Two-Body Effects and Neutron Polarization in the Reaction $^2H(\gamma,n)H$ at Low Energies

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The neutron polarization in the reaction $^2H(\gamma,n)H$ is calculated by incorporating the two-body charge- and current-density effects in the traditional theory. The calculations are carried out for a number of recent and old two-nucleon potentials. Clear evidence is provided that even at very low energies a serious discrepancy between theory and experiment exists.

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Ever since the calculation of the exchange-current operators due to the exchange of one pion and of vector mesons by Chemtob and Rho[1] and their successful application in the one-pion exchange approximation to the radiative n-p capture to explain the well-known 10% discrepancy in the cross section at thermal energies by Riska and Brown,[2] meson-exchange-current contributions have been included in calculations of magnetic moment,[3] radiative n-d captive,[4] electron-deuteron scattering,[5] and other nuclear structure calculations.[6] The contribution of the two-body charge density to the $E1$ operator had been first correctly calculated by Cambi, Mosconi, and Ricci,[7] who also studied the effect of the two-body charge and current densities in the deuteron photoabsorption sum rules. It has long been known that calculations and measurement on the polarization of the outgoing nucleons in the reaction $^2H(\gamma,n)H$,[8] especially at low energies,[9] will provide a rather sensitive way to examine the two-body effects. The object of this Letter is to make available the results of such a calculation and provide clear evidence that the two-body charge- and current-density effects, when included in the traditional theory, cannot account for the polarization data. The two-body charge effects are included here for the first time in such calculations.

The general notation and conventions regarding geometry employed here are the same as in Rustgi *et al.* (RZBA).[9] All the transitions, $E1$, $M1$, and $E2$, considered in the paper are included here. According to Chemtob and Rho,[1] the relevant exchange-current contributions to the transition amplitude may be written as

$$M^{\text{exch}}=\tfrac{1}{2}\mu_B\{(\vec{\tau}_p\times\vec{\tau}_n)_3[(\vec{\sigma}_p\times\vec{\sigma}_n)g_{\rm I}+T_{pn}{}^{\otimes}g_{\rm II}]+(\vec{\tau}_p-\vec{\tau}_n)_3[(\vec{\sigma}_p-\vec{\sigma}_n)h_{\rm I}+T_{pn}{}^{\ominus}h_{\rm II}]\},\tag{1}$$

where μ_B is the nucleon Bohr magneton and

$$T_{pn}{}^{\otimes}=(\vec{\sigma}_p\times\vec{\sigma}_n)\cdot\hat{r}\hat{r}-\tfrac{1}{3}(\vec{\sigma}_p\times\vec{\sigma}_n),\qquad T_{pn}{}^{\ominus}=(\vec{\sigma}_p-\vec{\sigma}_n)\cdot\hat{r}\hat{r}-\tfrac{1}{3}(\vec{\sigma}_p-\vec{\sigma}_n).\tag{2}$$

There is an isoscalar part as well which is neglected because it does not make any significant contribution to the present investigation. The functions $g_{\rm I}$, $g_{\rm II}$, $h_{\rm I}$, and $h_{\rm II}$ are real scalar functions of $|\vec{r}|=|\vec{r}_p-\vec{r}_n|$ and are given in Ref. 1 for the various meson exchange processes. For the $^2H(\gamma,n)H$ process, it is easy to show that

$$T_{pn}{}^{\otimes}=iT_{pn}{}^{\ominus},\quad(\vec{\tau}_p\times\vec{\tau}_n)_3=-i(\vec{\tau}_p-\vec{\tau}_n)_3,$$

and the two terms in the curly brackets in (1) can be combined into a single term:

$$M^{\text{exch}}=\mu_B[(\vec{\sigma}_p-\vec{\sigma}_n)(g_{\rm I}+h_{\rm I})+T_{pn}{}^{\ominus}(g_{\rm II}+h_{\rm II})],\tag{3}$$

where use has been made of the fact that for the ground state of the deuteron $T=0$.

On adding this operator to Eq. (4.1) of RZBA and carrying out a calculation for the amplitudes in the same fashion as described therein, one can calculate the contribution of these terms to the elements of the S matrix in Eqs. (18.1)–(18.3) of RZBA. These contributions amount simply to making the following changes in the radial integrals:

$$M_S\to M_S+2M_S{}^0-\tfrac{1}{3}(2\sqrt{2})M_D{}^0,\tag{4}$$

$$M_D\to M_D+2M_D{}^2-\tfrac{1}{3}(2\sqrt{2})M_S{}^2,\tag{5}$$

where the radial integrals M_S and M_D are defined in RZBA and

$$M_S{}^0=(\hbar/Mc)\gamma\int_0^\infty U\,{}^1\mathfrak{F}_0(kr)[g_{\rm I}+h_{\rm I}]dr,\tag{6}$$

$$M_S{}^2=(\hbar/Mc)\gamma\int_0^\infty U\,{}^1\mathfrak{F}_2(kr)[g_{\rm II}+h_{\rm II}]dr,\tag{7}$$

$$M_D{}^0=(\hbar/Mc)\gamma\int_0^\infty W\,{}^1\mathfrak{F}_0(kr)[g_{\rm II}+h_{\rm II}]dr,\tag{8}$$

$$M_D{}^2=(\hbar/Mc)\gamma\int_0^\infty W\,{}^1\mathfrak{F}_2(kr)[(g_{\rm I}+h_{\rm I})+\tfrac{1}{3}(g_{\rm II}+h_{\rm II})]dr.\tag{9}$$

In numerical calculations all the meson exchange diagrams consider by Hadjimichael[8] including the ones for the exchange of heavy mesons such as ρ and ω are considered.

The additional contribution due to the effect of the two-body charge density is included by adding to Eq. (4.1) of RZBA the following[7]:

$$\vec{D}^{(2)} = -(f^2/2M)Y_1(\mu r)\{[\mu_s \vec{\tau}_p\cdot\vec{\tau}_n + \tfrac{1}{2}\mu_v(\vec{\tau}_p+\vec{\tau}_n)_z]\hat{r}\times(\vec{\sigma}_p\times\vec{\sigma}_n) - \tfrac{1}{2}\mu_v(\vec{\tau}_p-\vec{\tau}_n)_z(\vec{\sigma}_p\vec{\sigma}_n\cdot\hat{r}+\vec{\sigma}_n\vec{\sigma}_p\cdot\hat{r})\}$$
$$-\tfrac{1}{2}(f^2/4M)(\vec{\tau}_p-\vec{\tau}_n)_z[Y_1(\mu r)(\vec{\sigma}_p\vec{\sigma}_n\cdot\hat{r}+\vec{\sigma}_n\vec{\sigma}_p\cdot\hat{r}) - \tfrac{1}{3}\hat{r}(e^{-\mu r}\vec{\sigma}_p\cdot\vec{\sigma}_n + Y_2(\mu r)S_{12})], \quad (10)$$

where S_{12} is the usual tensor operator, $f^2=0.081$, $\hat{r}=\vec{r}/r$, $\mu_s=\mu_p+\mu_n$, $\mu_v=\mu_p-\mu_n$, μ is the pion mass, and the functions Y_1 and Y_2 are defined by Cambi, Mosconi, and Ricci.[7] The contribution of these terms, for example, resulting in the $T=1$ final state, can be incorporated in the elements of the S matrix simply by making the following changes in the radial integrals in Eqs. (18.4)–(18.12) of RZBA:

$$I_0 = \gamma\int_0^\infty \left\{U\left[r+\frac{f^2}{6M}\varphi_0 - \frac{2}{3}\frac{f^2}{M}\varphi_2 - \frac{6}{M}\left(\mu_v\varphi - \frac{\varphi_1}{2}\right)\right] - \sqrt{2}W\left(r+\frac{f^2}{6M}\varphi_0 - \frac{2}{3}\frac{f^2}{M}\varphi_2\right)\right\}\mathfrak{F}_0{}^1(kr)\,dr, \quad (11)$$

$$I_1 = \gamma\int_0^\infty \left\{U\left[r+\frac{f^2}{6M}\varphi_0 + \frac{f^2}{3M}\varphi_2 + 4\frac{f^2}{M}\left(\mu_v\varphi - \frac{\varphi_1}{2}\right)\right] + \frac{W}{\sqrt{2}}\left[r+\frac{f^2}{6M}\varphi_0 + \frac{f^2}{3M}\varphi_2 - \frac{2f^2}{M}\left(\mu_v\varphi - \frac{\varphi_1}{2}\right)\right]\right\}\mathfrak{F}_1{}^1(kr)\,dr, \quad (12)$$

$$I_{u\tau}{}^2 = \gamma\int_0^\infty \left\{U\left(r+\frac{f^2}{6M}\varphi_0 - \frac{f^2}{15M}\varphi_2\right) - \frac{W}{5\sqrt{2}}\left[r+\frac{f^2}{6M}\varphi_0 - \frac{11}{3}\frac{f^2}{M}\varphi_2 - 30\frac{f^2}{M}\left(\mu_v\varphi - \frac{\varphi_1}{2}\right)\right]\right\}\mathfrak{U}_2{}^\tau\,dr, \quad (13)$$

$$I_{v\tau}{}^2 = \gamma\int_0^\infty \left[W\left(r+\frac{f^2}{6M}\varphi_0 - \frac{f^2}{3M}\varphi_2\right) + \sqrt{2}\,\frac{f^2}{3M}\varphi_2 U\right]\mathfrak{V}_2{}^\tau\,dr, \quad (14)$$

where $\tau=(\alpha,\beta)$.

These modified amplitudes are then incorporated in Eqs. (9.5)–(9.7) of RZBA to calculate the polarization of the outgoing neutrons. The prescription for obtaining the amplitudes from the outgoing neutrons from protons is described in RZBA. The numerical calculations are carried out for the Yale,[10] Hamada-Johnston,[11] Paris,[12] and supersoft-core[13] (versions B and C) potentials. The results are compared with the experimental data[14-16] in Figs. 1 and 2.

In Fig. 1, the results for the supersoft-core potential B (SSC-B) without and with the two-body effects are labeled as 1 and 2, respectively. The results for all other potentials, namely, the SSC-C, Paris, and Yale potentials, are reported only for calculations which have the two-body effects in them and are labeled as 3, 4, and 5, respectively. The results for the Hamada-Johnston potential were also calculated and fall between 4 and 5 but are not displayed. Our results for the Hamada-Johnston potential are higher than those of Hadjimichael,[8] who has included only the two-body current effects in his calculations.

There is no doubt that in Fig. 1, the results for the SSC-B give the best agreement with the data of Holt, Stephenson, and Specht[15] and Nath, Firk, and Schultz[14] when two-body effects are excluded (labeled as 1). In view of the large uncertainties in the data, and their rather scattered nature, it is not possible to claim that curve 2,

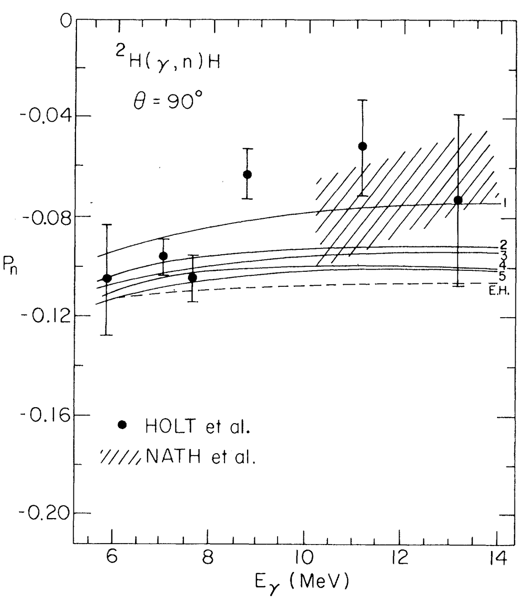

FIG. 1. Comparison of the present work with the data of Holt, Stephenson, and Specht (Ref. 15) indicated by solid circles, and Nath, Firk, and Schultz (Ref. 14) denoted by the hatched region. The dashed curve is the result of the calculation of Ref. 8 for the Hamada-Johnston potential and includes only the meson exchange currents. Curve 1 is the result of the present investigation for the potential SSC-B excluding two-body charge and current effects. Curves 2, 3, 4, and 5 display the results of the SSC-B, SSC-C, Paris, and Yale potentials, respectively, when two-body charge and current effects are included.

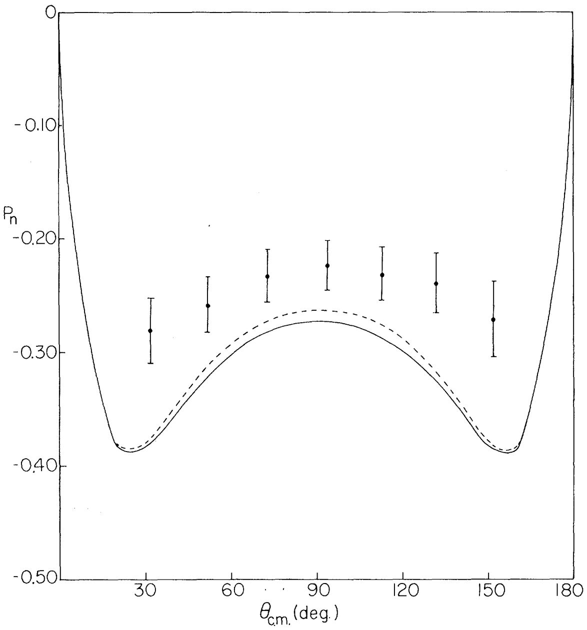

FIG. 2. Comparison of the present calculation for the SSC-B potential without (dashed curve) and with (solid curve) the two-body charge and current effects with the data of Jewell *et al.* (Ref. 16).

which has the two-body effects in it, does any worse. However, in general, it is found that the two-body contributions have the effect of increasing the polarization which is found to be negative. This is so because the neutron polarization from the reaction $^2H(\gamma,n)H$ at an angle of 90° and at low energy is known to be negative and dominated by $E1$-$M1$ interference terns.[9] The two-body current effects enhance the value of the radial integral defined in Eq. (4), resulting in more (negative) polarization.

The situation is, however, much clearer when comparison is made with the data of Jewell *et al.*[16] at 2.75 MeV. There is unambiguous disagreement between theory and experiment and the two-body effects make the agreement worse. The agreement can, however, be obtained if the solid curve is multiplied by an arbitrary factor of ~0.82. Similarly, all the published low-energy (2.5 to 16.0 MeV) angular distribution data, too many to be discussed in this Letter, can be fitted within reported error bars only with a normalization factor for each energy. The normalization factor is needed because the published cross sections are relative and not absolute. One cannot, however, draw any definite conclusions from work published earlier by Hadjimichael and Saylor[17] at higher energies (20–80 MeV) because their model, invented to fit the total cross sections, is not yet developed enough to fit the angular distribution and polarization data.

It may be mentioned that the results for the angular distribution and polarization with and without the two-body charge effects but including the exchange-current terms are hardly distinguishable from each other at the lowest γ-ray energy considered here, differing in the fourth decimal place by one or two. This indicates that in the low-energy limit our results agree with those including exchange-current terms as in Ref. 2.

The disagreement obtained in Fig. 2 in conjunction with the general discussion above and conclusions that can be drawn from Fig. 1 clearly demonstrate the inadequacy of the present theory for the $^2H(\gamma,n)H$ process.

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Coexistence of Single-Particle, Collective-Quadrupole, and $\alpha + ^{14}C$ Molecular-Dipole Degrees of Freedom in ^{18}O

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All natural-parity states of ^{18}O have been studied with high accuracy with the $^{14}C(^7Li, t\gamma)^{18}O$ coincidence and $^{14}C(\alpha,\gamma)^{18}O$ radiative-capture reactions. The four-particle, two-hole 0_2^+, 1^-, 2_3^+, and 3_3^- states deexcite with consecutive enhanced $E1$ and $E2$ cross-over transitions having $B(E1) \simeq 10^{-2}$ Weisskopf units (W.u.) and $B(E2) \simeq 20$ W.u. These data suggest the existence of an $\alpha + ^{14}C$ dipole band in ^{18}O similar to those discussed recently by Iachello and Jackson.

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The $A = 18$ system is an attractive one for study of nuclear structure[1] in that it gives access to both charge symmetry and charge independence of the nucleon-nucleon interaction, as well as interplay of the single-particle and deformed collective-quadrupole degrees of freedom.[2-4] Indeed, the coexistence of core-excitation deformed states [e.g., the four-particle, two-hole (4p-2h) 0_2^+ state at 3.63 in ^{18}O] and simple two-particle shell-model states is now well established in ^{18}O.

Recently, it has been suggested[5] that certain nuclei may display an altogether new collective degree of freedom. When the nucleus can be described as a dinuclear molecular system, as this suggestion implies, the relevant degree of freedom is the separation vector of the nuclear centers and the pertinent variables are the length of this vector and two of the three Euler angles which define its spatial orientation. The dinuclear molecular system can be described by a classical geometrical description[6] as well as a group-theoretical algebraic picture.[5] In this latter case the molecular spectra are considered to be generated by one S boson and three P_μ ($-1 \leq \mu \leq 1$) bosons, the generators of U(4).

Two quite distinct physical situations are possible. In the first, the participant nuclei do not themselves deform and only the length of the separation vector undergoes oscillation while the entire system can rotate about its center of mass. Such motion would be expected to lead to a conventional vibration-rotation spectra involving rotational bands having spin sequences 0^+, 1^-, 2^+, 3^-, 4^+, etc.; additionally it would be expected that enhanced collective $E1$ intraband transitions would be observed in non-self-conjugate systems.[7]

In the second physical situation the participant nuclei interpenetrate as the separation vector oscillates. Such motion can give rise to a spectrum having equidistant multiplets of dipole vibrational character and again enhanced collective $E1$ transitions (between states of different multiplets) are predicted in non-self-conjugate systems.[7]

Recently this enhancement of the radiative widths of transitions linking molecular states has been examined in a model-independent fashion[6] and sum rules have been derived for $E1$, $E2$, and $E3$ transitions. These sum rules together with the usual Wigner limits of reduced widths for particle decay of the presumed molecular states provide an effective signature for such structure as well as a measure of the degree of collective enhancement. They also provide a scale [10^{-2} Weisskopf units (W.u.)] for the $B(E1)$ enhancement.

The best candidate nuclear systems for exhibition of the suggested molecular states would ap-

