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STRETCHED EXCITATIONS AS A MEANS TO STUDY THE SPIN MODES OF THE NUCLEUS WITH ELECTROMAGNETIC AND HADRONIC PROBES

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Abstract - Excitation of spin transitions using electromagnetic and hadronic probes is discussed. Spectroscopic strengths are deduced from studies of (e,e'), (p,p'), (\pi,\pi') and (p,n) for selected states. The missing isovector and isoscalar strengths imply the existence of other important quenching mechanisms since the A-N~ mechanism does not contribute to high spin excitations or to isoscalar transitions. The stretched states are used as benchmarks to test the nucleon-nucleon tensor force in proton scattering and the pion-nucleon spin orbit force in pion scattering. The results are in striking agreement with predictions of the DWIA model.

INTRODUCTION

In this session on spin excitations in nuclei we are interested in the study of the nuclear response to nucleon operators of the \sigma type. It has been shown that effects on the spin-isospin operators due to pionic polarization of the nuclear medium are strongly momentum dependent. Studies of M1 and GT resonances have provided important and fundamental information on quenching mechanisms in the static limit of q = 0. Studies of high multipole magnetic resonances should provide valuable information on whether medium effects are important at high q. Yet, quenching effects such as delta particle-nucleon hole admixtures for the missing GT strength are not expected to strongly couple to the higher multipole magnetic resonances.

In this talk in this session I would like to present results from a study of the higher magnetic multipole resonances observed in (e,e') experiments at the Bates-MIT Linear Accelerator Center, (p,p') and (p,n) experiments at the Indiana University Cyclotron, and (\pi,\pi') experiments at the Los Alamos Meson Physics Facility that show high spin "stretched" states are also quenched compared to nuclear structure calculations by amounts similar to that observed in M1 and GT resonances. The absence of an explanation of the missing strength heightens the interest in the study of stretched excitations, particularly in view of their expected simple structure. For example, only one nucleon spin matrix element is expected to contribute to the electromagnetic and hadronic cross section, so orbital current contributions are expected in (e,e'), and two-body meson-exchange currents are small and calculable. Perhaps we have the opportunity here to study transitions where the
The principal quenching mechanism is mostly due to conventional nuclear structure. Since the subnucleonic effects are expected to be weak, attention can be focussed on the nucleon-nucleon correlations and core polarization effects. The presence of unexpectedly large ground state correlations will have general implications regarding most nuclear structure properties including M1 transitions. The stretched excitations are of further interest. Because of their simplicity they can be used as benchmarks to test the tensor component $V^T_T(q)$ of the nucleon-nucleon force and the spin orbit component, $t^S_T(q)$, of the pion-nucleon force in the nucleus. And ultimately we can use the hadronic reactions to deduce nuclear structure information on the isoscalar magnetic states which is unobtainable from $(e,e')$ data alone. Examples utilizing the complementary aspects of electromagnetic and hadronic reactions will be shown illustrating these features. The discussion will be in the spirit of material presented in Ref. 10.

Schematically, the relations between the weak, strong, and electromagnetic operators that are important in coupling the projectile to the nucleus in regions of momentum transfer that favor $l^+(q=0)$ and stretched high spin excitations ($q=2$ fm$^{-1}$) are given in Table 1.

Table 1. A tabulation of the weak, strong, and electromagnetic operators relevant for excitation of unnatural parity states of low and high spin.

<table>
<thead>
<tr>
<th>Probe</th>
<th>$J^\pi = 1^+$, $q=0$</th>
<th>$J^\pi = 4^-, 6^-, 8^-, \ldots$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\beta^-$-decay</td>
<td>$g_A^\sigma T_+ \sigma T_+$</td>
<td>$V^T_T(q) j_{J-1}(qr)(\bar{Y}<em>{J-1} \times \bar{\sigma}) T</em>+$</td>
</tr>
<tr>
<td>$(p,n)$</td>
<td>$V^C_T \sigma T_+ \sigma T_+$</td>
<td>$V^T_T(q) j_{J-1}(qr)(\bar{Y}<em>{J-1} \times \bar{\sigma}) T</em>+$</td>
</tr>
<tr>
<td>$(p,p')$</td>
<td>$V^C_T \sigma T_+ \sigma T_+$</td>
<td>$V^T_T(q) j_{J-1}(qr)(\bar{Y}<em>{J-1} \times \bar{\sigma}) T</em>+$</td>
</tr>
<tr>
<td>$(\pi,\pi')$</td>
<td>$t^S_T \sigma T_+ \sigma T_+$</td>
<td>$t^S_T(q) j_{J-1}(qr)(\bar{Y}<em>{J-1} \times \bar{\sigma}) T</em>+$</td>
</tr>
<tr>
<td>$(e,e')$</td>
<td>$g_s^\sigma T_+ \sigma T_+$</td>
<td>$g_s^\sigma j_{J-1}(qr)(\bar{Y}<em>{J-1} \times \bar{\sigma}) T</em>+$</td>
</tr>
</tbody>
</table>

The different sensitivities to different regions of $q$ arises because the higher order spherical Bessel function $j_{J-1}(qr)$ for $J>1$, when folded in with the nuclear wave function and the interaction, samples the nucleus at higher regions of $q$.

Another important observation that distinguishes the stretched states from the magnetic dipole states is that the former only involve the nuclear spin for the probes in Table 1, while the latter also includes the orbital momentum in electron scattering. To illustrate the different sensitivities in $(e,e')$ and $(p,p')$ for excitation of $1^+$ states, Fig. 1 shows a comparison of $B(M1)\uparrow$ strength obtained from $(e,e')$ with the mirror image cross sections obtained from $(p,p')$ taken at $\theta_{Lab} = 4^\circ$.
Fig. 1 - A comparison of B(M1) strength obtained from (e,e') with mirror image cross sections obtained from (p,p') for $^{50}$Ti, $^{52}$Cr, and $^{54}$Fe.

for several different even-even nuclei in the fp shell. The differences are most striking for the states in $^{52}$Cr. Although some of these differences in extracted strengths arise because of different sensitivities of the spin operator to neutron and proton components (or isoscalar and isovector), the additional orbital component for proton excitations in (e,e') cannot be neglected. A most dramatic comparison is shown for $^{51}$V in Fig. 2. The structure near 10 MeV excitation observed in (p,p')

Fig. 2 - A comparison of (e,e') and (p,p') spectra at low q for the odd A nucleus $^{51}$V.

at 3° is not at all visible in the electron spectra. A possible explanation is that the weaker excitations in (e,e') arises from possible spin and orbital destructive interference. Such a possibility becomes increasingly more important with increasing $\xi$. The single particle magnetic dipole matrix element for the spin and orbital operators are plotted versus $\xi$ in Fig. 3. At $\xi = 4$ the proton orbital recoupling matrix element $<\xi+1/2|\xi|\xi+1/2>$ is about 14 nuclear magnetons, which is as large as the isovector spin-flip matrix element between spin-orbit partners $<\xi+1/2|\xi-1/2>$. Spin and orbital destructive interference is also a possible explanation why strong M1 transitions are not observed in (e,e') in heavy nuclei like $^{90}$Zr and $^{208}$Pb. For closed
shell nuclei such a result implies that there would be significant ground state correlations breaking the simple shell model structure in order for this mechanism to be important. Calculations of this type should be performed in order to test the validity of this possibility.

Since the interest in this session is focussed on the spin modes of nuclear excitations, we shall avoid discussing multipole magnetic transitions that knowingly contain orbital contributions to the cross sections, which will eliminate the ambiguities discussed above.

ELECTRON, PROTON, AND PION INELASTIC SCATTERING CROSS SECTIONS TO STATES OF STRETCHED CONFIGURATION

Inelastic electron, proton, and pion scattering cross sections for unnatural parity excitations involve the magnetic multipoles of the electromagnetic interaction, the spin dependent central, tensor, and spin-orbit components of the nucleon-nucleon interaction, and the spin-orbit component of the pion-nucleon interaction and thus depend on the isoscalar and isovector spin and orbital transition densities of the target nucleus. The general expressions for the cross section have been given in Ref. 10 and 14. Here, we summarize in plane wave Born approximation (PWBA) the expressions for the inelastic differential cross section for $0^+ \rightarrow J^0$ "stretched" excitations in electron, proton, and pion-nucleus scattering. It is assumed that the "stretched" particle-hole configuration is of the form $(j_a^{-1} j_b)_{J^0_{max}}$, where $J_{max} = j_a + j_b$, $J_a = j_a + 1/2$, $J_b = j_b + 1/2$, and $j_a$ and $j_b$ are the largest angular momentum found in the last filled shell and the first open shell, respectively. This configuration is unique in a space excluding lp-lh excitations with $E \neq 3\hbar\omega$, so there is no mixing with other lp-lh configurations. Of course, the "stretched" configurations can mix with multi-nucleon-multi-hole configurations within the same shell. Such mixing produces physical ground and excited states that are not pure closed shell and particle-hole wave functions, respectively; however, the additional components in the wave functions will not be connected by the one body spin and orbital current operators. The only effect is then a reduction of the transition strength with the cross sections remaining proportional to a single particle matrix element corresponding to the stretched configuration. An additional simplification that arises is that the orbital transition density and the spin transition density associated with the $J+1$ order Bessel function vanish as a result of angular momentum restrictions on the single particle matrix element.

For inelastic magnetic electron scattering the PWBA expression for the differential cross section is

$$\frac{d\sigma}{d\Omega} = \frac{Z^2 \sigma_n}{\pi} \left( \frac{1}{2} + \tan^2 \frac{\theta}{2} \right) r_T^2(q)$$

(1)
where $q_M$ is the Mott cross section, $\eta$ is a recoil factor, $\theta$ is the scattering angle, $q$ is the momentum transfer, $F_T^2$ is the transverse form factor, which is given by

$$F_T^2(q) = \frac{4\pi}{Z^2} (J+1) \left| \frac{z}{2mc} \right|^2 \frac{1}{2} \gamma_\tau \rho_{J\gamma-1}(q) \right|^2$$

(2)

for transitions to pure isovector or isoscalar stretched states where $\tau = 1$ or 0, respectively and $\gamma^\tau$ are the isovector and isoscalar spin g-factors, respectively. An explicit expression for the spin transition density in momentum space is given by

$$\rho_{J\gamma-1}(q) = (-1)^\gamma b \left( 4\pi \right)^{-1/2} C_{J\gamma} \beta_{J\gamma} \beta_{J\gamma-1} \left| j_a^\gamma j_b^\gamma \right| \left| j_a^\gamma j_b^\gamma \right|$$

(3)

where $C_{J\gamma} = \sqrt{2}$ for $\gamma = 0, 1$. For harmonic oscillator radial wave functions and the stretched state conditions $n_a = n_b = 1$, $\gamma_a = \gamma_b + 1$, and $J = \gamma_a + \gamma_b$, the radial overlap in eq. (3) is given by

$$\left| n_a^\gamma j_a^\gamma \right| \left| n_b^\gamma j_b^\gamma \right| = \frac{(bq)^{-1/2} \exp(-b^2 q^2/4)}{2(J-1)! J^1!! \left((J-1)!!(J+1)!!\right)^{1/2}}$$

(4)

where $b$ is the oscillator length parameter. A useful analytic expression for the transverse electron scattering form factor is obtained by substituting eq. (3) and eq. (4) into eq. (2), obtaining

$$F_T^2(q) = \frac{\rho_{cm}(q) \rho_{NS}(q)}{Z} \frac{z}{2mc} \frac{1}{2} C_{J\gamma} \left| j_a^\gamma j_b^\gamma \right| \left| j_a^\gamma j_b^\gamma \right|$$

(5)

The terms $\rho_{cm}$ and $\rho_{NS}$ are center of mass and nucleon finite size correction factors, respectively. The coefficients $Z_{J\tau}$ are spectroscopic amplitudes proportional to the one body density matrix elements

$$Z_{J\tau} = \left| \left( J_{T f} \right) \left| \left( a_{j_a} \right. \left| a_{j_b} \right\right) \right| \right|^2$$

(6)

where $T_{f}, T_{f},$ and $\tau$ are the initial state, final state, and transition isospin, respectively.

For inelastic proton nucleus scattering the corresponding PWBA expression for the differential cross section to stretched states in the spirit of the impulse approximation is given by

$$\frac{d\sigma}{d\Omega} = 4\pi \frac{M_N}{2\pi q^2} \left( J |\varepsilon_{\tau}(q)|^2 + (J+1) |\varepsilon_{\tau}(q)|^2 + |\varepsilon_{\tau}(q)|^2 \right)$$

(7)

where $M_N(k_p)$ is the reduced energy (wave number) of the incident nucleon, the factor $\alpha$ arises from the decomposition of the relative nucleon-nucleon momentum $p = p_p - \alpha p_p - \alpha_T p_T$, $\varepsilon_{\tau}(q)$ refers to the spin orbit component of the effective nucleon-nucleon interaction, and
are the longitudinal and transverse combinations of the central and tensor interaction components. The bars on the interaction components are to indicate that they are defined to include, approximately, contributions associated with knockout exchange.\textsuperscript{16-18} The component $\mathbf{v}_S$ is solely due to tensor exchange.\textsuperscript{16-18}

For inelastic pion-nucleus scattering at pion incident energies near the (3,3) resonance the corresponding PWBA expression for the differential cross section to stretched states is given by\textsuperscript{14,193}

$$d\sigma^\pi/d\omega = 4\pi\left(\frac{\pi}{2}\right)^2 (J+1)\frac{1}{2} k_0^4 q^{-\delta} \sin^2\theta \alpha^2 t L_S(q) t_L S(q) S_{J=1}(q)^2$$

(10)

where $M_1$ and $K_1$ are the reduced energy and incident wave number of the incident pion, $\alpha_1$ is as in eq. (7) and $t L_S$ is the isovector ($\tau=1$) or isoscalar ($\tau=0$) spin-orbit component of the effective pion-nucleon interaction. In this framework pion-nucleus scattering determines uniquely the transverse spin density. In the important incident energy region of the (3,3) resonance, the pion-nucleon interaction is very non-local. This gives rise to corrections in eq. (10) which produce coupling to orbital current and spin-current transition densities.\textsuperscript{21,22} These terms are not expected to be important for the transition of stretched configurations near the peak of the angular distributions.

In the calculation of the inelastic proton and pion cross sections the center of mass and nucleon finite size correction factors are not included in the spin density $\rho_{J=1}^S(q)$. Instead the center of mass correction is made by modifying\textsuperscript{23} the harmonic oscillator parameter $b$ by $(A/A-1)^{1/2}$ and the overall normalization factor by $\left[\frac{(A-1)}{A}\right]^{1/2}$.

In electron scattering $g_S^S>g_O^S$ so this reaction provides mainly information on transitions where $\rho_{J=1}^S$ is large. For $q = 2$ fm$^{-1}$, the approximate momentum transfer for observing high spin excitations, $v_1$ is dominant,\textsuperscript{14,16-18} so nucleon-nucleus scattering is also preferentially sensitive to isovector spin-flip states of high spin. This reaction also provides information on isoscalar high spin states because both $v_1$ and $v_2$ are appreciable at $q = 2$ fm$^{-1}$ with most of $v_0$ coming from knockout exchange.\textsuperscript{14,16-18} For the important incident energy region corresponding to the (3,3) resonance, $t_{O=S}^S = 2t_{O=S}^S$ so pion-nucleus scattering favors isoscalar transitions.

### STRETCHED TRANSITIONS VIA ELECTRON SCATTERING

The high spin stretched magnetic transitions are readily observed in electron scattering at high $q$-2 fm$^{-1}$ and at a backward angle $\theta>140^\circ$ where non-magnetic transitions are retarded. In light p shell nuclei stretched $M_4$ ($d_5/2P_3/2$) transitions have been observed in all stable nuclei from $^{12}$C through $^{16}$O. The distribution of $M_4$ strength\textsuperscript{24,25,26,27} in $^{12}$C, $^{13}$C, $^{14}$C, and $^{15}$N as observed in electron scattering at $\theta = 180^\circ$ is shown in Fig. 4. The decomposition of the distribution of strength in terms of spin and isospin is important to study because it provides information on the spin and isospin dependent parts of the effective interaction in nuclei by providing well-defined systematics that should be reproduced by any realistic model. The nucleus $^{13}$C illustrates the feature that only one strong isovector stretched transition is observed\textsuperscript{24,26} in nuclei whose ground states have zero isospin and zero angular momentum. This is also true for $M_4$ transitions\textsuperscript{24,28} in $^{16}$O and $M_6$ transitions\textsuperscript{31,32,33} in $^{24}$Mg and $^{28}$Si. A second important feature illustrated\textsuperscript{28} by $^{14}$C is that in neutron excess nuclei with ground state isospin $T_0$, the strength is fragmented into several $4^-$ states with isospin $T_0$ and one strong $4^-$...
Fig. 4 - Electron spectra from targets $^{12}$C, $^{13}$C, $^{14}$C and $^{14}$N at $q = 2$ fm$^{-1}$ and $\theta = 180^\circ$ illustrating preferential excitation of M4 transitions.
state with $T + 1$. A similar distribution of $M_6$ and $M_8$ strength has been observed\textsuperscript{34,35} in $^{26}$Mg and $^{54}$Fe in the sense that one $T_0 + 1$ state is observed and several $T_0$ states are observed. The Hillenburg-Kurath\textsuperscript{36} $1\hbar\omega$ shell model calculation on $^{14}$C and the Lawson\textsuperscript{35} $(f_7/2^3p_{3/2})_6^-$ and $(g_9/2^1f_{7/2})_8^-$ shell model calculations for $^{26}$Mg and $^{54}$Fe, respectively, are in rough agreement with these observations that only one $T_0 + 1$ level should be observed. This appears to be the $J^\pi = 4^-, T=2$ level at 24.3 MeV in $^{14}$C, the $J^\pi = 6^-, T=2$ level at 18.1 MeV in $^{26}$Mg, and the $J^\pi = 8^-, T=2$ level at 13.26 MeV in $^{54}$Fe.

A third feature in which $^{14}$N serves as the only example\textsuperscript{26,27} so far is that for ground states with zero isospin and spin 1, the isovector $M_4$ strength is fragmented into the expected $3^-$, $4^-$, and $5^-$ states. And in $^{13}$C where we have neither zero angular momentum nor zero isospin in the ground state the $M_4$ strength\textsuperscript{25} appears to fragment into $J = 7/2$, $9/2$ and $T = 1/2$, $3/2$ states which poses a stronger theoretical challenge.

Because these stretched levels are located in the continuum, and difficult to unravel from unresolved neighboring levels, some care must be taken in identifying multipolarities and deducing strengths. There is always the possibility that more than one level is being observed. In addition, many of the levels are unbound to either proton or neutron emission and, therefore, continuum wave functions\textsuperscript{38} in a realistic finite potential well such as the Wood-Saxon form (WSWF) should be considered. With exception of testing the sensitivity of the cross section to Wood-Saxon radial shapes for related examples, the question of realistic wells is ignored here. Instead, the results of a systematic study of most of the strong isovector levels observed in $(e,e')$ based on the use of harmonic oscillator radial wave functions (HOWF) is presented. With this assumption the form factor for the $1\hbar\omega$ stretched particle-hole configurations has the simple $q$ dependence, indicated in eq. (5),

$$F_T^2(q) = q^{2J} \exp(-b^2q^2/2)$$

(11)

which peaks at $q = \sqrt{2J}/b$. This simplicity is an important asset in the identification of the multipolarity. The square of the transverse form factor $F_T^2$ is shown versus $q_{eff}$ for several transitions in various nuclei in Fig. 5. Distortion effects due to the electron-nucleus coulomb attraction are taken into account by correcting the momentum transfer using the relation

$$q_{eff} = q(1 + 3\frac{Za_c}{2}\frac{E_0}{E_R})$$

(12)

The curves through the data points in Fig. 5 were determined by a least square fit of the form factor calculated assuming the extreme single particle-hole (ESPHM) model. The oscillator parameter, $b$, and a normalization factor, $S_2^{ESPHM} = Z^2^{(exp)}/Z^2^{(ESPHM)}$ were used as fitting parameters. The resulting values of $b$ and the normalizing factor $S_2^2$ are summarized for several isovector transitions in Table 2. The values obtained for $b$ are close to that expected from the $A^{1/6}$ rule. It is no surprise that the factors $S_2^{(ESPHM)}$ are consistently much less than one. These factors are intended to serve as a unit in which to conveniently compare deduced spectroscopic strengths. $S_2^{(ESPHM)} = 1$ corresponds to the maximum isovector strength expected in a single level in the ESPHM. For example in $^{160}$ the $M_4 (d_5/2p_3/2)$ strength in the 18.98 MeV level is $S_2^{(ESPHM)} = 0.41$ or 41\% of the ESPHM prediction.

To test the sensitivity of the extracted $S_2^2$ to shapes of the radial wave functions, calculations of the form factors were performed using bound state WSWF. For $^{14}$C, $^{14}$N, $^{28}$Si, and $^{54}$Fe, $S_2^{(ESPHM)}$ was found to be 0.51, 0.61, 0.37, and 0.57 for the corresponding transitions listed in Table 2. To test the $q_{eff}$ approximation in nuclei $^{160}$, $^{26}$Mg, $^{28}$Si, and $^{54}$Fe, DWBA calculations, using the code HEIMAG,\textsuperscript{43} modified to use HOWF, were performed by fitting the cross sections instead of the form factors. The normalization factors were the same as those deduced in PWBA fits to within 3\%, but the size parameter $b$ was systematically larger by 2-4\%. 

\begin{table}
\centering
\begin{tabular}{|c|c|c|c|c|}
\hline
Nucleus & $b$ & $S_2^{ESPHM}$ & $S_2^{(exp)}$ & $S_2^{(WSWF)}$ \\
\hline
$^{14}$C & 1.9 & 0.41 & 1.0 & 0.51 \\
$^{14}$N & 2.1 & 0.61 & 1.0 & 0.61 \\
$^{28}$Si & 1.8 & 0.37 & 1.0 & 0.37 \\
$^{54}$Fe & 2.0 & 0.57 & 1.0 & 0.57 \\
\hline
\end{tabular}
\caption{Values of the oscillator parameter $b$ and the normalizing factors $S_2^{ESPHM}$, $S_2^{(exp)}$, and $S_2^{(WSWF)}$.}
\end{table}
There is ample evidence that the ground state orbits are not fully occupied as assumed in ESPHM. The column labeled S2(THY) contains the results with Z2(THY) deduced from various theoretical models that take into account core polarization effects and correlations in the ground state. For example, Millener-Kurath wave functions (MKWF) are used to calculate the M4 strength in 12C, 14C, 14N, and 16O. In 12C, 14C, and 14N the complete p shell space is used to calculate the ground state wave function. The space is further enlarged to allow one nucleon in the s-d shell to calculate the odd parity states. Some provision is also made to allow for 2ω excitations from the 1s to 1p shell, but not from the 1p to the 2s1d shell. The results of this calculation for the three strong 4− states in 14C are shown in Table 3.

Significant improvement is obtained for the T=1 states at 11.7 and 17.3 MeV using the MKWF. Still, however, the calculated strength is about 40% too large. Very little improvement is obtained for the T=2 state at 24.3 MeV. When two-body meson-exchange currents are included in the calculation of the M4 form factor, the calculated cross section increases by about 15% and, therefore, inclusion of MEC increases the subsequent disagreement between experiment and theory by about 15%. The inability to fit the M4 strengths clearly indicates a more realistic treatment of the excited state and ground state is required, for example, allowing for 3ω configuration admixtures in the excited state as well as 2p-4h ground state admixtures. In 16O where better agreement is obtained, multi-particle-multi-hole con-

Fig. 5 - The square of the transverse factor F^2_T(q) is shown versus q_eff for several stretched transitions of different multipolarity in various nuclei.
Table 2. Deduced spectroscopic strength $S^2$ from $(e,e')$ for stretched states with strong isovector components.

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>$E_x$ (MeV)</th>
<th>$J^\pi, T$</th>
<th>CONF</th>
<th>b</th>
<th>$S^2$ (ESPHM) $^+$</th>
<th>$S^2$ (THY)</th>
<th>Ref.</th>
<th>$Z^2$ (ESPHM)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{12}$C</td>
<td>19.50</td>
<td>4$, 1$</td>
<td>$d_{5/2}^{-1}$</td>
<td>1.50</td>
<td>0.37±0.04</td>
<td>0.60</td>
<td>24</td>
<td>1</td>
</tr>
<tr>
<td>$^{14}$C</td>
<td>24.30</td>
<td>4$, 2$</td>
<td>$d_{5/2}^{-1}$</td>
<td>1.54±0.08</td>
<td>0.37±0.07</td>
<td>0.42</td>
<td>8</td>
<td>1/2</td>
</tr>
<tr>
<td>$^{14}$N</td>
<td>16.91</td>
<td>5$, 1$</td>
<td>$d_{5/2}^{-1}$</td>
<td>1.54±0.03</td>
<td>0.39±0.02</td>
<td>0.41</td>
<td>27</td>
<td>11/27</td>
</tr>
<tr>
<td>$^{16}$O</td>
<td>18.98</td>
<td>4$, 1$</td>
<td>$d_{5/2}^{-1}$</td>
<td>1.63±0.03</td>
<td>0.41±0.02</td>
<td>0.71</td>
<td>29, 30</td>
<td>1</td>
</tr>
<tr>
<td>$^{24}$Mg</td>
<td>15.05</td>
<td>6$, 1$</td>
<td>$f_{7/2}^{-1}$</td>
<td>1.85±0.04</td>
<td>0.27±0.02</td>
<td>0.47</td>
<td>31</td>
<td>2/3</td>
</tr>
<tr>
<td>$^{26}$Mg</td>
<td>18.1</td>
<td>6$, 1$</td>
<td>$f_{7/2}^{-1}$</td>
<td>1.82±0.03</td>
<td>0.40±0.02</td>
<td>0.41</td>
<td>34</td>
<td>1/3</td>
</tr>
<tr>
<td>$^{28}$Si</td>
<td>14.36</td>
<td>6$, 1$</td>
<td>$f_{7/2}^{-1}$</td>
<td>1.77±0.02</td>
<td>0.31±0.01</td>
<td>0.55</td>
<td>33</td>
<td>1</td>
</tr>
<tr>
<td>$^{48}$Ca</td>
<td>9.28</td>
<td>8$, 4$</td>
<td>$g_{9/2}^{-1}$</td>
<td>1.78±0.03</td>
<td>0.22±0.01</td>
<td>----</td>
<td>40</td>
<td>1</td>
</tr>
<tr>
<td>$^{54}$Fe</td>
<td>13.26</td>
<td>8$, 2$</td>
<td>$g_{9/2}^{-1}$</td>
<td>1.90±0.02</td>
<td>0.51±0.02</td>
<td>0.72</td>
<td>35</td>
<td>3/8</td>
</tr>
<tr>
<td>$^{58}$Ni</td>
<td>12.50</td>
<td>8$, 2$</td>
<td>$g_{9/2}^{-1}$</td>
<td>1.93±0.03</td>
<td>0.18±0.01</td>
<td>0.31</td>
<td>39</td>
<td>1/2</td>
</tr>
<tr>
<td>$^{90}$Zr</td>
<td>7.36</td>
<td>10$, 5$</td>
<td>$h_{11/2}^{-1}$</td>
<td>2.08±0.04</td>
<td>0.26±0.03</td>
<td>----</td>
<td>41</td>
<td>1</td>
</tr>
<tr>
<td>$^{208}$Pb</td>
<td>6.43</td>
<td>12$, 22$</td>
<td>$j_{15/2}^{-1}$</td>
<td>2.18±0.05</td>
<td>0.54±0.04</td>
<td>1.00</td>
<td>42</td>
<td>1</td>
</tr>
<tr>
<td>$^{208}$Pb</td>
<td>7.02</td>
<td>12$, 22$</td>
<td>$j_{13/2}^{-1}$</td>
<td>2.30±0.03</td>
<td>0.56±0.02</td>
<td>1.00</td>
<td>42</td>
<td>1</td>
</tr>
<tr>
<td>$^{208}$Pb</td>
<td>6.74</td>
<td>14$, 22$</td>
<td>$j_{15/2}^{-1}$</td>
<td>2.22±0.04</td>
<td>0.56±0.03</td>
<td>1.00</td>
<td>42</td>
<td>1</td>
</tr>
</tbody>
</table>

$^+$Harmonic oscillator parameter b and normalization $S^2$ were determined by a least square fit of the DWBA calculated cross section to the data. The center of mass and finite size effects have been included in the form factor during the fitting procedure.

$^+S^2 = Z^2(\text{EXP})/Z^2(\text{ESPHM})$ where $Z^2(\text{ESPHM})$ is defined in eqs. (3) and (6). For pure proton or neutron configurations, $Z^2=1$.

Table 3. A comparison of the deduced isovector ($d_{5/2}P_{3/2}$) M4 strength with the predictions of the ESPHM and a truncated shell model using MKWF.

<table>
<thead>
<tr>
<th>$E_x$ (MeV)</th>
<th>$J^\pi, T$</th>
<th>$S^2$ (ESPHM) without MEC</th>
<th>$S^2$ (ESPHM) with MEC</th>
<th>$S^2$ (MKWF) without MEC</th>
<th>$S^2$ (MKWF) with MEC</th>
</tr>
</thead>
<tbody>
<tr>
<td>11.7</td>
<td>4$^-$ 1</td>
<td>0.19</td>
<td>0.16</td>
<td>0.59</td>
<td>0.51</td>
</tr>
<tr>
<td>17.3</td>
<td>4$^-$ 1</td>
<td>0.22</td>
<td>0.19</td>
<td>0.63</td>
<td>0.53</td>
</tr>
<tr>
<td>24.3</td>
<td>4$^-$ 2</td>
<td>0.37</td>
<td>0.32</td>
<td>0.42</td>
<td>0.36</td>
</tr>
</tbody>
</table>

Figurations are allowed both in the ground state and the excited state. In $^{24}$Mg and $^{28}$Si open shell RPA calculations predict too much M6 strength by a factor of two. Again the ground state is treated realistically allowing for 1d$_{5/2}$.2s$_{1/2}$, 1d$_{3/2}$ admixtures while the excited state only allows for one nucleon in the f$_{7/2}$ orbit. Recent shell model calculations by Amusa and Lawson which use a space of
Fig. 6 - Electron spectra from targets of $^{54,56}\text{Fe}$ and $^{58,60}\text{Ni}$ illustrating preferential excitation of M8 transitions.
the type \((d_5/2s_1/2)^{11} f_7/2\) yields similar quenching factors. In the extreme deformed limit of the collective model, the \(M_8\) strength is fragmented over seven rotational bands from \(K=0\) to \(K=6\). Each \(K=0\) rotational band would get \(2/13\) of the extreme single particle-hole model predictions. Although this deformed limit is unrealistic for \(^{28}\text{Si}\), it does suggest a mechanism for producing fragmentation.\(^{46,47}\) More realistic fragmentation has been predicted by Emery\(^{46}\) by introducing Coriolis coupling between bands but predicts too much strength in a single state for realistic deformations.

In the \(fp\) shell the fragmentation of stretched \(M_8\) strength is being studied as a function of proton and neutron number. Electron scattering measurements have been taken on \(^{54}\text{Fe}\), \(^{56}\text{Fe}\), \(^{58}\text{Ni}\), and \(^{60}\text{Ni}\), at Bates and also on \(^{52}\text{Cr}\) at NIKEF.\(^{49}\) A comparison of spectra are shown in Fig. 6 indicating the striking selectivity with which the \(8^-\) levels are excited. Form factors have been extracted for \(^{54}\text{Fe}\), \(^{58}\text{Ni}\), and \(^{60}\text{Ni}\), and analysis is still in progress on \(^{56}\text{Fe}\). Two characteristic features of the data are: (1) the total \(M_8\) strength is about \(1/3\) of the ESPHM prediction for each nucleus and (2) the \(M_8\) strength is split into two groups with isospin \(T_0\) and \(T_0 + 1\). This feature is particularly emphasized for \(^{56}\text{Fe}\) and \(^{60}\text{Ni}\) where there is a clear energy gap \(\Delta E = \nu_0 (T_0 + 1)/A\), which is consistent with an asymmetry energy\(^{50}\) of \(\nu_0 = 109 \pm 10\ \text{MeV}\), which is close to the simple Lane model value of 100 MeV. Among these nuclei, the strongest \(M_8\) transition occurs for the \(E_X = 13.26\ \text{MeV} \ T=2\) state in \(^{54}\text{Fe}\), which gives \(S^2(\text{ESPHM}) = 0.51\). Calculations by Metsch\(^{48}\) which include configurations of the type \((g_9/2p_3/2f_7/2)^{11} \ 8^-\) using the modified surface delta interaction (MSDI) yield \(S^2(\text{THY}) = 0.72\) for the same transition. However, the comparison between experiment and theory is still poor for the \(T=1\) states in \(^{54}\text{Fe}\) as is shown in Fig. 7. Even when the \(2p-4h\) configurations are included, only two \(T=1\) \(8^-\) states are predicted and are much too strong. Only those calculated \(8^-\) states that contain more than 1% of the \(M_8\) strength are shown in Fig. 7. In \(^{58}\text{Ni}\) calculations\(^{48}\) using MSDI in a shell model space containing particle-hole configurations of the type \(r g_9/2\) and \(r^2 g_9/2 f_7/2\), where \(r\) stands for one of the orbits \(2p_3/2\), \(2p_1/2\), and \(1f_5/2\) yield \(S^2(\text{THY}) = 0.31\) for the strongest \(T=2\) \(M_8\) transition in \(^{58}\text{Ni}\). Results for \(T_0 \rightarrow T_0 \ M_8\) transitions in \(^{58}\text{Ni}\) are also predicted to be too large by factors of two to three. More sophisticated calculations are required since existing results predict far too much strength.

In \(^{208}\text{Pb}\) one might expect the shell model assumptions to be most adequate, measurements of the \((e,e')\) cross sections to the \(J^\pi = 12^-\), \(E_X = 6.43\) and 7.02 MeV states and to the \(J^\pi = 14^-\), \(E_X = 6.74\ \text{MeV} \) state are about half the values predicted by the ESPHM as shown in Table 2. In deducing the strength,\(^{49}\) it is assumed that the structure of each level is given by the \(v_{j15/2} l_{13/2} \ n_{13/2} \ n_{11/2}\) and \(v_{j15/2} l_{13/2} \ n_{13/2} \ n_{11/2}\) configurations.
configurations. This assumption is confirmed by recent core polarization and particle-vibration calculations. The core polarization calculations by Hammamoto et al.⁵¹ which include a sum of $1p-1h$, $1\hbar\omega$, 2$\hbar\omega$, 3$\hbar\omega$, etc., configurations also predict the correct strength. However, the calculations employ a zero range delta force ($V_0 = 170 \text{ MeV-fm}^2$), which is known to be too strong at large $q$ in comparison to realistic interactions. Recent core polarization calculations⁵² including the effects of meson and rho exchange of the $12^-$ and $14^-$ electron scattering form factors by Suzuki predict too much strength by factors of two and do not fit the shape of the form factor. On the other hand particle-vibration coupling calculations by Krewald and Speth⁵³ show a 50% reduction in M12 and M14 strength to the $E_x = 6.43$, 7.02, and 6.74 MeV levels in agreement with experiment. However, the missing strength in these states appears as additional M12 and M14 strength in a group of states at 8.0–9.0 MeV excitation, which have not been identified experimentally. Clearly, the situation in $^{208}\text{Pb}$ needs further experimental and theoretical effort. It does appear, however, that the best agreement between theory and experiment obtained to date for these high spin states occurs for $^{208}\text{Pb}$.

For most of the $T_0 + 1$ excitations in Table 2 there is improved agreement between theory and experiment when the model space is enlarged to include more realistic ground state wave functions. Using the ESPHM, the average $S^2 = 0.38$ and for the more realistic theories $S^2 = 0.66$ with HWF. Extracted strengths are 15% to 20% higher in light nuclei when WSWF are employed and about 15% lower when MEC corrections are included. If both effects were simultaneously included, the results would be in effect similar to those in Table 2. Somewhat larger, spectroscopic strengths have been observed for the related stretched $0\hbar\omega$ excitations observed in $(p,n)$ charge exchange experiments.⁵⁴ The spectroscopic strengths from Table 2 are, however, in contrast with the much larger single particle strengths extracted for the largest (stretched) contributing multipole in elastic magnetic electron scattering⁵⁵ on odd $A$ targets with ground state spins $j_z = j + 1/2$. The difference in the inelastic and elastic magnetic strengths suggests that ground state and excited state correlations are very important, since the interfering backward going amplitude in RPA calculations for inelastic transitions would not be present in elastic scattering.⁵⁵ For the $T_0$ stretched magnetic excitations in non-self-conjugate nuclei, the situation appears more complex. For example, there is even greater fragmentation and less strength found experimentally than predicted by current theory (see Refs. 8, 34, 35, and 36) regarding $^{14}\text{C}$, $^{26}\text{Mg}$, $^{51}\text{Fe}$, and $^{58}\text{Ni}$ with the exception of $^{208}\text{Pb}$. Subnuclieonic mechanisms like $A$ particle-nucleon hole admixtures in the wave functions,⁶ which have been proposed to explain the quenching of the Gamow-Teller strength in nuclei, do not seem to be a likely explanation for the quenching of high spin states. There is always the possibility that there are many more weak transitions that cannot be discerned from the background. Although this may be true, still current theory is far from predicting such large amounts of fragmentation.

**HIGH SPIN STRETCHED STATES AS BENCHMARKS: COMPARISON WITH HADRONIC PROBES**

Although these pure spin transfer, stretched transitions are not yet completely understood in terms of current nuclear structure models, they are still very useful as "benchmarks" in testing direct, one step inelastic scattering models for hadronic reactions such as $(p,p')$, $(p,n)$, and $(n,p')$. The reason for this is that within the spirit of the DWIA model, the same spin transition density appears in the transition amplitude for all of these reactions as was displayed in eq. (2), eq. (7), and eq. (10). Distortion effects are taken into account through the use of DWBA instead of PWBA and the exchange amplitudes for the nucleon-nucleon reactions are included exactly through the use of the code DWBAP70. The comparisons of interest here provide information on the isovector tensor part of the effective nucleon-nucleon interaction and the spin-orbit component of the effective pion-nucleon interaction. To illustrate the former, the modulus of $S_1^V(q)$, $S_2^V(q)$ and $S_{12}^V(q)$ for 140 MeV protons as deduced from the free nucleon-nucleon amplitudes in Ref. 16 is shown in Fig. 8. An important feature of the interaction of Ref. 16 is that the radial dependence is given as sums of Yukawa terms with the long range parts of the central and tensor
Fig. 8 - Modulus of the isovector 140 MeV Love-Franey t-matrix in momentum space.

components matched to OPEP. Clearly, \( J^\pi_T(q) \) is dominant in the range of \( q \) where the electron scattering form factors displayed in Fig. 5 peak. This shows very nicely that the comparison of \((e,e')\) and \((p,p')\) for isovector stretched state transitions provides information on the strength and \( q \) dependence of the isovector tensor interaction in the approximate range \( q = 1-2.5 \) fm\(^{-1}\). For pictures of the pion-nucleon spin-orbit interaction at the \((3,3)\) resonance, the reader is referred to Refs. 14, 19, and 22.

The differential cross section for the \( 4^-,T=1 \) and the \( 6^-,T=1 \) isovector stretched states in \( 160 \) and \( 28\text{Si} \), respectively, have been measured using the \((e,e'),(p,p'),(p,n),\) and \((\pi,\pi')\) reactions. The data and calculations\(^{20,61,62,63}\)

for the \( 6^-,T=1 \) state in \( 28\text{Si} \) are shown as a function of \( q \) in Fig. 9. The transition density oscillator parameter \( b \) and normalization \( S^2 \), corresponding to fitting the cross section to the peak in Fig. 10 are summarized in Table 4 together with several other transitions. The value of \( S^2 \) from the \((p,n)\) and \((p,p')\) reaction is about 15% lower than the values of \( S^2 \) from \((\pi,\pi')\) and \((e,e')\). Additional results for the \( 6^-,T=1 \) level in \( 28\text{Si} \) observed in \((p,p')\) at \( E_p = 333 \) and 500 MeV give \( S^2 = 0.25 \) and 0.29.\(^{68}\) This indicates the energy dependence of the isovector tensor interaction as given by the Love-Franey nucleon-nucleon t matrix is approximately correct in the \( \text{q} \) region sampled in these comparisons.\(^{18,91}\) This is not the case for other effective interactions as discussed by Emery.\(^{46}\) Similar systematics are obtained for the \( 4^-,T=1 \) state in \( 160 \) in Table 4 with slightly more scatter than in \( 28\text{Si} \). Using the Love-Franey nucleon-nucleon t-matrix\(^{16}\) for \( 28\text{Si} \) and \( 54\text{Fe} \) and that of Ref. 70 for \( 160 \) and \( 24\text{Mg} \), \( S^2(p,p')/S^2(e,e') \) averages 0.84.

Near the peak of the form factor
Table 4. A comparison of $S^2$ deduced from hadronic probes to $S^2$ from $(e,e')$ for strong isovector stretched states.

<table>
<thead>
<tr>
<th>Target, $p$, $p'$</th>
<th>$E_X$ (MeV)</th>
<th>$J^\pi, T$</th>
<th>CONF</th>
<th>$S^2$ (ESPHM)</th>
<th>$b$</th>
<th>Ref.</th>
<th>$S^2/(x,x')$ $^{+\dagger}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{16}_0$</td>
<td>18.98</td>
<td>$4^-, 1$</td>
<td>$(d_{5/2}^{-1}/2^3/2)$</td>
<td>0.33</td>
<td>1.70</td>
<td>56, 57</td>
<td>0.81±0.04</td>
</tr>
<tr>
<td>$^{16}_0$</td>
<td>18.98</td>
<td>$4^-, 1$</td>
<td>$(d_{5/2}^{-1}/2^3/2)$</td>
<td>0.50</td>
<td>1.62</td>
<td>20, 58</td>
<td>1.22±0.06</td>
</tr>
<tr>
<td>$^{16}_0$</td>
<td>18.98</td>
<td>$4^-, 1$</td>
<td>$(d_{5/2}^{-1}/2^3/2)$</td>
<td>0.50</td>
<td>1.62</td>
<td>20, 58</td>
<td>1.22±0.06</td>
</tr>
<tr>
<td>$^{4}_0$</td>
<td>6.37</td>
<td>$4^-, 1$</td>
<td>$(d_{5/2}^{-1}/2^3/2)$</td>
<td>0.50</td>
<td>1.72</td>
<td>59</td>
<td>0.80±0.06</td>
</tr>
<tr>
<td>$^{24}_0$</td>
<td>15.05</td>
<td>$6^-, 1$</td>
<td>$(f^{-1}/2/2^3/2)$</td>
<td>0.24</td>
<td>1.86</td>
<td>60</td>
<td>0.88±0.06</td>
</tr>
<tr>
<td>$^{28}_0$</td>
<td>14.36</td>
<td>$6^-, 1$</td>
<td>$(f^{-1}/2/2^3/2)$</td>
<td>0.24</td>
<td>1.78</td>
<td>61</td>
<td>0.78±0.03</td>
</tr>
<tr>
<td>$^{28}_0$</td>
<td>14.36</td>
<td>$6^-, 1$</td>
<td>$(f^{-1}/2/2^3/2)$</td>
<td>0.32</td>
<td>1.76</td>
<td>20, 62</td>
<td>1.03±0.03</td>
</tr>
<tr>
<td>$^{28}_0$</td>
<td>14.36</td>
<td>$6^-, 1$</td>
<td>$(f^{-1}/2/2^3/2)$</td>
<td>0.32</td>
<td>1.76</td>
<td>20, 62</td>
<td>1.03±0.03</td>
</tr>
<tr>
<td>$^{28}_0$</td>
<td>14.36</td>
<td>$6^-, 1$</td>
<td>$(f^{-1}/2/2^3/2)$</td>
<td>0.25</td>
<td>1.74</td>
<td>63</td>
<td>0.81±0.03</td>
</tr>
<tr>
<td>$^{48}_0$</td>
<td>9.31</td>
<td>$8^-, 4$</td>
<td>$(g^{-1}/2/2^3/2)$</td>
<td>0.22</td>
<td>1.90</td>
<td>64</td>
<td>1.00±0.05</td>
</tr>
<tr>
<td>$^{54}_0$</td>
<td>13.26</td>
<td>$8^-, 2$</td>
<td>$(g^{-1}/2/2^3/2)$</td>
<td>0.45</td>
<td>1.92</td>
<td>65, 66</td>
<td>0.88±0.03</td>
</tr>
<tr>
<td>$^{208}$</td>
<td>6.74</td>
<td>$14^-, 22$</td>
<td>$(i_{15/2}^{-1}/2^{13/2})$</td>
<td>0.50</td>
<td>2.30</td>
<td>67</td>
<td>1.12±0.06</td>
</tr>
<tr>
<td>$^{208}$</td>
<td>6.43</td>
<td>$12^-, 22$</td>
<td>$(i_{15/2}^{-1}/2^{13/2})$</td>
<td>0.80</td>
<td>2.30</td>
<td>67</td>
<td>0.67±0.05</td>
</tr>
<tr>
<td>$^{208}$</td>
<td>7.02</td>
<td>$12^-, 22$</td>
<td>$(i_{13/2}^{-1}/2^{11/2})$</td>
<td>0.20</td>
<td>2.30</td>
<td>67</td>
<td>2.88±0.88</td>
</tr>
</tbody>
</table>

$^+$ the center of mass correction factor has been included in the final determination of $S^2$ and $b$.

$^{+\dagger}$ errors on ratios only include uncertainties in $S^2/(e,e')$. 

for these pure isovector transitions we are only sensitive to the tensor force. These results could imply that the free nucleon-nucleon tensor force is too strong since the extracted strength obtained from $(p,p')$ is 16% weaker than that from $(e,e')$. However, a word of caution concerns the correction to the one body spin density extracted from $(e,e')$ data due to two body meson exchange currents. Meson-exchange current calculations for $^{16}_C$ using Millener-Kurath and ESPHM wave functions show that the meson exchange current contribution has the effect of reducing $S^2/(e,e')$ by about 15%. From this, one might conclude that the spectroscopic factor $S^2$ deduced from $(e,e')$ should be reduced before comparing with $S^2$ deduced from hadron scattering. This would bring the electron results in better agreement with the Love-Franey parametrization of the tensor force. Inelastice proton scattering studies to stretched $6^-$ states in $^{28}_N$ are also in progress. The overall general consistency in results for $(e,e')$, $(p,p')$, and $(n,n')$ discussed so far is quite striking and is supportive of the DWIA description of the hadronic scattering process. This gives us a good measure of the strength of the tensor and spin-orbit components of nucleon- and pion-nucleon interactions in the range of $q$ near the peak cross section in a nuclear medium.

The results for the two $12^-$ states in $^{208}_Pb$ are not consistent. One of these levels is primarily an $i_{13/2}^{-1}/2^{11/2}$ proton configuration and the other is primarily a $j_{15/2}^{-1}/2^{13/2}$ neutron configuration. Only the first is stretched and since there are two levels, this is not a clean example of the unique stretched excitations which
have been the primary topic of discussion here. It was shown in Ref. 68 that a consistent explanation of the \((e,e')\) and \((p,p')\) data for these \(12^+\) levels could not be achieved by considering configuration mixing between the two levels alone. A possible problem is that the proton cross section for the \(6.43\) MeV neutron \(12^+\) level may contain a contribution from an unresolved \(13^+\) level. For pure neutron excitations like the \((89/2^+_7/2^-)\) in \(^{48}\)Ca and the \((j_{15/2}^+;_{13/2}^-)\) in \(^{208}\)Pb, the isoscalar spin-orbit and tensor contribute as well as the isovector tensor component of the force. It would be highly desirable to determine the isoscalar spin density from an independent experiment in order to isolate the effects of the force. Although agreement between the \((e,e')\) and \((p,p')\) result for \(^{48}\)Ca indicates that the difficulty may not be in the uncertainties in the force. Since \(^{208}\)Pb is such an important shell model nucleus, it is essential that it be examined more carefully and completely with regard to the high spin states.

There are some difficulties in fitting the \((p,p')\) and \((p,n)\) at higher \(q\) for the \(6^-\) in \(^{28}\)Si and the \(4^-\) in \(^{16}\)O. The possible appearance of a second maximum in \(^{16}\)O(p,p') and \((p,n)\) suggest several possibilities. The tensor component gets rapidly weaker at high \(q\) and the isovector spin-orbit interaction which is not as well determined begins to dominate. Effects due to the shapes of radial wave functions and the nuclear medium may also be important here. The local nucleon-nucleon \(t\) matrix excludes medium effects such as Pauli blocking, consideration of the Fermi motion of target nucleons, and the information on the off-shell nature of the nucleon-nucleus interaction. Careful studies of the energy dependence of the \(t\)-matrix and other effective interactions which include medium effects related to spin observables are required and are underway and may become more important at higher \(q\).

CONCLUSIONS

From a comparison of \((e,e')\), \((p,p')\), and \((p,n)\) excitations of \(1^+\) states at low \(q\), it is evident that there are significant spin and orbital interference effects, which, presumably, could be used to derive nuclear orbital current information as well as spin. On the other hand, excitations of stretched states via electromagnetic and hadronic probes exclusively provide information on the nuclear spin currents. Comparison to state of the art structure calculations indicate that there is significant quenching of the isovector spin-flip strength to high spin stretched states (38 to 66%). The absence of subnucleonic quenching mechanisms suggests that the missing strength may be found in conventional nuclear configuration mixing calculations that more realistically treat nuclear correlations, polarization, and nuclear deformation effects. Since similar isovector spectroscopic strengths are also obtained from proton and pion scattering, the missing strength is unrelated to any specific probe or whether the interaction is strong or electromagnetic. After approximately correcting the electron data for two-body meson-exchange currents, agreement between \((e,e')\) and \((p,p')\) for selected isovector stretched states is within a few per cent using the Love-Franey pa-rametrization of the tensor force. This particular form is very close to that predicted using appropriate mixtures of one pion exchange and rho exchange microscopic forces. Continuation of the complimentary studies extending the nuclear arena to incorporate the isoscalar magnetic observables should allow for a broader interpretation of quenching phenomena and a more fundamental picture of nuclear behavior.

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