Self-weakening of the tensor interaction in a nucleus

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We examine several "landmarks" for the effects of the tensor interaction on the properties of light nuclei. These were usually discussed in the context of small-space shell-model calculations. We show, using G matrices derived from a realistic nucleon-nucleon potential, that when the model space is small (e.g., $0\hbar\omega$), these effects are overestimated, indicating that the tensor interaction is too strong. However, when larger spaces are used, there is a diminishing of these effects, which, in general, leads to better agreement with experiment.

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I. INTRODUCTION

The tensor interaction in a nucleus is a rather elusive beast. There are no first-order contributions to the binding energy of a closed LS nucleus like ^{4}He, ^{16}O, or ^{40}Ca due to the tensor interaction. Likewise, there are no first-order contributions to the single-particle energies for a closed LS shell plus or minus one nucleon. There are, however, important second-order contributions.

To see the effects of the tensor interaction in first order, the simplest thing to do is to go to systems with *two* quasiparticles. To this end, we will consider the following cases: (a) 1p-1h (one-particle–one-hole) system: the $J^\pi=0^-$, $T=0$, and $T=1$ states in ^{16}O, (b) 2h system: the beta decay of ^{14}C to ^{14}N, and (c) 2p system: the $E2$ and $M1$ moments of ^{6}Li. For these cases we perform shell-model calculations first in a small space ($0\hbar\omega$) and then in larger spaces in which $2\hbar\omega$ and sometimes even higher excitations are included.

We employ Brueckner reaction matrices G [1] calculated according to

$$G(E_s) = v_{12} + v_{12}\frac{Q}{E_s - (h_1 + h_2 + v_{12})}v_{12}, \qquad (1)$$

using the method introduced in Ref. [2]. In the above equation, v is the bare NN potential for which we adopt a new Nijmegen local NN interaction (NijmII) [3]; $h = t+u$ is the one-body Hamiltonian with u chosen to be the harmonic-oscillator (HO) potential $u(r) = \frac{1}{2}m\omega^2r^2$; E_s is the starting energy, which, for an initial two-particle state $|12\rangle$ in the ladder diagrams, is taken to be

$$E_s = \epsilon_1 + \epsilon_2 + \Delta\,, \qquad (2)$$

where ϵ's are the HO single-particle energies,

$$\epsilon_i = \left(2n_i + l_i + \frac{3}{2}\right)\hbar\omega \equiv \left(N_i + \frac{3}{2}\right)\hbar\omega\,, \qquad (3)$$

thus $(\epsilon_1 + \epsilon_2)$ is the unperturbed energy of the initial two-particle state $|12\rangle$ in the ladder diagrams. The parameter Δ in Eq. (2) can be thought of as the interaction energy between the two particles. The reader is referred to Ref. [4] for a partial justification for using the prescription in Eq. (2) for the starting energy. Note that such a state-dependent choice for E_s leads to a non-Hermitian G matrix but the non-Hermiticity is found to be small, and we here use an average of G and its Hermitian conjugate, $\frac{1}{2}(G+G^\dagger)$, as our effective interaction $v^{\rm eff}$. For all the nuclei under consideration, we fix the basis parameter $\hbar\omega$ at 14 MeV and the starting-energy parameter Δ at -50 MeV.

The Pauli operator Q in Eq. (1) is defined to prevent the two nucleons from scattering into the intermediate states which are either occupied (therefore Pauli forbidden) or inside the model space (to avoid double counting). It is therefore related to the choice of the model space. As we increase the size of the model space, we enlarge the $Q=0$ region to avoid double counting. We show in Fig. 1, as an example, our definition of Q which is used

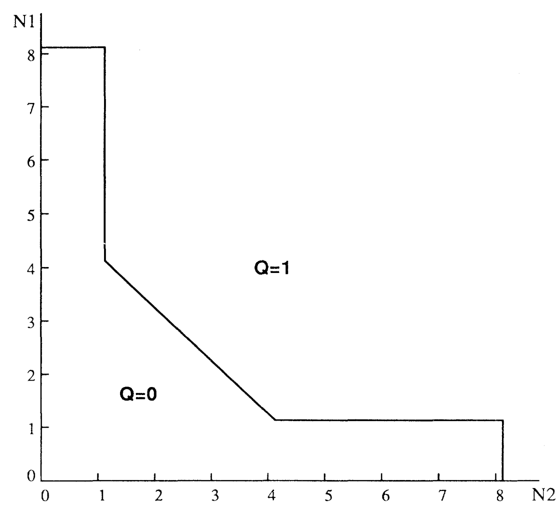

FIG. 1. The Pauli operator Q for a full $3\hbar\omega$ calculation of ^{16}O in a model space consisting of five major shells (s+p+sd+pf+sdg). $Q=0$ for $N_1 + N_2 \leq 5$ or $N_1 \leq 1$ or $N_2 \leq 1$. Here $N_i = 2n_i + l_i$ is the principal quantum number for the harmonic oscillator single-particle state $|i\rangle$. It is 0 for the lowest major shell (0s) and 1 for the 0p major shell, etc. The wings extend out to include the ninth (N=8) major shell.

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to calculate the G matrix for a full $3\hbar\omega$-space calculation of ^{16}O. We cut off the "wings" at the edge of the $N=8$ major shell. We emphasize that in this work we have used different definitions of Q for different model spaces and/or nuclei.

The matrix diagonalization is performed for the shell-model Hamiltonian

$$H_{\rm SM} = \left(\sum_{i=1}^{A} t_i - T_{\rm c.m.}\right) + \sum_{i<j}^{A} v_{ij}^{\rm eff} + \lambda(H_{\rm c.m.} - \tfrac{3}{2}\hbar\omega), \quad (4)$$

where $T_{\rm c.m.}$ is the center-of-mass (c.m.) kinetic energy and the last term (with $\lambda \gg 1$) is added to remove the spurious effects of the c.m. motion from the low-lying states. We have not included the Coulomb interaction. It should be pointed out that our calculations involve no phenomenological single-particle energies. These are implicitly generated from the two-body G matrix elements as well as the one-body kinetic energy in the matrix diagonalization.

II. THE 1p-1h CASE: ISOSPIN SPLITTING OF $J^\pi=0^-$ STATES IN ^{16}O

We previously considered in some detail the isospin splitting of $J^\pi=0^-$ states in ^{16}O [5]. We will here take the opportunity to discuss a few points. It was shown by Blomquist and Molinari [6] and Millener and Kurath [7] that without the tensor interaction, the energy difference ΔE between the $J^\pi=0_1^-$, $T=1$ and $J^\pi=0_1^-$, $T=0$ states in ^{16}O would be very small. Experimentally, the $J^\pi=0_1^-$, $T=0$ state is at 10.952 MeV and the $J^\pi=0_1^-$, $T=1$ is at 12.797 MeV so that the value of ΔE is 1.845 MeV. The results for ΔE using various model spaces, using the G matrices derived from the NijmII potential, are given in Table I.

In the $1\hbar\omega$ 1p-1h space, in which the dominant configuration is $(1s_{1/2}0p_{1/2})^{J^\pi=0^-}$ with small admixture of $(0d_{3/2}0p_{3/2})^{J^\pi=0^-}$, the value of ΔE is too large, i.e., 2.809 MeV as compared to the experimental value of 1.845 MeV. The situation is not improved much when $3\hbar\omega$ 3p-3h configurations are included in the matrix diagonalization, i.e., when, besides the $1\hbar\omega$ 1p-1h already there, we allow two additional nucleons to be excited from the $0p$ shell to the $1s$-$0d$ shell. The splitting ΔE changes from 2.809 MeV to 2.678 MeV; see Table I.

However, when $3\hbar\omega$ 2p-2h admixtures [i.e., $(0s)^{-1}(0p)^{-1}(1s0d)^2$ and $(0p)^{-2}(1s0d)^1(1p0f)^1$] are introduced, the situation improves dramatically: ΔE goes down to 1.641 MeV. This number is in better agreement with experiment, but we have an overshoot. When we furthermore include $3\hbar\omega$ 1p-1h configurations [i.e., $(0s)^{-1}(1p)^1$ and $(0p)^{-1}(2s1d0g)^1$] to make our model space complete for a full $(1+3)\hbar\omega$ calculation, the splitting becomes 1.931 MeV, in very good agreement with experiment.

This is one example of the "self-weakening" mechanism for the effective tensor interaction. The 2p-2h diagrams which contribute to ΔE in second-order perturbation theory are shown in Fig. 2. We classify them as particle-hole (or bubble) [Fig. 2(a)], hole-hole [Fig. 2(b)], and particle-particle [Fig. 2(c)] diagrams. In order to show how different parts of the interaction contribute to the isospin splitting ΔE, we present the perturbation-theory results in Table II for a schematic interaction which was introduced in Ref. [8] for an easy control of the strengths of the spin-orbit and tensor interactions:

$$V = V_c + xV_{\rm s.o.} + yV_t , \quad (5)$$

where c=central, s.o.=spin-orbit, t=tensor. For $x=1$ and $y=1$, the matrix elements of this interaction approximately resemble G matrix elements derived from a realistic NN potential like Bonn A [9]. By setting $x=0$ (1), we switch the spin-orbit interaction off (on); by setting $y=0$ (1), we switch the tensor interaction off (on).

As noted in Ref. [5], the particle-hole diagram [Fig. 2(a)] is the most important for getting a large, negative contribution for ΔE, but this only occurs when the tensor interaction is turned on. From Table II, we also note that the particle-particle and hole-hole diagrams are of the same sign and they act against the particle-hole diagrams. For $x=1$ and $y=1$, whereas the particle-hole diagrams contribute an amount of -1.728 MeV to the splitting ΔE, the particle-particle and hole-hole diagrams together give a contribution of $+0.901$ MeV, leading to a net result of $\Delta E=-0.826$ MeV.

We wish to note, and this point has not been made before, that for the particle-hole diagrams of Fig. 2(a), the most important contribution comes from the case in

TABLE I. The isospin splitting ΔE of the lowest 0^- states in ^{16}O obtained from shell-model diagonalizations in various model spaces. The binding energy (E_B) for the ground state and the excitation energy for the 0_1^-T states (T is the isospin) are also listed. The Coulomb interaction is not included. (The experimental binding energy listed in the table is Coulomb corrected.) All energies are in MeV.

Model space		E_B	$E(0_1^-,0)$	$E(0_1^-,1)$	ΔE
0^+: $0\hbar\omega$	0^-: $1\hbar\omega$ 1p-1h	118.933	15.943	18.752	2.809
0^+: $(0+2)\hbar\omega$	0^-: $1\hbar\omega$ 1p-1h + $3\hbar\omega$ 3p-3h	114.404	17.747	20.425	2.678
0^+: $(0+2)\hbar\omega$	0^-: $1\hbar\omega$ 1p-1h + $3\hbar\omega$ 3p-3h + $3\hbar\omega$ 2p-2h	124.692	15.946	17.587	1.641
0^+: $(0+2)\hbar\omega$	0^-: Full $(1+3)\hbar\omega$	124.692	15.646	17.577	1.931
Experiment		~ 142	10.952	12.797	1.845

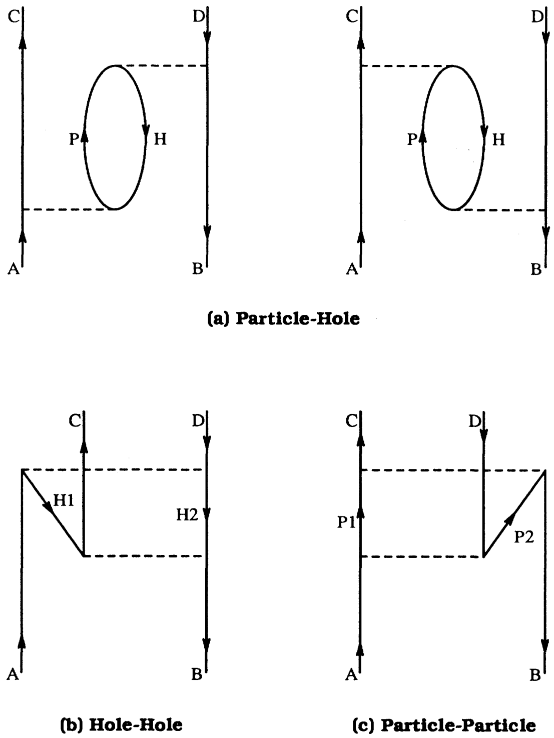

FIG. 2. The 2p-2h admixtures to the 1p-1h configuration: (a) particle-hole (bubble) diagrams, (b) hole-hole diagram, and (c) particle-particle diagram.

which the lower vertex involves a central interaction and the upper vertex involves a tensor interaction. Indeed, when the PH in Fig. 2(a) is equal to $(1s)(0s)^{-1}$, the lower vertex cannot involve a tensor interaction at all because the matrix element for the lower vertex involves only l=0 orbits. Thus Fig. 2(a) is approximately linear in the strength of the tensor interaction, and it tends to act *against* the first-order 1p-1h tensor interaction.

It should not be construed as a theorem that for any central, spin-dependent, interaction, the isospin splitting ΔE is small. Indeed, there is considerable discussion of this splitting in the classic paper by Elliott and Flowers [10]. The authors considered a central interaction of the following form:

$$V = V_0(W + MP_x - HP_\tau + BP_\sigma)\frac{\exp(-r/a)}{r/a}. \quad (6)$$

For the l=even, T=0, and S=1 channel ("deuteron channel"), V is proportional to $(W+M+H+B)$ which is -1 by convention. For the l=odd, T=0, and S=0 channel, V is proportional to $X = W - M + H - B$. It was found that ΔE is very sensitive to X. As X increases, the energy of the J^π=0^-, T=0 state increases linearly, and the energy of the J^π=0^-, T=1 state decreases linearly. These two states become degenerate when X is about 0.75. For a larger X, one gets the erroneous result that the T=1 state comes below the T=0 state. This is the case for the Rosenfeld interaction (X = 1.8) which has a strong p-wave repulsion. Therefore, if the central interaction in nature were close to a Rosenfeld interaction, one would need a strong tensor interaction as a corrective to invert the order of the T=0 and T=1 states. However, the Rosenfeld interaction was rejected on the basis of neutron-proton scattering data. At intermediate energies, the Rosenfeld interaction leads to a much larger differential cross section at 180° (in the center-of-mass frame) than at zero degree. Experimentally the two are nearly equal. So the data go more in the direction of a Serber interaction (for which X = 0) where the p-wave interaction vanishes.

We see from the above discussion that if one were allowed to adjust the exchange mixture at will, one would be able to obtain the experimental isospin splitting even with a strong tensor force. So it is vital to use realistic interactions before one can make quantitative, or even qualitative, statements about the role of the tensor interaction in a nucleus.

III. THE A=14 SYSTEM (2 HOLES): ALLOWED BUT INHIBITED BETA DECAY

Another landmark example of the effects of the tensor interaction in a nucleus is the famous A=14 beta decay: $^{14}C(J^\pi = 0^+, T = 1) \to ^{14}N(J^\pi = 1^+, T = 0)$. The quantum numbers involved in this transition are just right for an allowed Gamow-Teller (GT) transition, but the transition is very strongly suppressed. The matrix element B(GT) is essentially zero.

For the two holes in the $0p$ shell, the wave functions of the initial and final states in the LS coupling can be written as [8]

$$\psi_i = C_i^S\ |^1S_0\rangle + C_i^P\ |^3P_0\rangle\ ,$$
$$\psi_f = C_f^S\ |^3S_1\rangle + C_f^P\ |^1P_1\rangle\ + C_f^D\ |^3D_1\rangle\ . \quad (7)$$

It was shown by Inglis [11] (see also Ref.[12]) that it is im-

TABLE II. The contributions (in MeV) of the two-particle, two-hole perturbation-theory diagrams of Fig.2 to the isospin splitting of the lowest $J = 0^-$ states in ^{16}O.

Interaction		Diagram			
Spin-orbit	Tensor	Particle-Hole	Particle-Particle	Hole-Hole	Total
Off	Off	0.012	0.252	0.049	0.312
On	Off	-0.024	0.249	0.057	0.282
Off	On	-1.739	0.639	0.385	-0.714
On	On	-1.728	0.593	0.308	-0.826

possible to get $B(\mathrm{GT})=0$ with the above wave functions unless there is a tensor interaction present. Since the GT operator $\sum \sigma_\mu(i)t_+(i)$ cannot change the orbital angular momentum, one way (but not the only way) of getting $B(\mathrm{GT})=0$ would be to have $\psi_f = |^3D_1\rangle$. In general, the GT amplitude is $[B(\mathrm{GT}) = |A(\mathrm{GT})|^2]$

$$A(\mathrm{GT}) = \sqrt{6}\left(C_i^S C_f^S - \frac{1}{\sqrt{3}} C_i^P C_f^P\right). \tag{8}$$

The fact that the one-body spin-orbit force for holes is minus that for particles (whereas the hole-hole two-body interaction is equal to the particle-particle two-body interaction) leads to a large admixture of $|^3D_1\rangle$ in the $J^\pi=1^+$ ground state of ^{14}N.

Are higher-shell effects important here? To answer this question, we give in Table III the results for $B(\mathrm{GT})$ along with a few other observables, obtained in the $0\hbar\omega$ and $2\hbar\omega$ calculations using the NijmII [3] G matrices. [As we mentioned in Sec. I, when going from the $0\hbar\omega$ space to the $2\hbar\omega$ space, we have modified the Pauli operator Q in the G matrix equation (1) to exclude the model-space states from the intermediate spectrum.] As a contrast we also show in Table III our results for a one-hole system ($A=15$) for which there is much less sensitivity to the $2\hbar\omega$ configuration mixing.

The results in Table III for $B(\mathrm{GT})$ require some explanation. The value of $B(\mathrm{GT})$ for $A=14$ that we obtain in the $0\hbar\omega$ calculation is 3.967, much larger than zero. However, in terms of the tensor interaction, this means that we have an *overshoot.* As discussed in Ref. [8], the amplitude $A(\mathrm{GT})$ is large in magnitude when there is no tensor force. As we turn on the tensor interaction and gradually increase its strength, $A(\mathrm{GT})$ decreases in magnitude, goes through zero and changes sign. This happens well before we come to the full tensor strength. When we further strengthen the tensor interaction towards the full scale, $A(\mathrm{GT})$ becomes large in magnitude again. Therefore, the large result obtained for $B(\mathrm{GT})$ in the $0\hbar\omega$ space again indicates that the tensor interaction is too strong.

When we go from the $0\hbar\omega$ space to the full $2\hbar\omega$ space, the value of $B(\mathrm{GT})$ decreases by about 55% from 3.967 to 1.795. This latter value is still far from satisfactory, but it is much closer to the experimental answer of nearly zero.

It should be pointed out that, unlike the example of the isospin splitting between $J^\pi=0^-$ states in ^{16}O which we discussed in Sec. II, in the present case of the $A=14$ beta decay, the spin-orbit interaction also plays an important role. An alternate way of getting $B(\mathrm{GT})=0$ is to keep the tensor interaction at its full strength but *increase* the strength of the spin-orbit interaction. However, in a previous work [5], we have seen that the higher-shell effects do not have a significant effect on the spin-orbit interaction as they do on the tensor interaction.

The higher-shell admixtures also have a large effect on the magnetic dipole ($M1$) moment μ of the ground state ($J^\pi=1^+$, $T=0$) in ^{14}N, which is $0.768\mu_N$ in the $0\hbar\omega$ calculation and $0.554\mu_N$ in the $2\hbar\omega$ calculation. These values are obtained using the bare g factors: $g_l(p)=1$, $g_l(n)=0$, $g_s(p)=5.586$, $g_s(n)=-3.826$. The experimental result is $0.404\mu_N$ [13, 14].

For the $A=15$ system, we note that the higher-shell admixtures do not have a significant effect on $B(\mathrm{GT})$. We can understand this from the theorem which says that there are no first-order corrections to $B(\mathrm{GT})$, or to the $M1$ moment for a system consisting of a closed LS shell plus or minus one nucleon. For the two-hole system, on the other hand, the higher-shell effects can, in part, renormalize the particle-particle (or hole-hole) interaction between the two quasiparticles.

IV. THE $A=6$ SYSTEM (2 PARTICLES): THE MAGNETIC DIPOLE MOMENT AND ELECTRIC QUADRUPOLE MOMENTS OF ^{6}Li

One more landmark signature of the tensor interaction, although one that is somehow not recognized by many people, is the fact that the electric quadrupole ($E2$) moment of the ground state in ^{6}Li is *negative.* To show this, we again use the schematic interaction previously described. With the bare values of $e_p=1$ and $e_n=0$,

TABLE III. Properties of $A=14$ and $A=15$ nuclei from the $0\hbar\omega$-space and $2\hbar\omega$-space shell-model matrix diagonalizations. In the Table, we also give the binding energy $E_B(J^\pi,T)$ for the ground state and the excitation energy $E_x(J^\pi,T)$ for the excited state involved. Bare electromagnetic operators are used.

A	Observable	$0\hbar\omega$	$2\hbar\omega$	Expt.
14	$B(\mathrm{GT})(0^+\,1 \to 1^+\,0)$	3.967	1.795	~ 0
	$B(M1)(0^+\,1 \to 1^+\,0)$ (μ_N^2)	9.737	4.998	
	$\mu(1^+\,0)$ (μ_N)	0.768	0.554	0.404
	$Q(1^+\,0)$ $(e\,\mathrm{fm}^2)$	1.236	2.151	1.56
	$E_B(1^+,0)$ (MeV)	82.928	89.521	104.64
	$E_x(0^+,1)$ (MeV)	2.142	1.836	2.313
15	$B(\mathrm{GT})(\frac{1}{2}^- \to \frac{1}{2}^-)$	0.333	0.326	0.270
	$\mu(^{15}\mathrm{N})$ (μ_N)	-0.264	-0.277	-0.283
	$\mu(^{15}\mathrm{O})$ (μ_N)	0.638	0.655	0.719
	$E_B(\frac{1}{2}^+,\frac{1}{2})$ (MeV)	98.496	104.778	115.476
	$E_x(\frac{3}{2}^+,\frac{1}{2})$ (MeV)	4.102	5.444	6.324

TABLE IV. The results for the electric quadrupole moment (in $e\,\mathrm{fm}^2$) and magnetic dipole moment (in μ_N) and the binding energy (in MeV) of the ground state in ^{6}Li from $0\hbar\omega$, $2\hbar\omega$, and $4\hbar\omega$ shell-model calculations. Bare electromagnetic operators are used.

Space	Q	μ	E_B
$0\hbar\omega$	-0.360	0.866	26.49
$2\hbar\omega$	-0.251	0.848	27.58
$4\hbar\omega$	-0.0085	0.846	30.03
Expt.	-0.082	0.822	31.989

the quadrupole moment Q of the $J^\pi=1^+$ ground state of ^{6}Li assumes the following values when we fix the spin-orbit interaction at its full strength (x=1) and vary the strength (y) of the tensor interaction:

$$Q = \;\; 0.106\; e\,\mathrm{fm}^2 \quad \text{for} \quad y = 0\,,$$
$$Q = -0.135\; e\,\mathrm{fm}^2 \quad \text{for} \quad y = 0.5\,,$$
$$Q = -0.358\; e\,\mathrm{fm}^2 \quad \text{for} \quad y = 1\,.$$

That is to say, when the tensor interaction is switched off, the quadrupole moment is positive. As we increase the strength of the tensor interaction, Q decreases from being positive to being negative.

We also calculate the $E2$ moment Q of ^{6}Li using the NijmII G matrices. The results from the shell-model diagonalization are shown in Table IV for three model spaces: $0\hbar\omega$, $2\hbar\omega$, and $4\hbar\omega$. We also give the results for the $M1$ moment μ obtained using again the bare g factors.

In the $0\hbar\omega$ space, the calculated value of Q is $-0.360\; e\,\mathrm{fm}^2$, which is much more negative than the experimental value of $-0.082\; e\,\mathrm{fm}^2$. Again, this could be interpreted as being due to the fact that the effective tensor interaction in this $0\hbar\omega$ space is too strong. However, as we enlarge the model space, the magnitude of Q comes down. In the $2\hbar\omega$-space calculation, the value of Q is $-0.251\; e\,\mathrm{fm}^2$ and in the $4\hbar\omega$ space, there is an overshoot: we obtain $-0.0085\; e\,\mathrm{fm}^2$.

For the $M1$ moment μ, the experimental value is $0.822\mu_N$. Because the ground state of ^{6}Li has isospin zero, this is an isoscalar magnetic moment. The experimental value lies between the value for the jj limit ($0.627\mu_N$) and the value for the LS limit ($0.880\mu_N$). Going from the $0\hbar\omega$ to the $2\hbar\omega$ space, the calculated value of μ changes from $0.866\mu_N$ to $0.848\mu_N$. We are still above the experimental value even in the $4\hbar\omega$ calculation. The deviation is only $0.024\mu_N$ or 2.9%. However, for the *isoscalar* moments, one generally has a higher standard than for the isovector moments. The experimental isoscalar moments lie much closer to the Schmidt limit than do the isovector ones, and they are less sensitive to configuration mixing.

We feel that ^{6}Li deserves further study. It is the most elementary example of a system with two nucleons embedded in a nuclear medium. The medium corrections can be calculated to a higher precision than in heavier nuclei.

V. ADDITIONAL REMARKS

There are other approaches for dealing with light systems such as cluster calculations for ^{6}Li performed by Lehman *et al.* [15], Eskandarian *et al.* [16], and Schellingerhout *et al.* [17]. As compared to our shell-model results, the above authors seem to get better agreement for the isoscalar magnetic moment of the ground state, but the quadrupole moment of this state comes out positive. One can argue that the cluster approach is more physical. On the other hand, the Brueckner shell model is not self-limiting. With improved technology the calculations can always be extended in a systematic manner. This does not necessarily mean that perfect agreement with experiment will be reached since the assumption that we can describe a nucleus solely in terms of neutrons and protons interacting via a two-body interaction may not be correct. But to see this requires very high quality calculations.

Another point to be made is that, in contrast to this work where we claim that many anomalies in *nuclear structure* relating to the tensor interaction can be explained by simply admixing higher-shell configurations, in the realm of nucleon-nucleus scattering, there have been many works which put forth the idea that the tensor interaction in the nuclear medium must be considerably modified. These include experimental analyses of Hintz *et al.* [18] and Stephenson *et al.* [19]. They find that the polarization anomalies in proton-nucleus scattering can be removed by adopting the theoretical "universal scaling" ideas of Brown and Rho [20] that all mesons except the pion are less massive in the nuclear medium. Our next task will be to consider the fact that we do not seem to need medium modifications of the tensor interaction for nuclear structure, but we do need them for nucleon-nucleus scattering.

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