# Measurements of Gamow-Teller Strength Distributions in Masses 13 and 15

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The differential cross section and the transverse spin-flip probability have been measured for the dominant transitions in $^{13}C(p,n)^{13}N$ and $^{15}N(p,n)^{15}O$ at $E_p = 160$ MeV. The Gamow-Teller transition strengths deduced from the data show that the major $\frac{1}{2}^- \rightarrow \frac{3}{2}^-$ transitions are strongly quenched relative to the $\frac{1}{2}^- \rightarrow \frac{1}{2}^-$ mirror transitions, in strong disagreement with simple shell-model expectations.

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Allowed Fermi and Gamow-Teller beta-decay transition rates provide a special class of nuclear model information because of the simple relationship between the model description of the nucleus and the transition process. The Fermi (F) operator changes only the isospin projection of a nucleon. The Gamow-Teller (GT) operator changes the projection of both isospin and spin. Mirror-state transitions between $T = \frac{1}{2}$ nuclei have been a favored testing ground because of the model simplification resulting from the fact that the parent and daughter states differ only in isospin projection. The transition rate between mirror states is the incoherent sum of the rates for the Fermi and Gamow-Teller components. All of the Fermi strength appears in the mirror-state transition, but, because of the spin-orbit interaction, the GT strength is distributed between the spin-orbit pair states. Typically, only a fraction of the total GT strength is contained in the mirror-state transition. For $L$-$S$ closed-shell $\pm$ 1 nuclei the model is further simplified, and that fraction is easily calculable.

A systematic survey has shown that the GT transition rates deduced from the measured $ft$ values are in most cases smaller than the calculated rates.[1] This evidence of missing GT strength in mirror transitions has been known for many years.[2] In most cases that can be explored through beta decay, the fraction of the total GT strength that appears in the mirror transition is small. Therefore, deductions about missing strength are more model sensitive than would be the case if a large part of the sum-rule strength were seen.

The $(p,n)$ reaction can be used to explore the distribution of the remaining strength,[3] and studies using the $(p,n)$ reaction have shown that the missing GT strengths is a general feature of nuclear structure. For strong GT transitions typically (50–60%) of the calculated strength is actually observed.[4] A notable exception to this observation occurs in mass 13. The ratio of measured[1] to calculated[5] GT strength for the $\frac{1}{2}^- \rightarrow \frac{1}{2}^-$ mirror transition is 0.66. This ratio is consistent with the "typical" GT quenching factor. However, the calculated $B$(GT) value of the strongest transition,[5] that from the ground state of $^{13}$C to the $\frac{3}{2}^-$, 3.51-MeV level in $^{13}$N, is $B(\mathrm{GT})=2.38$, while the value deduced from $(p,n)$ measurements[6] is $B(\mathrm{GT})=0.85\pm0.03$, a discrepancy of nearly a factor of 3. The calculated $B$(GT) summed over all levels is 3.95. The discrepancy, if parametrized in terms of "quenching" of GT strength, suggests a stronger quenching of the $\frac{1}{2}^- \rightarrow \frac{3}{2}^-$ transition than the $\frac{1}{2}^- \rightarrow \frac{1}{2}^-$ transition.

The essence of the problem is exposed even more clearly in the mass-15 data. Mass 15 is only one nucleon removed from the simultaneous $L$-$S$ and $j$-$j$ shell closure at mass 16. If we hold only to the restriction that the model space be limited to the $p$ shell, unlike the situation in mass 13, the total GT strength in mass 15 and the distribution of strength between the $p_{1/2}$ and $p_{3/2}$ hole states are independent of the spin-orbit splitting and of the two-body residual force. The total strength in this model is $B(\mathrm{GT})=3$, with $B(\mathrm{GT})=\frac{8}{3}$ going to the $p_{3/2}$ hole state and $B(\mathrm{GT})=\frac{1}{3}$ going to the $p_{1/2}$ hole state.

We have recently measured the spin-flip probability $S_{NN}(0°)$ for $(p,n)$ reactions on $^{13}$C and $^{15}$N. As explained below, this observable can be used to obtain a determination of the GT strengths for excited states in $^{13}$N and $^{15}$O that is independent of the method used in Ref. 6. In addition, we point out that both methods are independent of the absolute normalization of the $(p,n)$ cross sections and make use only of relative cross sections, which can be determined quite reliably.

The cross sections and transverse spin-flip probabilities for $^{13}\mathrm{C}(p,n)$ and $^{15}\mathrm{N}(p,n)$ were measured at 0° by means of a 160-MeV polarized proton beam from the Indiana University Cyclotron Facility and a neutron polarimeter consisting of bars of plastic scintillator at the end of a 60-m flight path. The polarimeter is described briefly elsewhere[7] and will be described more fully in a future publication. The targets were pressed wafers of carbon (> 95% $^{13}$C) and Melamine ($C_3H_6{}^{15}N_6$, > 99% enrichment in $^{15}$N). The results of our measurements are displayed in Table I, and the $(p,n)$ spectra are shown in Fig. 1. The carbon contribution to the Melamine-target spectrum was subtracted by making use of data obtained under the same experimental conditions with a natural carbon target. Table I also shows calculated and experimentally deduced values of $B$(GT). The shell-model calculations are based on the assumption that these nuclei can be characterized as $p$-shell nuclei.[5] The ground-state–to–ground-state $B$(GT) values are deduced from beta-decay $ft$ values. We have used values from Raman *et al.*[1]

Excited-state transition strengths can be extracted by first decomposing the ground-state cross section into Fermi Gamow-Teller parts by use of the observed relationship between GT and F transitions[6]:

$$\sigma_{\mathrm{GT}}/\sigma_{\mathrm{F}} = [E_p/(55\pm1\ \mathrm{MeV})]^2[B(\mathrm{GT})/B(\mathrm{F})], \quad (1)$$

TABLE I. Cross sections, spin-flip probabilities, and GT transition strengths for $^{13}\mathrm{C}(p,n)$ and $^{15}\mathrm{N}(p,n)$ at $\theta=0°$ and $E_p=160$ MeV.

| Final state | $\sigma_{\mathrm{c.m.}}(0°)$ [a] (mb/sr) | $S_{NN}(0°)$ | $B(GT)$ Expt. | Model |
|---|---|---|---|---|
| $^{13}\mathrm{N}(0.0,\frac{1}{2}^-)$ | $4.2\pm0.1$ | $0.46\pm0.2$ | $0.206\pm0.004$ [b] | 0.323 [e] |
| $^{13}\mathrm{N}(3.51,\frac{3}{2}^-)$ | $10.5\pm0.1$ | $0.66\pm0.02$ | $0.83\pm0.03$ [c] | 2.38 |
| | | | $0.75\pm0.05$ [d] | |
| $^{15}\mathrm{O}(0.0,\frac{1}{2}^-)$ | $4.5\pm0.1$ | $0.53\pm0.03$ | $0.261\pm0.006$ [b] | $\frac{1}{3}$ |
| $^{15}\mathrm{O}(6.18,\frac{3}{2}^-)$ | $10.8\pm0.1$ | $0.70\pm0.03$ | $1.00\pm0.03$ [c] | $\frac{8}{3}$ |
| | | | $0.80\pm0.06$ [d] | |
| $^{15}\mathrm{O}(8\text{–}12,\frac{3}{2}^-)$ | $3.2\pm0.1$ | $0.68\pm0.04$ | $0.30\pm0.02$ [c] | . . . |
| | | | $0.26\pm0.02$ [d] | |

[a]Statistical uncertainty only. Absolute normalization uncertainty is ±15%.
[b]Transition strength determined from beta-decay $ft$ values (Ref. 1).
[c]$B$(GT) determined from $(p,n)$ data and Eqs. (2) and (4). Note that the value of $B$(GT) for $^{13}$N(3.51 MeV) differs slightly from that in Ref. 7 because we have not averaged in values obtained at other energies.
[d]$B$(GT) determined from $(p,n)$ data and Eqs. (3) and (4).
[e]Shell-model transition strengths, Cohen-Kurath "POT" wave functions (Ref. 5).

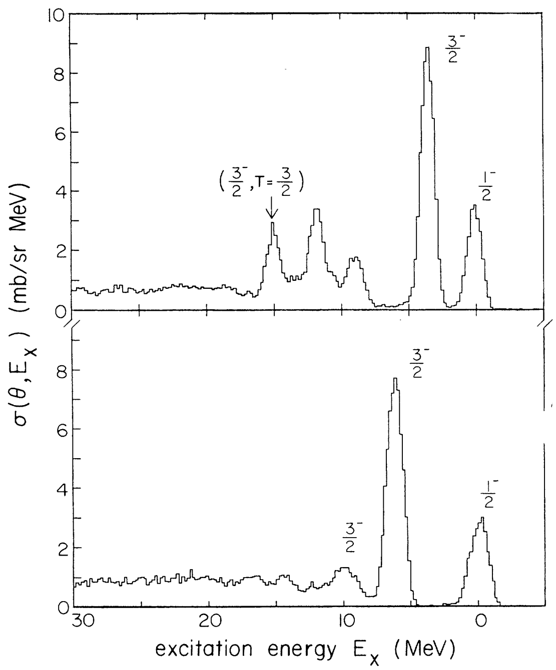


FIG. 1. Differential cross section vs excitation energy for (top) $^{13}C(p,n)^{13}N$ and (bottom) $^{15}N(p,n)^{15}O$ at $\theta=0^0$ and $E_p=160$ MeV. The $^{15}N(p,n)$ spectrum was obtained from measurements with a Melamine ($C_3H_6{}^{15}N_6$) target and has had the carbon contribution subtracted. The $^{12}C(p,n)$ contaminant transitions would appear in the $^{15}N$ spectrum at excitation energies greater than 14.6 MeV. The vertical bars illustrate the calculated GT strength distributions of Ref. 5. The dashed portions of the ground-state bars represent the Fermi contributions to the cross sections. The scale is chosen so that the ground-state bars would match the data peaks in height if the beta-decay $B$(GT) values were used instead of the calculated values. The bars for the major $\frac{3}{2}^-$ peaks are reduced by a factor of 3 to keep them on scale. The spin assignments of the unlabeled experimental peaks cannot be made on the basis of this experiment alone.

where $B(\mathrm{F})=N-Z$. Once the fraction $f_{\mathrm{GT}}=\sigma_{\mathrm{GT}}/(\sigma_{\mathrm{GT}}+\sigma_{\mathrm{F}})$ of the cross section attributable to the GT part of the mirror transition has been determined, the cross section per $B$(GT) for that target is known. Values of $B$(GT) for excited states are then extracted with this proportionality factor.

The spin-flip probability measurements give a second (independent) determination of the GT fraction in the ground-state cross sections. The spin-flip probability for these transitions is the weighted sum of the pure GT value of $S_{NN}(0°,\mathrm{GT})=0.66\pm0.03$ and the Fermi value of zero. The "pure GT" value represents an average obtained from measurements of $S_{NN}$ for many GT transitions at 160 MeV.[7] This value is consistent with the value $\frac{2}{3}$ expected for a pure $L=0$ transition. The uncertainty represents not only experimental uncertainties, but also real deviations from the nominal value that can be attributed to $L\neq0$ amplitudes in the transition.

The two procedures for extraction of the GT fraction in the ground state and $B$(GT) for excited states can be summarized in the following formulas:

$$f_{\mathrm{GT}}=[1+B(\mathrm{F})/B_M(\mathrm{GT})R^2]^{-1} \tag{2}$$

or

$$f_{\mathrm{GT}}=S_{NN}(0°,M)/S_{NN}(0°,\mathrm{GT}) \tag{3}$$

and

$$B_x(\mathrm{GT})=(\sigma_x/\sigma_M)f_{\mathrm{GT}}^{-1}B_M(\mathrm{GT})F(q), \tag{4}$$

where $R=E_p/(55\pm1\ \mathrm{MeV})$, $M$ refers to the mirror (ground-state) transition, and $x$ refers to the excited state. The factor $F(q)$ is a correction for the momentum-transfer dependence of the differenial cross sections and is constrained by $1.00\leq F(q)\leq1.10$ for the cases studied here. The values of $B$(GT) for the excited states shown in Table I indicate that the methods of Eqs. (2) and (3) are more or less consistent. If anything, the spin-flip probability procedure makes the quenching of the $\frac{3}{2}^-$ transitions look even greater.

For the strongest transitions in both masses 13 and 15, the values of $B$(GT) extracted by the above procedures are reduced from the $1p$ shell-model values by factors much larger than the typical GT quenching. A striking feature is that the model predicts a ratio of 8:1 for the $\frac{1}{2}\rightarrow\frac{3}{2}$ to $\frac{1}{2}\rightarrow\frac{1}{2}$ transitions in mass 15, and yet the observed ratio is only about 4:1. Similarly, in mass 13, a shell-model calculation[5] predicts a ratio of 7.5:1 for the ratio of the strongest $\frac{1}{2}\rightarrow\frac{3}{2}$ transition to the mirror state $\frac{1}{2}\rightarrow\frac{1}{2}$ transition, while the observed value is about 4:1. Stated in other words, it appears that the major $\frac{1}{2}\rightarrow\frac{3}{2}$ transitions are significantly more quenched than the $\frac{1}{2}\rightarrow\frac{1}{2}$ transitions.

Some fragmentation of the $\frac{3}{2}^-$ hole strength for mass 15 has been observed in pickup reactions.[8] In Fig. 1, the "peak" centered at about 10-MeV excitation in $^{15}O$ can be plausibly associated with the $T_z=-\frac{1}{2}$ component of this remaining strength. This strength in the 8–12-MeV region should be included for the comparison to the model value of $B(\mathrm{GT})=\frac{8}{3}$ if one assumes that the $p_{3/2}$ strength is spread into nearby $\frac{3}{2}^-$ levels that belong to more complicated configurations such as three-hole–two-particle and five-hole–four-particle states.

We also note that there are a number of positive-parity levels that, if excited in the $(p,n)$ reaction would not be resolved from the prominent GT peaks in the present data. Transitions to these levels involve angular momentum transfers of $L\geq1$ and should be weak at 0°. Corrections for these unresolved transitions would decrease the values of $B$(GT) deduced for the

$\frac{3}{2}^-$ transitions. In that sense the $\frac{3}{2}^-$ transition strengths for $^{15}$N in Table I may be regarded as upper limits.

Since the data span the complete energy region of the $p$ shell, the model failure points to the necessity of an enlargement of the model space beyond the $p$ shell. The nature of the discrepancy also suggests that the failure cannot be parametrized simply by the introduction of an effective axial-vector coupling constant. Two suggestions in the literature for enlarging the shell-model space to account for missing GT strength involve coupling with high-lying excitations. The suggestions involve, on the one hand, coupling to delta-particle–nucleon-hole states and, on the other hand, coupling to high-lying nucleon particle-hole states that are ignored in the usual truncations used in the shell-model calculations.[9] Towner and Khanna[10] present calculations of several types of corrections that can be applied to the simple shell model of mass 15. These corrections have the right qualitative behavior with respect to the data, but they are smaller than the observed discrepancies.

The surprise in the data is not that a simple $p$-shell model is imperfect for a description of the nuclei discussed. It is rather that the simple model that works so well for energy-level structure fails so badly for describing GT strength distributions in nuclei that are expected to be rather simple. The model failure here is very similar to that seen in $\frac{3}{2}^- \rightarrow \frac{1}{2}^-$ $M1$ transition rates for masses 13–15, where for six transitions the measured strengths are about (50–60)% of the strengths calculated wih the Cohen-Kurath interaction.[11] The $p$-shell model also yields a poor description of the transverse form factors measured in backward electron scattering in masses 13 (Hicks *et al.*[12]) and 15 (Singhal *et al.*[13]). It is also reported that the $p_{1/2}$ and the $p_{3/2}$ hole strengths as seen in $^{16}$O$(e,e'p)$ are slightly more than 50% of the shell-model, single-particle values.[14]

In summary, we have measured the cross sections and transverse spin-flip probabilities for $^{13}$C$(p,n)$ and $^{15}$N$(p,n)$ at $E_p = 160$ MeV. The $(p,n)$ data show clearly and simply the comparison of GT transition strengths to different members of a spin-orbit pair. The data suggest that the $\frac{1}{2} \rightarrow \frac{3}{2}$ GT transitions are more quenched than the $\frac{1}{2} \rightarrow \frac{1}{2}$ transitions when compared to simple shell-model calculations restricted to the $p$ shell. The $(p,n)$ data provide one more piece of convincing evidence that a $p$-shell model is inadequate for some nuclei generally considered to be $p$-shell nuclei. The total evidence is now compelling that the model space must be enlarged perhaps even much beyond that shell. The magnitude of the required corrections seems to be large enough to negate some of the appealing simplicity of the nuclear shell model in providing a guide to a valid truncation of the model space. Even for one of the simplest shell-model nuclei, the simple version of the model seems to fail.

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[1]S. Raman, C. A. Houser, T. A. Walkiewicz, and I. S. Towner, At. Data Nucl. Data Tables **21**, 567 (1978).

[2]M. Rho, Nucl. Phys. **A231**, 493 (1974).

[3]C. D. Goodman, C. A. Goulding, M. B. Greenfield, J. Rapaport, D. E. Bainum, C. C. Foster, W. G. Love, and F. Petrovich, Phys. Rev. Lett. **44**, 1755 (1980).

[4]C. D. Goodman and S. D. Bloom, in *Spin Excitations in Nuclei,* edited by F. Petrovich, G. E. Brown, G. T. Garvey, C. D. Goodman, R. A. Lindgren, and W. G. Love (Plenum, New York, 1984), p. 143.

[5]T.-S. H. Lee and D. Kurath, Phys. Rev. C **21**, 293 (1980); D. Kurath, private communication; S. Cohen and D. Kurath, Nucl. Phys. **73**, 1 (1965).

[6]T. N. Taddeucci, J. Rapaport, D. E. Bainum, C. D. Goodman, C. C. Foster, C. Gaarde, J. Larsen, C. A. Goulding, D. J. Horen, T. Masterson, and E. Sugarbaker, Phys. Rev. C **25**, 1094 (1982); T. N. Taddeucci, in *The Interaction Between Medium Energy Nucleons in Nuclei—1982,* edited by H. O. Meyer, AIP Conference Proceedings No. 97 (American Institute of Physics, New York, 1983), p. 228.

[7]T. N. Taddeucci, T. A. Carey, C. Gaarde, J. Larsen, C. D. Goodman, D. J. Horen, T. Masterson, J. Rapaport, T. P. Welch, and E. Sugarbaker, Phys. Rev. Lett. **52**, 1960 (1984).

[8]V. Bechtold, L. Friedrich, P. Doll, K. T. Knoplfe, G. Mairle, and G. J. Wagner, Phys. Lett. **72B**, 169 (1977).

[9]See papers by M. Rho and by A. Arima, in *Spin Excitations in Nuclei,* edited by F. Petrovich, G. E. Brown, G. T. Garvey, C. D. Goodman, R. A. Lindgren, and W. G. Love (Plenum, New York, 1984), pp. 7, 111, and 249, for discussions of these points of view.

[10]I. S. Towner and F. C. Khanna, Nucl. Phys. **A399**, 334 (1983).

[11]D. J. Millener and E. K. Warburton, prviate communication.

[12]R. S. Hicks, J. Dubach, R. A. Lindgren, B. Parker, and G. A. Peterson, Phys. Rev. C **26**, 339 (1982).

[13]R. P. Singhal, J. Dubach, R. S. Hicks, R. A. Lindgren, B. Parker, and G. A. Peterson, Phys. Rev. C **28**, 513 (1983).

[14]M. Bernheim, A. Bussiere, J. Mougey, D. Royer, D. Tarnowski, S. Turck-Chieze, S. Frullani, S. Boffi, C. Giusti, F. D. Pacati, G. P. Capitani, E. De Sanctis, and G. J. Wagner, Nucl. Phys. **A375**, 381 (1982).