ues of the electro-optic coefficient. Such large values may, in part, account for the structure due to excitons in electroreflectance spectra.[15]

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ISOSPIN-NONCONSERVING DECAYS IN MIRROR NUCLEI*

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We have measured branching ratios for the isospin-nonconserving particle decays of the 15.1-MeV $T=\frac{3}{2}$ states in ^{13}N and ^{13}C to the $(0.0, 0^+)$ and $(4.44, 2^+)$ states of ^{12}C. By comparing ratios of reduced widths with the 0^+ and 2^+ states of ^{12}C for the mirror decays, we show that the form of the isospin impurity in the 15.1-MeV states depends on T_z.

One of the more intriguing ideas in nuclear physics is the suggestion that the isospin impurities observed in nuclear wave functions may be useful in understanding the charge dependence of the internucleon force itself.[1] That the exploitation of this suggestion has not been entirely successful is largely due to the scarcity of sufficiently detailed experimental information on the size and form of these impurities. Our present knowledge of the magnitude of isospin admixtures is primarily obtained from three sources: (1) β-decay matrix elements,[2] (2) electromagnetic selection rules (especially for $\Delta T=0$, $E1$ transitions in self-conjugate nuclei),[3,4] (3) comparison of isospin-allowed and -forbidden reaction rates.[5-7]

We have employed a new technique for studying the form of isospin impurities in nuclear wave functions by comparing isospin-nonconserving decays from analog levels in mirror nuclei. In this Letter we present results for the decays of the lowest $T=\frac{3}{2}$ levels in ^{13}C and ^{13}N. Similar experiments are presently under way in mass 9 and mass 17.

Decay branching ratios (BR's) to the ground and 4.44-MeV states of ^{12}C were measured using the reactions ^{11}B(^{3}He,n)^{13}N(15.07)$\rightarrow p+{}^{12}$C and ^{11}B(^{3}He,p)^{13}C(15.11)$\rightarrow n+{}^{12}$C. Coincidences between neutrons and protons were stored in a 64 $\times$64-channel array, one variable being neutron flight time, the other proton energy. Protons were detected in a 1000-μ surface-barrier counter, neutrons in a 5.08-cm thick×12.70-cm diam heavily shielded plastic scintillator. The "start" pulse for the time-to-amplitude converter was taken from the neutron detector, the "stop" pulse from the fast current pulse of the solid-state counter. The neutron detection efficiency was determined to an accuracy of 10% using the known cross sections for the reaction D(d,n).

Since the lowest $T=\frac{3}{2}$ states in $A=13$ have spin $\frac{3}{2}$, the decays are not isotropic. To simplify the integration over solid angle in computing BR's, the protons or neutrons populating the $T=\frac{3}{2}$ levels were detected at 0°. This served to align the residual nuclei and simplified the angular correlations. In this case the correlation is axially symmetric around the beam direction and the θ dependence is given by $W(\theta)=a_0+a_2P_2(\cos\theta)$.

In Fig. 1(a) we display a singles pulsed-beam time spectrum from the reaction ^{11}B(^{3}He,n). The stop pulses in this case were obtained from the beam pulsing system. From such spectra we found the number of neutrons populating the ^{13}N($T=\frac{3}{2}$) state per monitor count [the d_0 group from the reaction ^{11}B(^{3}He,d) was used to normalize the singles and coincidence spectra]. Then the beam-pulsing system was turned off and a coincidence spectrum was accumulated with identical geometry. In Figs. 1(b) and 1(c) we show the projections onto the t_n (neutron time-of-flight) axis of the kinematically allowed regions in the E_p, t_n plane for which the outgoing protons leave the ^{12}C residual nucleus in its ground and first-excited states. The numbers of $T=\frac{3}{2}$ decays per monitor count to ^{12}C(0.0; 0$^+$) and ^{12}C(4.44; 2$^+$) were taken from the projected spectra, and the BR's computed from the ratio of decays to populations. To ensure that the "$T=\frac{3}{2}$ decays" were indeed due to a resonance in the ^{12}C$+p$ channel and not in ^{12}C$+n$, we have covered a wide range of decay angles in both forward and backward hemispheres and avoided beam energies and decay angles where ambiguities from known states in ^{13}C may arise.

The technique used to measure the ^{13}C decays was quite similar, except that the singles spectra were taken with a continuous beam. Since the proton energy resolution was degraded by the foil used to "range out" the incident beam, we used a magnetic spectrometer with an energy resolution of 35 keV to obtain the number of protons populating the ^{13}C(15.11) state per monitor count, in a separate experiment.

In Fig. 2 we present the measured differential

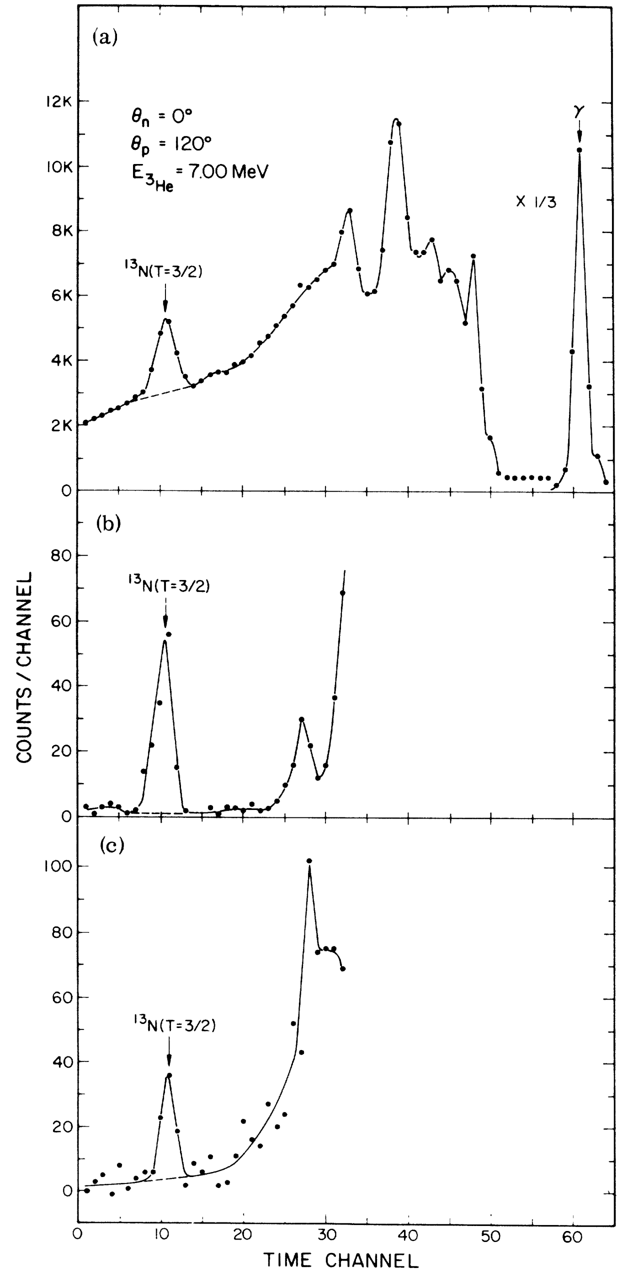

FIG. 1. Neutron time spectra from the reaction ^{11}B(^{3}He,n)^{12}C. (a) Pulsed-beam singles time spectrum. The $T=\frac{3}{2}$ state at 15.11-MeV excitation occurs in time channel 11, corresponding to a neutron energy of 1.1 MeV. Peaks near the center of the spectrum correspond to lower lying ^{13}N states. (b) Neutron spectrum in coincidence with protons from decays to the ground state of ^{12}C. (c) Neutron spectrum in coincidence with protons from decays to the 4.44-MeV state of ^{12}C.

BR's (BR's per unit solid angle) for ^{13}N and ^{13}C plotted versus $\cos^2\theta_{\text{c.m.}}$. On such a plot the angular correlations should be straight lines. As a

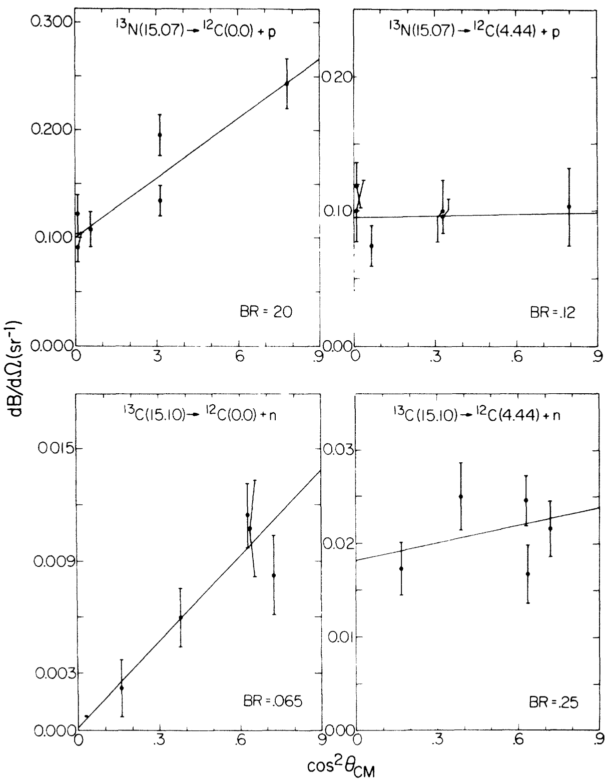

FIG. 2. Differential branching ratio measurements for decays of the 15.1-MeV $T=\frac{3}{2}$ states in ^{13}N and ^{13}C to the ground and 4.44-MeV states of ^{12}C. Branching ratios were obtained by integrating over $\theta_{c.m.}$.

check on our experimental technique, we have measured two BR's which should be close to unity and obtained 1.05 ± 0.11 and 0.98 ± 0.05 for the $^{13}C(6.87) \to {}^{12}C(0.0) + n$ and $^{16}F(3.58) \to {}^{15}O(0.0) + p$ decays, respectively. Our BR's for the $T=\frac{3}{2}$ states are listed in Table I, along with values for Γ, Γ_{γ_0}, Γ_{p_0}, Γ_{n_0}, Γ_{p_1}, and Γ_{n_1} obtained by combining our results with those of recent resonance and coincidence studies.[8-11]

The most striking feature of our results are the marked differences between the mirror transitions in ^{13}N and ^{13}C. To interpret the observed asymmetry we must first remove the trivial charge-dependent effects due to differences in phase space and barrier penetration. We presumably account for such effects by dividing the branching ratios by the $l=1$ penetrabilities ($r = 3.4$ F was used) to obtain the ratios of reduced widths,

$$\theta^2(0^+)/\theta^2(2^+) = [\mathrm{BR}(0^+)/P(0^+)] \times [\mathrm{BR}(2^+)/P(2^+)]^{-1}.$$

However, the large asymmetry is still present in the ratio of reduced widths, also listed in Table I.

We shall use first order perturbation theory to derive an expression for the ratios of reduced widths. We expect that the isospin-nonconserving interaction can reasonably be treated as a perturbation since presumably the $T=\frac{1}{2}$ admixtures are small [$\Gamma = 1.17 \pm 0.21$ keV (Refs. 8 and 10) in ^{13}N; $\Gamma = 6.2 \pm 1.1$ keV (Refs. 10 and 11) in ^{13}C] and the $T=\frac{3}{2}$ components of these states have very similar wave functions. This can be tested by comparing the $M1$ radiation widths of the $15.1 \to 0.0$ transitions in ^{13}N and ^{13}C. If the 15.1- and 0.0-MeV states were eigenstates of isospin, these $\Delta T = 1$ transitions would have identical widths (see Ref. 4). Recent electron scattering studies[12] give an $M1$ width in ^{13}C of 25 ± 7 eV. Using our value for Γ_{p_0}/Γ and measurements of $\Gamma_{p_0}\Gamma_{\gamma_0}/\Gamma$ and the $E2/M1$ amplitude ratio from the reaction[9] $^{12}C(p,\gamma_0)$, an $M1$ width of 27 ± 5 eV is found for the state in ^{13}N. The excellent agreement supports our assumption that isospin impurities in the 15.1-MeV states are small. To our knowledge this is the first case where the equality of $\Delta T = 1$ γ-radiation widths has been tested.

In first-order perturbation theory the initial (mass 13, $T=\frac{3}{2}$) and final (^{12}C + nucleon) states may be described in an obvious notation as

$$|i\rangle = |a, T=\tfrac{3}{2}, T_z\rangle + \textstyle\sum_i \beta_i(T_z)|b_i, T=\tfrac{1}{2}, T_z\rangle,$$

$$|f\rangle = \{|c, T=0, 0\rangle + \textstyle\sum_j \delta_j |d_j, T=1, 0\rangle\}|T=\tfrac{1}{2}, T_z\rangle,$$

Table I. Some properties of the decays of the 15.1-MeV $T=\frac{3}{2}$ states in ^{13}C and ^{13}N. See text for references to other work.

	BR(0^+)	BR(2^+)	$\theta^2(0^+)/\theta^2(2^+)$[a]	Γ (keV)	Γ_{γ_0} (eV)	Γ_{p_0} or Γ_{n_0} (keV)	Γ_{p_1} or Γ_{n_1} (keV)
$^{13}N(15.07)$	0.202 ± 0.020	0.121 ± 0.015	1.18 ± 0.11	1.17 ± 0.21[b]	27 ± 5	0.24 ± 0.05	0.14 ± 0.03
$^{13}C(15.11)$	0.065 ± 0.014	0.250 ± 0.036	0.17 ± 0.03	6.2 ± 1.1[c]	25 ± 7	0.40 ± 0.11	1.55 ± 0.35

[a]Error in ratio of θ^2's is less than that obtained by combining BR's because of systematic effects which cancel in ratio of θ^2's.

[b]Obtained by combining results from Refs. 8-10, and present work.

[c]Obtained by combining results from Refs. 10-12.

where the sums run over all possible admixed states. In the case where $T_z = \pm\frac{1}{2}$,

$$\langle f|i\rangle = \sum_i \beta_i(T_z)\langle c, T=\tfrac{1}{2}|b_i, T=\tfrac{1}{2}\rangle + (\tfrac{2}{3})^{1/2}\sum_j \delta_j^*\langle d_j, T=\tfrac{3}{2}|a, T=\tfrac{3}{2}\rangle \pm O(\beta\delta).$$

We will assume that there are no fortuitous cancellations in first order and neglect higher order terms.

The T_z dependence of the β_i can be made explicit:

$$\beta_i(T_z) = (E_a - E_{b_i})^{-1}\langle b_i, T_z|H_{NC}|a, T_z\rangle.$$

We then expand the isospin-nonconserving operator in spherical tensor operators, $H_{NC} = T_0^0 + T_0^1 + T_0^2$, and then use the Wigner-Eckart theorem to yield

$$\beta_i(T_z = \pm\tfrac{1}{2}) = (E_a - E_{b_i})^{-1} \times [(\tfrac{2}{3})^{1/2}\langle b_i\|T^1\|a\rangle \pm (\tfrac{2}{5})^{1/2}\langle b_i\|T^2\|a\rangle].$$

The ratio of reduced widths to final states c and c' becomes

$$\frac{\theta_c^{\,2}}{\theta_{c'}^{\,2}} = \left|\frac{\langle f|i\rangle}{\langle f'|i\rangle}\right|^2 = \left|\frac{\sum_i \beta_i(T_z)\langle c|b_i\rangle + (\frac{2}{3})^{1/2}\sum_j \delta_j^*\langle d_j|a\rangle}{\sum_i \beta_i(T_z)\langle c'|b_i\rangle + (\frac{2}{3})^{1/2}\sum_j \delta_j^{*\prime}\langle d_j'|a\rangle}\right|^2.$$

We can see immediately that a difference in the ratio of reduced widths for $T_z = \pm\frac{1}{2}$ implies the following:

(1) Isospin impurities must be important in the mass-13 $T=\frac{3}{2}$ levels; $T=1$ impurities in the ^{12}C states alone cannot produce the observed effect.

(2) If impurities in the ^{12}C wave functions can be neglected, the impurity in the $T=\frac{3}{2}$ states must arise from the admixture of more than one state. [It is very likely that $\delta_i \ll \beta_i$ since $(E_c - E_{dj}) \gg (E_a - E_{bi})$.][12]

(3) If we further assume that the energy denominators $(E_a - E_{bi})$ are independent of T_z, the isovector and isotensor matrix elements must have comparable strength. The validity of the assumption that the energy denominators do not differ significantly for $T_z = \pm\frac{1}{2}$ is difficult to assess, as the $T=\frac{1}{2}, J^\pi = \frac{3}{2}^-$ states in this energy region have not been located experimentally and shell-model calculations[13] predict a $T=\frac{1}{2}, J^\pi=\frac{3}{2}^-$ state very close to the lowest $T=\frac{3}{2}$ level.

It would, of course, be interesting to see if Coulomb mixing alone can account for the measured BR's. Exploratory calculations by Maqueda[14] using spherical shell-model wave functions and excitation energies do not reproduce our BR's. Reliable mixing calculations, of course, cannot be made until more of the $T=\frac{1}{2}, J=\frac{3}{2}^-$ states in the $A=13$ nuclei have been identified. It appears that, at present, the most promising avenue for additional work lies in careful experimental and theoretical studies of these $T=\frac{1}{2}$ levels.

It is interesting to note that the form of isospin impurities in $T=2$ states of $T_z=0$ nuclei may be studied by comparing neutron and proton decays of these levels with analogous $T=\frac{1}{2}$ states. If impurities in the $T=\frac{1}{2}$ states can be neglected, a difference in reduced widths for neutron and proton decays could be interpreted as arising from an isospin admixture containing both $T=0$ and $T=1$ components.

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NEUTRON PARTICLE-HOLE STATES OBSERVED BY INELASTIC PROTON SCATTERING FROM ^{136}Xe†

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Neutron particle-hole states have been observed in the decay of isobaric analog resonances formed in the reaction ^{136}Xe plus proton. Analysis of the on-resonance inelastic data has given inelastic proton partial widths, spectroscopic factors, spins, and parities of several of the observed particle-hole states.

We have recently completed an extensive series of measurements of elastic and inelastic proton scattering from ^{136}Xe in the bombarding energy region 9.77 to 12.98 MeV. Analysis of the elastic-scattering data[1] shows that well-defined isobaric analog resonances occur at bombarding energies of 10.270 ($f_{7/2}$), 10.874 ($p_{3/2}$), 11.255 ($p_{1/2}$), 11.583 ($f_{5/2}$), and 11.810 MeV ($f_{5/2}$) and that the corresponding compound states in ^{137}Cs are analogs of states in ^{137}Xe formed by addition of a neutron to the closed $N=82$ neutron shell in the configurations indicated above.

Virtually all the inelastic excitation functions show resonant structure. Many have an even more distinctive behavior in that they exhibit sharp peaks at some, but not all, of the analog resonances described above. These peaks in the excitation functions are symmetric and typically have a height which at angles greater than 90° is 10 to 20 times higher than the nearby off-resonance background. In common with other workers who have studied similar nuclei,[2] we identify these inelastic transitions as being due to the excitation of neutron particle-hole states in ^{136}Xe. The motivation for this Letter is to show that it is possible not only to make an identification as to the nature of these states but also to use the observed data in conjunction with a simple theory to extract information about their spins and parities and to some extent about their configurations.

The neutron particle-hole states which we wish to discuss are expected to have configurations in which the particle is in a level above $N=82$ ($2f_{7/2}$, $3p_{3/2}$, $3p_{1/2}$, $2f_{5/2}$), and the hole is in a level below $N=82$ ($2d_{3/2}$, $3s_{1/2}$, $1h_{11/2}$, $2d_{5/2}$). The possibility of $h_{11/2}$ or $d_{5/2}$ hole states will be ignored in the present analysis because of the high angular momentum involved in a $h_{11/2}$ transition and because these two levels lie further from the Fermi surface in ^{136}Xe. Particle-hole states will therefore be considered to be populated by the inelastic emission of a $d_{3/2}$ or $s_{1/2}$ proton from the analog state formed in the reaction.

The theoretical expression used in our analysis[3] is derived from the R-matrix theory as developed by Lane and Thomas.[4] We assume that the single-level approximation is valid and further that the background matrix R_0L_0 is zero. This second assumption is equivalent to the assumption that on resonance, the direct contribution to the inelastic cross section may be neglected in comparison to the compound contribution. This assumption is only partially justified since there is a small but clearly present off-resonance yield in the inelastic excitation functions. We justify this assumption on the basis that our fits at backward angles are reasonably good and that, to our knowledge, no theory which properly takes account of the direct contribution is avail-