J_iT_i}\rangle}{(2I+1)^{1/2}(2\tau+1)^{1/2}}.\quad(4)$$

In Eqs. (2) and (4), Φ_{JT} is a target nuclear state of total angular momentum J and isospin T. $A_{IN,\tau\rho}$ is the single-particle transition operator[8] for definite angular-momentum transfer I with projection N and isospin transfer τ with projection ρ:

$$A_{IN,\tau\rho}(j_1j_2)=\sum_{m_1m_2\alpha_1\alpha_2}C(j_1j_2I;\,m_1\,-m_2\,-N)\times C(\tfrac{1}{2}\tfrac{1}{2}\tau;\,\alpha_1\,-\alpha_2\,-\rho)\times(-1)^{j_1-m_1+\frac{1}{2}-\alpha_1}a_{j_2m_2\alpha_2}{}^{\dagger}a_{j_1m_1\alpha_1}.\quad(4')$$

[9] G. R. Satchler, Nucl. Phys. **77**, 481 (1966).
[10] N. K. Glendenning and M. Veneroni, Phys. Letters **14**, 228 (1965); Phys. Rev. **144**, 839 (1966).
[11] W. Tobocman, *Theory of Direct Nuclear Reactions* (Oxford University Press, London, 1961).

The spin-dependent part of the interaction Eq. (3) contributes only to terms in Eq. (1) for which the spin transfer $I'=1$, while the spin-independent term contributes terms for which $I'=0$. The 9-j coefficient in Eq. (2) along with Eq. (4) imply the following restrictions on the various quantum numbers:

$$|J_f-J_i|\leq I\leq J_i+J_f,\quad(5)$$

$$|j_2-j_1|\leq I\leq j_1+j_2,\quad(6)$$

$$0\leq I'\leq 1,\quad(7)$$

$$|l_2-l_1|\leq L\leq l_1+l_2,\quad(8)$$

$$|L-I'|\leq I\leq L+I',\quad(9)$$

$$|T_f-T_i|\leq 1\leq T_i+T_f,\quad(10)$$

while the reduced matrix element in Eq. (2) and conservation of parity in the nuclear states give the parity restriction

$$\Pi_f\Pi_i=(-1)^{l_1+l_2}=(-1)^L.\quad(11)$$

The cross section, Eq. (1), consists of an incoherent sum of contributions from the angular-momentum transfers I', L, and I, each of which consists of a coherent sum of angle-dependent, single-particle amplitudes weighted with the coefficients $D_{j_1j_2}$. These coefficients contain information about the single-particle coupling and about the details of the nuclear wave function in the spectroscopic-amplitude factor.

For $0^+\rightarrow0^+$ isobaric analog transitions, only the spin-independent term in the interaction Eq. (3) can contribute. Furthermore, only one term need be considered since the orbital angular-momentum transfer can only be $L=0$. For this case $I=I'=L=0$, which, according to Eq. (2), implies that there are only transitions for $j_1=j_2$. The transition operator Eq. (4′) for charge-exchange elastic scattering is simply

$$A_{00,10}(jj)=(2j+2)^{-1/2}[N_j(n)-N_j(p)],\quad(12)$$

where $N_j(n)$ is the number operator for j-shell nucleons. The spectroscopic amplitude Eq. (4) reduces to

$$S(JJ0;\,TT1;\,jj)=[6(2j+1)]^{-1/2}\times[(2J+1)(2T+1)]^{1/2}\times[C(T1T;\,T0T)]^{-1}\langle N_j-Z_j\rangle,\quad(13)$$

where J and T are the target spin and isospin and the last factor is the expectation value of j-shell neutron and proton numbers. The coefficient $D_{jj}(000)$ can be calculated in a straightforward way to give

$$D_{j_1j_2}(000)=\frac{b}{\sqrt{\pi}}\frac{[2(2J+1)]^{1/2}}{N-Z}\langle N_j-Z_j\rangle,\quad(14)$$

$$\frac{d\sigma}{d\Omega}=\left(\frac{2m}{4\pi\hbar^2}\right)^2\frac{k_f}{k_i}\frac{b^2}{\pi}\frac{1}{N-Z}\times\Big|\sum_j\langle N_j-Z_j\rangle F_{00}^{jj}\Big|^2.\quad(15)$$

It is usually true that for the important j shells in terms of their contribution to the reaction, the single-particle amplitude is roughly independent of j. Let us then take $F_{00}{}^{jj} \approx F_{00}$, where F_{00} is an average amplitude. The j sum in Eq. (15) can be done to give

$$\frac{d\sigma}{d\Omega}=\left(\frac{2m}{4\pi\hbar^2}\right)^2\frac{k_f}{k_i}\frac{b^2}{\pi}(N-Z)|F_{00}|^2. \tag{16}$$

The approximation made in obtaining Eq. (16) is particularly good for the $1p_{3/2}$ and $1p_{1/2}$ shell. Equation (16) also applies to the spin-independent $L=0$ term for analog transitions between states with nonzero J.

The important point here is that for $0^+ \rightarrow 0^+$ analog transitions the cross section is nearly independent of the phases and extent of configuration mixing. Thus $0^+ \rightarrow 0^+$ transitions are particularly favorable for testing the reaction model because of lack of sensitivity to the wave function.

The similarity between β decay, γ decay, charge exchange, and inelastic scattering has long been recognized.[7] The relationship between β decay and the (p,n) reaction is given below.

The essential point in the comparison is that the (p,n) reaction treated in distorted-wave Born approximation (DWBA) and β decay are both first-order processes involving one-body operators in nuclear coordinates. In beta decay the ft value is

$$ft=\frac{2\pi^3\hbar^7\ln 2}{m^5c^4\langle|M|^2\rangle}, \tag{17}$$

where $\langle|M|^2\rangle$ is the squared nuclear matrix element averaged over initial and final nuclear, nuclear-spin projections. All nuclear information is contained in this matrix element. The calculation for allowed β decay is essentially the same as that for charge exchange for $L=0$, Fermi decay corresponding to pure charge exchange and Gamow-Teller to charge exchange with spin exchange. A single formula can be given for both processes; the transition *rate* is

$$R=\frac{1}{2J_i+1}C(T_iT_f1;\,P_i-P_f1)^2\sum_I C_I(2I+1) \times\Big|\sum_{j_1j_2}\delta_{l_1l_2}[(2j_1+1)(2j_2+1)]^{1/2}W(I\tfrac{1}{2}j_2l_2;\,\tfrac{1}{2}j_1) \times S(J_iJ_fI;\,T_iT_f1;\,j_1j_2)H^{j_1j_2}\Big|^2. \tag{18}$$

For β decay, R is the averaged-squared nuclear matrix element $(\langle|M|^2\rangle)$ of Eq. (17); $I=0$ for Fermi decay and $I=1$ for Gamow-Teller decay; C_I is the square of the coupling constant, and

$$H^{j_1j_2}=\int R_{j_2}(r)R_{j_1}(r)r^2dr. \tag{19}$$

For the (p,n) charge-exchange reaction, R is the differential cross section;

$$C_I=\left(\frac{2m}{4\pi\hbar^2}\right)^2\frac{k_f}{k_i}\left(\frac{2}{3\pi}\right)[\delta_{I0}a^2+\delta_{I1}b^2], \tag{20}$$

and $H^{j_1j_2}=F_{00}{}^{j_1j_2}$ is the single-particle transition amplitude appearing in Eq. (1). To the extent that $H^{j_1j_2}$ is independent of j_1 and j_2, the β decay and (p,n) rates are proportional when only one value of I (either 1 or 0) is allowed in the (p,n) reaction by the triangle inequalities satisfied by the spectroscopic amplitude Eq. (4). This formal connection between β decay and the (p,n) process also applies to the part of the contribution due to space exchange and tensor interactions which corresponds to transfer of zero units of orbital angular momentum and one unit of spin to the nucleus.

V. DWBA DIRECT-REACTION CALCULATIONS

In Eq. (1), $\psi_{k_f}{}^{(-)}$ and $\psi_{k_i}{}^{(+)}$ are, respectively, the distorted waves for the emitted neutron and incident proton. The optical-model parameters for generating the distorted waves were assumed to be the same for the incident proton and emitted neutron. Since the Q value of the (p,n) reaction is, in general, nonzero, this is equivalent to assuming, in addition to charge independence, that the optical-model parameters are energy independent. These assumptions are permissible since the sensitivity of the calculations to changes of optical parameter is not very great.

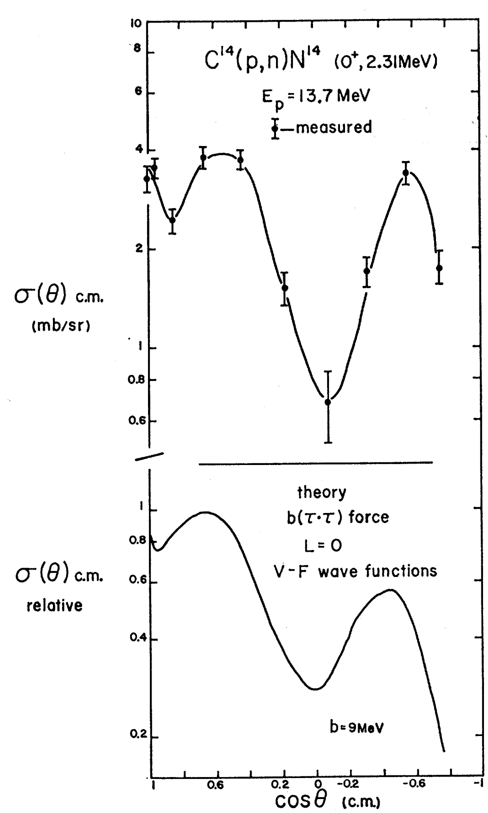

FIG. 4. Comparison of measurements and calculations for the $0^+ \rightarrow 0^+$ isobaric transition in $C^{14}(p,n)$. The proton bombarding energy was 13.7 MeV.

The parameters used were

$V_0 = 46$ MeV (real)	Woods-Saxon,
$W = 7.9$ MeV (imaginary)	Derivative Woods-Saxon,
$R = 3.01$ F	Nuclear radius,
$a_0 = 0.65$ F (real)	Diffuseness parameter,
$a_s = 0.47$ F (imaginary)	Diffuseness parameter.

The bound-state wave functions were generated with a Woods-Saxon real potential. The diffuseness (a_0) and nuclear radius (R) parameters are identical to those used in the scattering problem. However, the depths of the real potential were arbitrarily adjusted to yield separation energies of 9 MeV for the $p_{1/2}$ and 15.3 MeV for the $p_{3/2}$ bound states in C^{14}. The spin-orbit splitting of 6.3 MeV was taken from the $p_{1/2}-p_{3/2}$ splitting observed in N^{15} and O^{15}, while the 9-MeV separation energy is an average of the neutron separation energies in C^{14} and N^{14}. The $p_{3/2}$ and $p_{1/2}$ wave functions thus generated are therefore different. On the other hand, charge independence is assumed: the bound-state neutron and proton wave functions are identical. As a consistency check, the spin-orbit coupling strength needed to yield the observed splitting and the depth of the real potential were calculated. The spin-orbit coupling strength of 10 MeV appears reasonable; the real potential depth of approximately 51 MeV is about 10% larger than that used in the scattering problem.

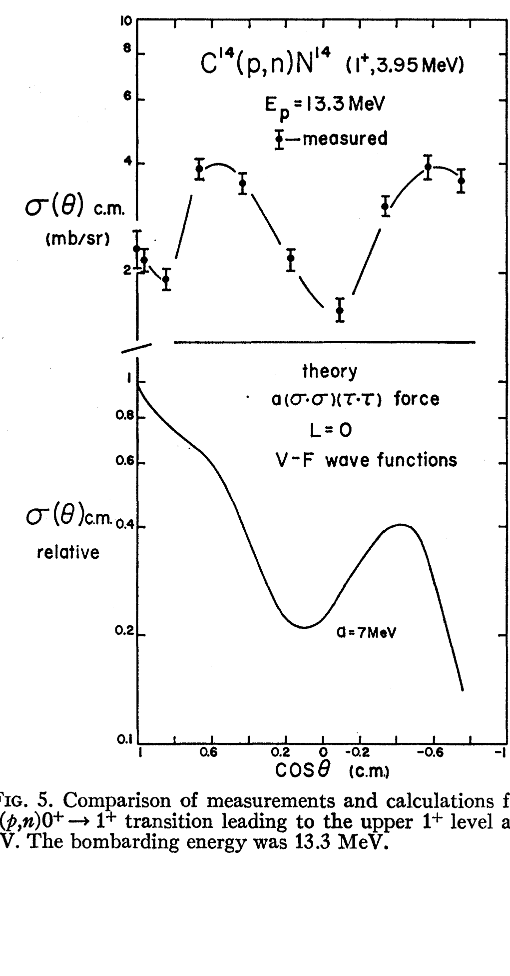

FIG. 5. Comparison of measurements and calculations for the $C^{14}(p,n)0^+ \rightarrow 1^+$ transition leading to the upper 1^+ level at 3.95 MeV. The bombarding energy was 13.3 MeV.

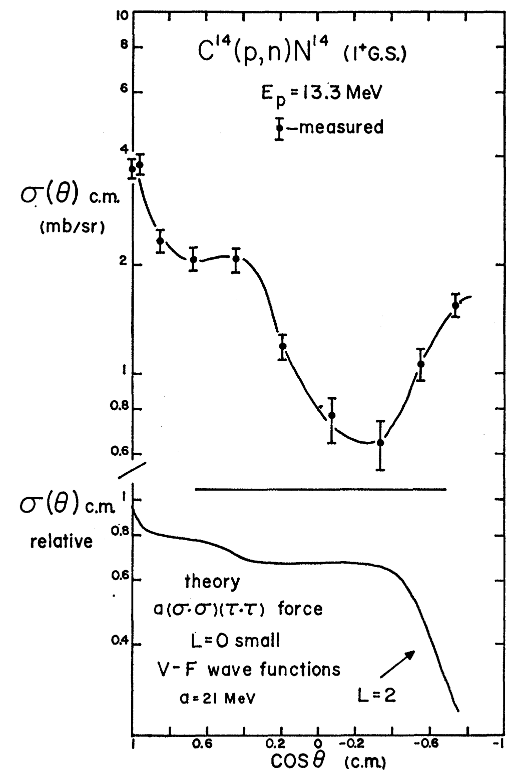

FIG. 6. Comparison of measurements and calculations for the $C^{14}(p,n)0^+ \rightarrow 1^+$ ground-state transition. The bombarding energy was 13.3 MeV.

For the (p,n) calculations, the Visscher-Ferrell wave functions[7] ($l-s$ representation) were transformed to the j-j representation to conveniently incorporate spin-orbit effects in the radial integrals.[12] The results are

1^+ N^{14} ground state:

$$0.926(p_{1/2})^2+0.362p_{1/2}p_{3/2}+0.119(p_{3/2})^2,$$

0^+ C^{14} ground state and

N^{14} first excited state:

$$0.968(p_{1/2})^2+0.250(p_{3/2})^2,$$

1^+ N^{14} second excited state:

$$-0.367(p_{1/2})^2+0.931(p_{1/2}p_{3/2})+0.0426(p_{3/2})^2.$$

A. $0^+ \rightarrow 0^+$ Transition

This transition is assumed to be induced by $b\,(\tau_0\cdot\tau_i)$, the pure charge-exchange operator. With Visscher-Ferrell wave functions used, the relevant $D_{j_1j_2}(I,I',L)$'s are $D_{1/2,1/2}(0,0,0)$ and $D_{3/2,3/2}(0,0,0)$. The calculated shape of the angular distribution for $E_p = 13.7$ MeV is shown at the bottom of Fig. 4. Normalizing the calcula-

[12] This procedure is not entirely correct. The introduction of spin-orbit differences in the radial wave functions, with no other changes in the Visscher-Ferrell wave functions, will increase the β-decay rate. However, we have calculated the radial integrals for β decay and find the spin-orbit splitting to have an effect on β decay which is very small compared to the inherent difference in radial integrals between β decay and the (p,n) reaction.

tions to the measured total integrated cross section yields a b value of 9 MeV. For $0^+ \to 0^+$ transitions, only $L=0$ is allowed. A comparison with the measured distribution (top of Fig. 4) shows that the main features of the distribution are well reproduced by the theory. As shown in Sec. IV, the prediction for the isobaric transition is very nearly configuration-independent, i.e., the assumption of a pure $(p_{1/2})^2{}_0$ configuration for C^{14} and N^{14} would yield nearly identical results.

B. $(0^+ \to 1^+)$ 3.95-MeV Transition

This transition is assumed to be induced by $a(\boldsymbol{\tau}_0 \cdot \boldsymbol{\tau}_i) \times (\boldsymbol{\sigma}_0 \cdot \boldsymbol{\sigma}_i)$. With Visscher-Ferrell wave functions, the relevant $D_{j_1j_2}$'s are $D_{1/2,1/2}(1,1,0)$, $D_{1/2,1/2}(1,1,2)$, $D_{1/2,3/2}(1,1,0)$, $D_{1/2,3/2}(1,1,2)$, $D_{3/2,1/2}(1,1,0)$, $D_{3/2,1/2}(1,1,2)$, $D_{3/2,3/2}(1,1,0)$, and $D_{3/2,3/2}(1,1,2)$. Only $L=0$ is shown in Fig. 5 since the $L=2$ contribution is smaller by a factor of 300. Normalizing the calculations to the total integrated cross section yields an a value of 7 MeV. With the exception of the forward angles, where through distortion effects the sensitivity to optical parameters is greatest, the trend of the measured angular distribution is roughly reproduced by the theory. A comparison with single-configuration calculations [assuming C^{14} is $(p_{1/2})^2{}_0$ and N^{14} (3.95 MeV) is $(p_{1/2}p_{3/2})_1$] shows that the use of Visscher-Ferrell wave functions suppresses the $L=2$ contribution relative to $L=0$ by a factor of 10.

C. $(0^+ \to 1^+)$ Ground-State Transition

The spin- and charge-exchange operator and relevant $D_{j_1j_2}$'s are the same as in Sec. V B. Calculations with Visscher-Ferrell wave functions show that, indeed as in β decay, the $L=0$ contribution is suppressed because of a cancellation of matrix elements.[7] Calculations show that the $F_{L,M}{}^{j_1j_2}$'s in Eq. (1) are nearly equal and $\sum_{j_1j_2} D_{j_1j_2} \approx 0$ for $L=0$. Compared to single-configuration calculations [i.e., C^{14} is $(p_{1/2})^2{}_0$, while N^{14} ground state is $(p_{1/2})^2{}_1$] use of Visscher-Ferrell wave functions suppresses the $L=0$ contribution by a factor of 5 000. The cancellation is not as complete as that in β decay (factor of 10^5 suppression). On the other hand, the $L=2$ contribution (shown in Fig. 6) is not suppressed relative to the single-configuration value. Comparison with measurements (Fig. 6) shows not only a poor fit to the shape of the angular distribution but also an a value three times larger than that for the higher 1^+ state. The spin- and charge-exchange operator therefore cannot fit the angular distribution nor give an a value consistent with that for the upper 1^+ state.

VI. CONCLUSIONS

Calculations with $b(\boldsymbol{\tau}_0 \cdot \boldsymbol{\tau}_i)$ adequately describe the shape of the $0^+ \to 0^+$ $C^{14}(p,n)$ angular distribution. The b value of 9 MeV can be compared to that deduced by Satchler[13] in fitting the $Zr(p,n)0^+ \to 0^+$ transition at 18.5 MeV. Satchler obtained a b value of 19 MeV which is much larger principally because of the shorter range of his Yukawa interaction. Making the range correction (from 1.4 F to 1.0 F) increases our b value from 9 to 22 MeV, which is then in good agreement with the Satchler value of 19 MeV.

The calculations presented have shown that a spin- and charge-exchange term $(\boldsymbol{\sigma}_0 \cdot \boldsymbol{\sigma}_i)(\boldsymbol{\tau}_0 \cdot \boldsymbol{\tau}_i)$ gives an a value of 7 MeV and a reasonable fit to the angular distribution for the transition to the 3.95-MeV 1^+ state. However, both the angular distribution and absolute cross section calculated on the basis of an interaction strength determined by fitting the upper state fail to fit the experimental data for the 1^+ ground state. The explanation of the data will require other spin-flip mechanisms such as particle exchange or tensor forces. Une *et al.*[14] have shown that for C^{13} the $L=2$ knockout-exchange amplitudes are greater than the direct ones, while for the $L=0$ transitions the exchange contribution is small. Since the $L=0$ transfers are inhibited as in β decay, only $L=2$ contributions are important for the ground-state transition. The exchange contribution from a Serber interaction would raise the ground-state cross section substantially. At the same time the 3.95-MeV 1^+ state, being produced almost entirely from $L=0$ transfer (the $L=2$ transfers are substantially retarded), is only slightly changed by the inclusion of exchange. A tensor interaction also affects $L=2$ contributions to a greater extent than $L=0$. Preliminary calculations, including a tensor interaction of the same strength and sign as the central spin-spin term ($a=7$ MeV), increase the ground-state cross section by about a factor of 3 while affecting the cross section for the 3.95-MeV state only slightly. It therefore appears possible to explain the cross section for both the upper and lower 1^+ transitions with inclusion of knockout and tensor interaction. Work is currently underway to determine whether consistency for these as well as other nonanalog transitions can be obtained in this manner. The heavy-particle stripping mechanism discussed by Banerjee and Pal[15] is another possible mechanism for contributing to the ground-state transition, but as yet there are no numerical calculations to indicate its relative importance.

---

[13] G. R. Satchler, Nucl. Phys. (to be published).
[14] T. Une, S. Yamaji and H. Yoshida, Progr. Theoret. Phys. (Kyoto) **35**, 1010 (1966).
[15] M. K. Banerjee and D. Pal, Nucl. Phys. **83**, 575 (1966).