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ISOBARIC SPIN IN STRIPPING

AND PICK-UP REACTIONS**

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IN NUCLEAR SCIENCE ABSTRACTS

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To introduce our discussion of isobaric spin in particle-transfer reactions, let us compare the reactions (d,t) and (d,He^3) performed with the same target. The first reaction removes a neutron from the target, the second removes a proton. Let us first neglect any effects of the Coulomb interaction and the neutron-proton mass difference. Figure 1 symbolizes the changes that take place in the isobaric spin and its z-component. T_A in Figure 1 is the isobaric spin of the target.

The (d,t) reaction excites states with isobaric spin $T_A - 1/2$ ($\equiv T_<$) and $T_A + 1/2$ ($\equiv T_>$). The isobaric analogue of each of the $T_>$ states should be excited by the (d,He^3) reaction. To compare the relative intensities with which pairs of analogue states are excited, we consider an expansion of the target wave function in terms of states of the first $A-1$ nucleus, coupled to the last nucleon:

* Invited paper at Conference on Isobaric Spin in Nuclear Physics, Tallahassee, Florida, March 17-19, 1966.

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$$\begin{aligned} \bar{\Psi}(A; I_A M_A, T_A T_A) = & \sum_{k, n \ell_J} (\beta_k I_k T_k; n \ell_J | \} I_A T_A) \left[\bar{\Psi}(A-1; \beta_k I_k, T_k T_k) \phi_A(n \ell_J; \frac{1}{2} \frac{1}{2}) \right]_{M_A}^{I_A} \\ & + \sum_{i, n \ell_J} (\alpha_i I_i T_i; n \ell_J | \} I_A T_A) \times \\ & (T_i \frac{1}{2} T_i - \frac{1}{2} | T_A T_A) \left[\bar{\Psi}(A-1; \alpha_i I_i, T_i T_i) \phi_A(n \ell_J; \frac{1}{2} - \frac{1}{2}) \right]_{M_A}^{I_A} \\ & + (T_i \frac{1}{2} T_i - \frac{1}{2} | T_A T_A) \left[\bar{\Psi}(A-1; \alpha_i I_i, T_i T_i - 1) \phi_A(n \ell_J; \frac{1}{2} \frac{1}{2}) \right]_{M_A}^{I_A} \end{aligned} \quad (1)$$

The square brackets indicate vector coupling of the angular momentum of the A^{th} nucleon to that of the first $A-1$ nucleons.

The symbol $(; | \}$ is a kind of fractional parentage coefficient,

and α_i symbolizes an index that distinguishes between daughter

states of the same angular momentum and isobaric spin. The

product $\left[\bar{\Psi}(A-1; \beta_k I_k, T_k T_k) \phi_A(n \ell_J; \frac{1}{2} \frac{1}{2}) \right]_{M_A}^{I_A}$

automatically has a total isobaric spin of $T_k + \frac{1}{2} (\equiv T_A)$.

However the products $\left[\bar{\Psi}(A-1; \alpha_i I_i, T_i T_i) \phi_A(n \ell_J; \frac{1}{2} - \frac{1}{2}) \right]_{M_A}^{I_A}$

and

$$\left[\bar{\Psi}(A-1; \alpha_i I_i, T_i T_i - 1) \phi_A(n \ell_J; \frac{1}{2} \frac{1}{2}) \right]_{M_A}^{I_A}$$

have to enter the sum in (1) with amplitudes given by the

vector-coupling coefficients $(T_i \frac{1}{2} T_i - \frac{1}{2} | T_A T_A)$ and

$(T_i \frac{1}{2} T_i - \frac{1}{2} | T_A T_A)$ in order to achieve a total isobaric spin of T_A .

From the wave function in (1) we immediately get the ratio of the spectroscopic factors for exciting the two analogues

$$\bar{\Psi}(A-1; \alpha_i I_i, T_i T_i) \quad \text{and} \quad \bar{\Psi}(A-1; \alpha_i I_i, T_i T_i - 1).$$

It is

$$\frac{S(I_A, T_A T_A \xrightarrow[n\ell_j]{(d,t)} \alpha_i I_i, T_> T_>-1)}{S(I_A, T_A T_A \xrightarrow[n\ell_j]{(d, He^3)} \alpha_i I_i, T_> T_>)} = \frac{(T_> \frac{1}{2} T_>-1 \frac{1}{2} | T_A T_A)^2}{(T_> \frac{1}{2} T_> -\frac{1}{2} | T_A T_A)^2} = \frac{1}{2T_A + 1}. \quad (2)$$

Now let $Z_{n\ell_j}$ and $N_{n\ell_j}$ be the numbers of $n\ell_j$ protons and neutrons in the target. Then

$$Z_{n\ell_j} = \sum_i S(I_A, T_A T_A \xrightarrow[n\ell_j]{d, He^3} \alpha_i I_i, T_> T_>) \quad (3)$$

so that

$$\sum_i S(I_A, T_A T_A \xrightarrow[n\ell_j]{d,t} \alpha_i I_i, T_> T_>-1) = \frac{Z_{n\ell_j}}{2T_A + 1}, \quad (4)$$

a sum rule for the total (d,t) spectroscopic factors leading

to daughter states with isobaric spin $T_>$. Further

$$N_{n\ell_j} = \sum_i S(I_A, T_A T_A \xrightarrow[n\ell_j]{d,t} \alpha_i I_i, T_> T_>-1) + \sum_k S(I_A, T_A T_A \xrightarrow[n\ell_j]{d,t} \beta_k I_k, T_< T_<) \quad (5)$$

which yields the corresponding sum rule for daughter states

with isobaric spin $T_<$:

$$\sum_k S(I_A, T_A T_A \xrightarrow[n\ell_j]{d,t} \beta_k I_k, T_< T_<) = N_{n\ell_j} - \frac{Z_{n\ell_j}}{2T_A + 1}. \quad (6)$$

Equations (4) and (6) are examples of a type of sum rule introduced by French and Macfarlane in 1961.

An attempt to check these sum rules was made by Sherr, Rost, Rickey, Hoot and myself. The (p,d) reaction was performed on a series of targets from T_i^{48} up to N_i^{62} , and an attempt was made to determine the sums (4) and (6) for the pick-up of an $f_{7/2}$ neutron. But to extract a spectroscopic factor from an ex-

perimental angular distribution one requires a D.W.B.A. calculation for the pick-up of an $f_{7/2}$ neutron from the well generated by all the other nucleons. This is essentially a single-particle calculation. The spectroscopic factor is the factor by which this single-particle transfer cross-section must be multiplied to yield the experimental cross-section. But to do the D.W.B.A. calculation of the pick-up of an $f_{7/2}$ neutron, one needs its radial wave function. The standard procedure for obtaining it is to assume it is the wave function in a Wood-Saxon well, of a depth chosen to give the picked-up neutron the correct separation energy. This will ensure that the asymptotic tail of the picked-up neutron has the right shape, which is important for what is believed to be a surface reaction. This means that the $f_{7/2}$ neutron radial wave function used for transitions to highly excited states will be squeezed in more tightly than the $f_{7/2}$ neutron radial wave function used for transitions to the low-lying states.

The summed spectroscopic factors so obtained are illustrated in Figure 2. The ratios S_{exp}/S_{theory} plotted there should ideally be unity. However, we would not be disturbed by a value different from unity, provided that it were constant. It is seen that there is a generally downward trend with increasing A , but what is more serious from the present viewpoint is the discrepancy between the ratios for the $T_>$ and $T_<$ summed spectroscopic factors. The $T_>$ spectroscopic factors are too large compared

with the $T_{<}$ ones, and this discrepancy is larger when there is a greater difference in excitation energy between the $T_{>}$ and $T_{<}$ states, that is when there is a greater difference between the tails of the picked-up neutron wave functions. We would decrease the $T_{>}$ spectroscopic factors relative to the $T_{<}$ ones if we would ignore the effect of binding energy on the radial wave functions, since this would allow the radial wave functions corresponding to excited states to extend further out, which would give a larger D.W.B.A. cross-section so that a smaller spectroscopic factor would be needed. The result of such an analysis is shown in the bottom half of Figure 2, where the same radial wave function is used in the D.W.B.A. calculation of low- and high-lying states. There is still a systematic decrease with increasing A, but at least the $T_{>}$ and $T_{<}$ spectroscopic factors are consistent with each other.

However, the situation cannot be allowed to rest there because the neutron separation energies for the low-lying and high-lying final states are different, and this difference should reflect itself in the radial wave functions. Moreover, there are other data which favor wave functions appropriate to the separation energy. Table I shows the analysis by Blair and Armstrong³ of their (He^3, d) data, where they added a proton to Fe^{54} , Ni^{58} and Ni^{60} to make $T_{>}$ and $T_{<}$ states, and compared this to the addition of a (d, p) neutron to make the

analogous $T_{\frac{1}{2}}$ states. Consider the case $Fe^{54}(d,p)Fe^{55}$ versus $Fe^{54}(He^3,d)Co^{55}$. States in Fe^{55} at 0, .417 and .935 MeV are seen in the (d,p) reaction, with angular momenta $3/2^-$, $1/2^-$ and $5/2^-$. The analogous Co^{55} states at 4.71, 5.14 and 5.72 MeV are seen in the (He^3,d) reaction. If we use the French-Macfarlane rule, given by equation (2), the neutron spectroscopic factors measured in the (d,p) reaction would imply proton spectroscopic factors of .27, .20, .20. If the proton spectroscopic factors are extracted from the (He^3,d) data using the separation energy procedure, values of .22, .26 and .27 are obtained. However if the radial wave functions are used that explain the (He^3,d) data to the low-lying Co^{55} states, one gets .27, .24 and .43, indicating that it is better to use the radial wave function appropriate to the correct separation energy. It should be noted that apart from the 4.71 MeV state in Co^{55} , all the $Co^{55} T_{\frac{1}{2}}$ states are unbound by a few hundred keV to a few MeV. Thus it is not surprising that we see greater sensitivity to the radial wave function than in the (p,d) experiment of Sherr et al, where the final states were all bound except for the case of $Ti^{50}(p,d)Ti^{49}$. Also since the He^3 is strongly absorbed by the nucleus, the (He^3,d) reaction may be more sensitive to the radial wave function at the nuclear surface than is the (p,d) reaction.

A way to understand the apparent insensitivity, observed by Sherr et al, of the neutron wave function to the separation energy was proposed by Pinkston and Satchler⁴. They noted that a neutron moving relative to an (A-1)-particle core with T_{ζ} sees a different effective potential than one moving relative to an (A-1)-particle core with T_{γ} . This is a consequence of the isobaric spin dependence of the nuclear interactions, and of the fact that the neutron must be antisymmetrized with respect to the neutrons in the core⁵. Lane has shown how this effect may be incorporated into the single-particle picture by the addition of a term proportional to $\vec{t}_A \cdot \vec{T}_C$ to the isobaric-spin-independent shell-model potential U_0 :

$$V = U_0 + U_1 \vec{t}_A \cdot \vec{T}_C \quad (7)$$

Here $\vec{t}_A$ refers to the last nucleon and $\vec{T}_C$ to the (A-1)-particle core. If this potential is used in the analysis of the (p,n) reaction to analogue states as a quasi-elastic process, the best fit is obtained if U_1 is positive and is peaked on the nuclear surface. As applied to the (p,d) reaction where T_{γ} and T_{ζ} refer to the isobaric spin of the (A-1)-particle core, (7) becomes

$$V = U_0 + \frac{1}{2} U_1 \times \begin{cases} T_A - \frac{1}{2} & \text{for } T_A(p,d) T_{\zeta} \\ -T_A - \frac{1}{2} & \text{for } T_A(p,d) T_{\gamma} \end{cases} \quad (7')$$

This has the consequence that the T_{γ} potential effectively has

a larger radius than the $T_<$ potential, and this counteracts the tendency of the higher $T_>$ separation energy to pull the radial wave function in. For the (He^3, d) reaction the "core" is the target, and $T_>$ and $T_<$ refer to the final state of target-plus-a nucleon:

$$V = U_0 + \frac{1}{2} U_1 \times \begin{cases} -T_c - 1 & \text{for } T_c(\text{He}^3, d) T_< \\ T_c & \text{for } T_c(\text{He}^3, d) T_> \end{cases} \quad (7'')$$

and here the well is smaller for the $T_>$ state so that the effect of the higher binding energy is reinforced by the change in the size of the well. Thus the suggestion of Pinkston and Satchler is able to reconcile the apparently conflicting results of Sherr et al and Blair and Armstrong.

However, the addition of a $\vec{t} \cdot \vec{T}_c$ term gives rise to more than a difference between the $T_>$ and $T_<$ effective single-particle potentials. It also converts the $T_>$ situation into a 2-channel problem, since the $t_+ T_-$ component in the potential can change the $(T_>, T_> + \text{proton})$ wave function into a $(T_>, T_> - 1 + \text{neutron})$ wave function. For example, let us return to the comparison of (d, t) and (d, He^3) reactions to analogue states, and consider in particular Cr^{52} as a target. If we write the Cr^{52} ground state wave function schematically as

$$\Psi(\text{Cr}^{52}) = \Psi(v^{51}) \phi(\frac{1}{2}, -\frac{1}{2}) + \Psi(\text{Cr}^{51*}) \phi(\frac{1}{2}, \frac{1}{2}), \quad (8)$$

the Schrodinger wave function takes the form given by Lane⁵
and by Robson,⁶

$$\begin{aligned} \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + U_c - \frac{5}{4} U_1 - E \right] \phi\left(\frac{1}{2}, -\frac{1}{2}\right) + \frac{\sqrt{5}}{2} U_1 \phi\left(\frac{1}{2}, \frac{1}{2}\right) &= 0 \\ \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + \Delta_c + \frac{3}{4} U_1 - E \right] \phi\left(\frac{1}{2}, \frac{1}{2}\right) + \frac{\sqrt{5}}{2} U_1 \phi\left(\frac{1}{2}, -\frac{1}{2}\right) &= 0 \end{aligned} \quad (9)$$

Here U_c is the coulomb energy term added to the proton shell-model potential, whereas Δ_c is the coulomb energy difference between V^{51} state and its isobaric analogue in Cr^{51} (see Figure 3).

$$\Delta_c = M(Cr^{51*}) + M_n - M(V^{51}) - M_p \approx 8.1 \text{ MeV} \quad (10)$$

Tamura,⁷ and Bondorff, Jägar and Lütken have solved coupled⁸ equations such as these for an energy such that the proton kinetic energy is positive but the neutron kinetic energy is negative. In our case we want both proton and neutron kinetic energies to be negative. If U_c were constant and equal to Δ_c , one would uncouple the equations (9) by taking linear combinations with the vector coupling coefficients that give Cr^{52} an isobaric spin of 2,

$$\begin{aligned} \chi_2 &\equiv \sqrt{\frac{5}{6}} \phi\left(\frac{1}{2}, -\frac{1}{2}\right) - \sqrt{\frac{1}{6}} \phi\left(\frac{1}{2}, \frac{1}{2}\right) \\ \chi_3 &\equiv \sqrt{\frac{1}{6}} \phi\left(\frac{1}{2}, -\frac{1}{2}\right) + \sqrt{\frac{5}{6}} \phi\left(\frac{1}{2}, \frac{1}{2}\right) \end{aligned} \quad (11)$$

$$\begin{aligned} \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 - \frac{7}{4} U_1 + \frac{5U_c + \Delta_c}{6} - E \right] \chi_2 + \frac{\sqrt{5}}{6} (U_c - \Delta_c) \chi_3 &= 0 \\ \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + \frac{5}{4} U_1 + \frac{U_c + 5\Delta_c}{6} - E \right] \chi_3 + \frac{\sqrt{5}}{6} (U_c - \Delta_c) \chi_2 &= 0, \end{aligned} \quad (12)$$

and this would lead directly to the French-Macfarlane ratio for the (d,t) and (d,He³) spectroscopic factors. To see this we note that if the coupling terms in (12) were absent, we could find a solution corresponding to the Cr⁵² ground state for which $\chi_2 \neq 0$ but $\chi_3 = 0$. Since $\chi_3 = 0$, $\phi(\frac{1}{2}, -\frac{1}{2}) = -\sqrt{5} \phi(\frac{1}{2}, \frac{1}{2})$ so that the (d,He³) spectroscopic factor is 5 times the (d,t) spectroscopic factor, as given by (2).

However Δ_c and U_c are not equal and an exact solution requires numerical integration of the coupled equations (9) or (12). Thus, to do a complete D.W.B.A. analysis of the hypothetical (d,t) - (d,He³) experiment, we must obtain the $T_>$ form factors by integrating the coupled equations. The $T_<$ form factor is obtained by integrating the single Schrodinger equation

$$\left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + \frac{3}{4} U_1 - \delta - E \right] \chi_< = 0, \quad (13)$$

where δ is the difference between neutron and proton separation energies in Cr⁵², as shown in Figure 3.

Taro Tamura is planning to do this calculation within the next few weeks. In a contribution to this conference, Bock, von Brentano, Duhm and Stock give some interesting speculations

about the effect of this coupling on the (He^3, α) form factor, and apply their conclusions successfully to an analysis of their $\text{Cr}^{52}(\text{He}^3, \alpha)\text{Cr}^{51}$ and $\text{Cr}^{50}(\text{He}^3, \alpha)\text{Cr}^{49}$ data.

If we take the view that the relevant nuclear symmetry in these medium-heavy nuclei is the analogue spin discussed by Dr. MacDonald¹⁰, the equations (12) automatically decouple. Let us suppose that the normalized functions $\Psi(1/2, \pm 1/2)$ satisfy

$$\begin{aligned} \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + U_c - \frac{7}{4} U_1 - E \right] \Psi(\frac{1}{2}, -\frac{1}{2}) &= 0 \\ \left[-\frac{\hbar^2 \nabla^2}{2m} + U_0 + \Delta_c - \frac{7}{4} U_1 - E \right] \Psi(\frac{1}{2}, +\frac{1}{2}) &= 0 \end{aligned} \quad (14)$$

with the same eigenvalue E . Following MacDonald, we define analogue spin operators $\vec{w}$ by

$$\begin{aligned} w_{\pm} \Psi(\frac{1}{2}, \mp \frac{1}{2}) &= \Psi(\frac{1}{2}, \pm \frac{1}{2}) \\ w_{\pm} \Psi(\frac{1}{2}, \pm \frac{1}{2}) &= 0 \\ w_z \Psi(\frac{1}{2}, \pm \frac{1}{2}) &= \pm \frac{1}{2} \Psi(\frac{1}{2}, \pm \frac{1}{2}). \end{aligned} \quad (15)$$

Then the function

$$-\sqrt{\frac{1}{6}} \Psi(Cr^{51}) \Psi(\frac{1}{2}, \frac{1}{2}) + \sqrt{\frac{5}{6}} \Psi(V^{51}) \Psi(\frac{1}{2}, -\frac{1}{2}) \quad (16)$$

will satisfy the Schrodinger equation, provided that the symmetry term is $U_1 \vec{w}_A \cdot \vec{w}_c$ and that

$$\begin{aligned} W_+ \Psi(c_r^{s_1*}) &= \sqrt{5} \Psi(v^{s_1}) \\ W_- \Psi(v^{s_1}) &= \sqrt{5} \Psi(c_r^{s_1*}) \end{aligned} \quad (17)$$

In dealing with the analogue spin, we absorb the difference $\Delta_c - U_c$ into the formalism so that no coupling is left over. It will be interesting to compare the results of this approach to that of the coupled equations.

Let us now turn to the relevance of isobaric spin to the transfer of two nucleons. Figure 4 is the two-nucleon counterpart of Figure 1.

The selection rules implied by Figure 4 have been used very effectively by Cerny, Detraz, Garvey, and Pehl in their study of the masses of members of $T > 1$ multiplets. As an example of their method, consider a comparison of the reactions $\text{Be}^4(p,t) \text{Be}^7$ and $\text{Be}^9(p,\text{He}^3) \text{Li}^7$. The former reaches Be^7 states with $T = 1/2$ and $3/2$, in both cases by transfer of two nucleons with $T = 1$ (and $T_z = 1$) and $S = 0$, assuming that the transferred nucleons are spacially symmetric. The $\text{Be}^9(p,\text{He}^3) \text{Li}^7$ reaction also reaches $T = 1/2$ and $3/2$ states in the daughter nucleus. The $T = 3/2$ states are again reached by $\Delta T = 1$, $\Delta S = 0$ transfer but the $T = 1/2$ states can be reached by both $\Delta T = 1$, $\Delta S = 0$ and $\Delta T = 0$, $\Delta S = 1$ transfer. Thus the tritons and He^3 particles that are emitted leaving Be^7 and Li^7 in analagous $T = 3/2$ states

should have exactly the same angular distributions, apart from phase space factors and appropriate isobaric spin vector coupling coefficients. On the other hand there is no reason to expect the triton and He^3 angular distributions to be simply related if the Be^7 and Li^7 are left in analagous $T = 1/2$ states. Thus a comparison of the (p,t) and (p,He^3) angular distributions makes it possible to identify the $T = 3/2$ states in the daughter nuclei. Figure 5 shows the near-identity of the $\text{Be}^9(p,t)\text{Be}^7$ and $\text{Be}^9(p,\text{He}^3)\text{Li}^7$ angular distributions associated with a $T = 3/2$ final state. The separation energies from Be^9 of two neutrons or a neutron and a proton are evidently close enough that one is not troubled by differences in radial form factors.

An interesting feature of Figure 4 is the possibility¹⁵ of reaching in the (p,t) reaction states whose isobaric spin exceeds T_z by 2. Table II summarizes the energies of the $T = T_z + 2$ "double analogue" states that have been seen¹⁶ until now. In every case where a $T = 2$ state has been sought in a $T_z = 0$ nucleus it has been found, with a width of less than 80 keV. No $T = T_z + 2$ states have been seen in nuclei with $T_z > 0$, because the fraction of the (p,t) strength that reaches $T = T_z + 2$ states is lower in these cases. An argument similar to that which led to (4) and (6) implies that for $(n\ell j)^2$ pick-up, this fraction is

$$\frac{\sum_i S(I_A T_A T_A \xrightarrow{(p,t)} \alpha_i I_i T_A^{+1}, T_A^{-1})}{\sum_i S(I_A T_A T_A \xrightarrow{(n,t)} \alpha_i I_i T_i, T_A^{-1})} = \frac{Z_{n_{eJ}} (Z_{n_{eJ}} - 1)}{N_{n_{eJ}} (N_{n_{eJ}} - 1)} \times \frac{1}{(T_A + 1)(2T_A + 1)}.$$

For $S^{34}(p,t) S^{36} T=2$ this is $\frac{1}{6}$, assuming that the neutrons are placed from doubly-closed shells. For $S^{36}(p,t) S^{34} T=3$ this is $\frac{1}{15}$.

Reactions involving the transfer of more than two nucleons have similar selection and intensity rules. It is clear that future research will make extensive use of their simplicity and generality to test our understanding of the reaction process, and to explore the intriguing relationship between isobaric analogue states.

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FIGURE CAPTIONS

1. Isobaric spin changes associated with one-nucleon pick-up.
2. Comparison of experimental and theoretical values for the summed $7/2^-$ spectroscopic factors. The upper part of the figure shows the results obtained by using a form factor determined for each state by the separation energy. The lower part of the figure shows the results obtained by using a form factor determined by the lowest $7/2^-$ state in each nucleus, and then kept constant for that nucleus.
3. Separation energies in the $Z = 24$, $N = 28$ system.
4. Spin and isobaric spin changes associated with two-nucleon pick-up. The ΔS and ΔT values are related by the assumed space-permutation symmetry of the transferred nucleons.
5. A comparison of the (p,t) and (p,He^3) reactions to isobaric analogue states of Be^7 and Li^7 (from reference 14).

TABLE 1

(1)	(2)	(3)	(4)	(5)	(6)	(7)	
State excited by (d,p) reaction (MeV)	J^π	Analogue state excited by (He ³ , d) reaction (MeV)	(C ² S) _{pred}	(C ² S) ₁	(C ² S) ₂	Estimated shell-model position (MeV)	
Fe ⁵⁵	0	Co ⁵⁵	4.71 ± 0.07	0.27	0.22	0.27	4.0
	0.417		5.14 ± 0.05	0.20	0.26	0.24	5.8
	0.935		5.72 ± 0.05	0.20	0.27	0.43	4.6
Ni ⁵⁹	0	Cu ⁵⁹	3.88 ± 0.05	0.17	0.18	0.37	1.5
	0.340		4.30 ± 0.05	0.32	0.33	0.72	2.2
	0.471		4.36 ± 0.05	0.20	0.22	0.29	3.3
	3.071		6.86 ± 0.10 a)				
Ni ⁶¹	0	Cu ⁶¹	6.40 ± 0.05	0.062	0.062	0.26	1.5
	0.069		6.47 ± 0.05	0.12	0.14	0.67	2.2
	0.290		6.65 ± 0.07 b)				
	2.133		8.42 ± 0.10 a)				

a) State is too far unbound to allow meaningful analysis.

b) Strongly interfering state present.

ENERGIES OF $T = 2$ STATES (MeV)

A	$T_Z = 0$	$T_Z = 1$
12	27.52	12.63
16	22.9	9.9
20	16.8	6.43
24	15.355	5.95
40	11.98	4.39
52	8.54	2.79
56	9.85	3.58

TABLE II

FIG. 1

FIGURE 2

FIG. 3

FIG. 4

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FIGURE 5

7/8/9/10/11/12/1/2/3/4/5/6
AM
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C.O.O.

ACTION	TO	REMARKS