mately 16% of the ground-state cross section. Even this result, however, would still imply considerable configuration mixing in the ground-state wave function. Furthermore, a careful examination of the spectrograph data indicates that the doublet members at this energy would necessarily lie within 10 keV of each other.

A full account of the present experiment will be given in a forthcoming paper. We wish to acknowledge helpful discussions with R. Broglia and to thank B. Bayman for the use of his DW code.

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ENERGY DEPENDENCE OF THE ISOSPIN-NONCONSERVING REACTIONS $^{12}C(d, \alpha)^{10}B(0^+, T = 1)$ AND $^{16}O(d, \alpha)^{14}N(0^+, T = 1)$†

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The resonancelike excitation function and strong forward peaking of the isospin-nonconserving $^{12}C(d,\alpha)^{10}B(0^+, 1)$ and $^{16}O(d,\alpha)^{14}N(0^+, 1)$ cross sections can be reconciled by an appropriate choice of reaction mechanism interpreted in the strong-absorption model.

The isospin-nonconserving reactions $^{12}C(d, \alpha)^{10}B(1.74, 0^+, 1)$ and $^{16}O(d, \alpha)^{14}N(2.31, 0^+, 1)$ recently have been found to proceed via some direct mechanism at incident deuteron energies above 11 MeV.[1-3] The major features of the data are the following: (1) The angular distributions are strongly forward-peaked (at about 20°) and vanish at 0° [Fig. 1(a)]. (2) The peak differential cross sections are large in magnitude, about 100 μb/sr, or 1% of typical T-allowed cross sections. (3) The excitation function for the reaction $^{12}C(d, \alpha)^{10}B^*$ is strongly energy dependent, exhibiting two large resonancelike maxima at E_d(lab) = 12.8 and 14.5 MeV, with widths on the order of 1 MeV [Fig. 1(b)]; similar behavior has been reported for the ^{16}O reaction.[3]

The comparatively large magnitude of the T-forbidden cross sections rules out mechanisms based on single-photon exchange, of which several have been proposed.[4-6] Since direct reactions involve relatively simple, low-order matrix elements of the interactions, large T-nonconserving cross sections suggest large isospin impurities in certain nuclear levels. There is thus need to reexamine previous estimates of isospin mixing in low-lying levels of light nuclei, which have mostly been based on first-order perturbation treatment of the electromagnetic interaction (i.e., on single-photon exchange).[7,8]

The strong forward peaking of the angular distribution indicates the presence of many partial waves in the amplitude and so precludes the compound nucleus mechanism. The data thus present a dilemma: how to reconcile an excitation function like that of a compound nucleus in the region of giant resonances with an angular distribution which unequivocally indicates a direct mechanism.

What can we learn from the general features of these reactions? Since the entrance- and exit-channel spins are respectively 1 and 0, first-order direct two-nucleon pickup is forbidden by spin and parity conservation, as well as by isobaric symmetry[1,9]; so we must consider mechanisms of second or higher order in perturbation theory. (The vanishing of the angular distribu-

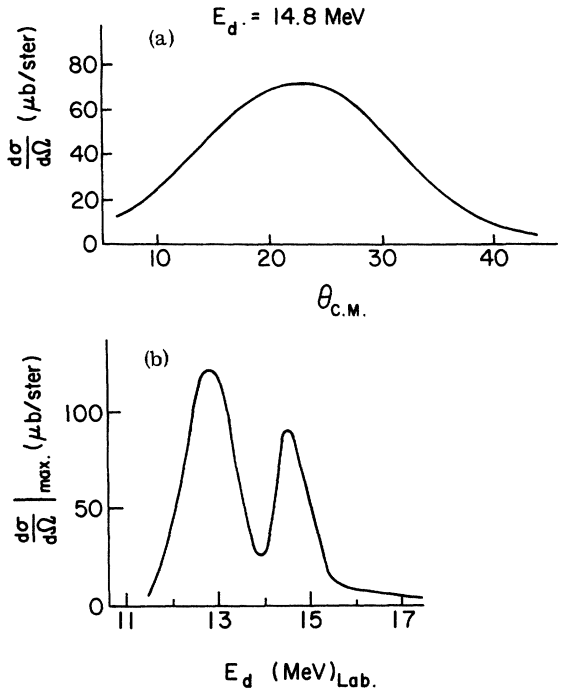

FIG. 1. (a) Typical angular distribution in $^{12}C(d, \alpha)^{10}B(0^+, 1)$ (Ref. 2). (b) Excitation function for the same reaction (Ref. 2).

tion at $\theta = 0$ is another consequence of the channel spins.) Only those deuterons polarized parallel or antiparallel to the beam direction contribute. Also, the dependence of the reaction matrix on azimuthal angle and on deuteron polarization is unobserved; so we can characterize the reaction by the amplitude $A(E, \theta)$ with the partial-wave decomposition[10]

$$A(E, \theta) = \sum_{j=1}^{\infty} (2j+1) A_j(E) d_{10}^{\ j}(\theta), \tag{1}$$

where E and θ are the barycentric energy and scattering angle, and $d_{10}^{\ j}(\theta)$ is the Wigner d-function.[10] The observed angular distributions can be understood with hardly any reference to the specific reaction mechanism. (By reaction mechanism we mean "lowest order perturbation-theoretic amplitude.") Since the incident and outgoing particles are strongly absorbed at these energies, we expect the partial-wave amplitude $A_j(E)$ to be small for $j \lesssim KR$, where R is some sort of strong-absorption radius. Conversely, for $j > KR$, the effects of entrance- and exit-channel distortion become negligible; so we can represent $A_j(E)$ reasonably well by its Born approximation $B_j(E)$. In a large class of models, the $B_j(E)$'s have the same sign and decrease with j approximately as

$$|B_j(E)| \sim \exp[-(j+\tfrac{1}{2})\Lambda(E)]. \tag{2}$$

This behavior, together with the well-known small-angle, large-j formula[11]

$$d_{10}^{\ j}(\theta) \sim J_1((j+\tfrac{1}{2})\theta),$$

enables us to write, for $\theta \ll \Lambda(E)$,

$$A(E, \theta) \simeq J_1(KR\theta) \sum_{j=KR-\frac{1}{2}}^{\infty} (2j+1) B_j(E). \tag{3}$$

Equation (3) is valid up to the first zero of the Bessel function J_1; at larger angles in this approximation scheme $A(E, \theta)$ is set to zero. Choosing $R = 4.5$ fm (reasonable for this mass region) we find that the forward peak in the ^{12}C data is well represented by Eq. (3). Furthermore, the forward displacement with energy of the first maximum in the angular distribution is reproduced by (3) with R fixed at 4.5 fm. The angular distributions for the ^{16}O reaction are not yet available, but we would be surprised to find them qualitatively different. We emphasize that Eq. (3) was derived under general assumptions and therefore conclude that the angular distribution contains little information about the mechanism. All the physical information resides in the magnitude and energy dependence of the differential cross sections, summarized by the factor appearing in Eq. (3),

$$F(E) = \sum_{j=KR-\frac{1}{2}}^{\infty} (2j+1) B_j(E). \tag{4}$$

We can now discuss the reaction mechanism. It must be at least second order (but probably cannot be third or higher order) in perturbation theory, and it must transfer two nucleons, not conserve isospin, and reproduce the observed energy dependence. The Coulomb-induced preferential spin-flip model, described in several previous papers[12] and illustrated in Fig. 2(a), was proposed as a theory of the results of Meyer-Schützmeister, von Ehrenstein, and Allas.[1] Unfortunately, the energy-dependence of this model is nothing like that reported recently by Jänecke et al.,[2,3] and therefore it must be regarded as experimentally disproved (although it may yet apply to the high-energy tail of the excitation function[2]).

The mechanism we now propose has several features. Its amplitude is represented by the "box" diagram in Fig. 2(b). The incident deuter-

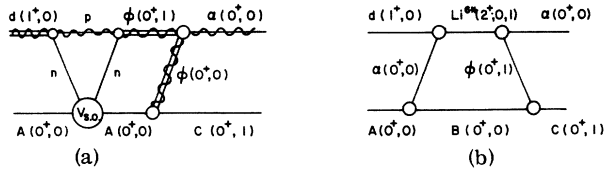

FIG. 2. (a) Diagrammatic representation of the preferential spin-flip process (Ref. 10). The superposed wavy lines indicate particles propagating in Coulomb-modified plane waves. (b) Diagram for the mechanism introduced in this paper.

on picks up an α particle from the target A forming ^{6}Li in an (isospin-mixed) excited state, plus the intermediate residual nucleus B. The ^{6}Li* then propagates and decays through its $T=1$ channel into the outgoing α particle (which is detected) plus a $^1S_0$ neutron-proton pair φ which is captured by B to form the $T=1$ residual nucleus C. There are two reasons for choosing an intermediate state containing ^{6}Li*. First, ^{6}Li has a 2^+ $T=0$, $T=1$ doublet of states at 4.6- and 5.4-MeV excitation[13] which may be strongly isospin mixed.[14] Second, the two prominent bumps in the excitation function (on both ^{12}C and ^{16}O) occur 1 to 2 MeV above the respective thresholds for producing the ^{6}Li states in the reaction

$$A+d\rightarrow B+{}^6\mathrm{Li}^*.$$

This correspondence is no accident. Finally, we see that the process illustrated in Fig. 2(b) dominates all other second-order processes because the ^{6}Li resonances are picked out by the $\Delta T=1$ process under consideration.

The next question is how does the mechanism described above account for the observed energy dependence? Roughly speaking, when the average energy in the intermediate state permits the d-α system to resonate (on-shell), the spatial extension characterizing the amplitude shown in Fig. 2(b) is much greater than when, on the average, no pair of particles resonates. Now the incident and outgoing particles with low total angular momentum (or, equivalently, at small impact parameters) are absorbed almost completely. The energy dependence arises from the fact that at some energies, the interaction mechanism extends outside the absorption volume, whereas at other energies it does not. The degree of extension determines the overlap with incident and outgoing states, thereby determining the magnitude of the amplitude. The key phrase in the preceding qualitative discussion is "average energy in the intermediate state." We can give it a quantitative meaning by noting that the function $F(E)$ defined by (4), which contains the energy dependence of the amplitude, can also be written as an integral:

$$F(E)=\int_0^\infty dE''\,\rho(E'',E)[\gamma_1(E''+E_1-E-i\gamma_1)^{-1}-\gamma_2(E''+E_2-E-i\gamma_2)^{-1}]. \quad (5)$$

The factor in square brackets, containing the difference of two Breit-Wigner terms, is an approximate representation of the isospin-nonconserving part of the d-α reaction matrix in the vicinity of the 2^+ resonances in ^{6}Li. E_1 and E_2 are the positions of the ^{6}Li resonances referred to an appropriate zero of energy, and the γ's are half their respective widths. Equation (5) was obtained by interchanging the order of summation and integration in (4). [The partial-wave amplitude $B_j(E)$ corresponding to the box diagram, Fig. 2(b), involves an integration over the kinetic energy of the intermediate state.] All the difficulties have therefore been subsumed into $\rho(E'', E)$, a strongly peaked function of E'' which involves the cutoff radius R, and the overlap integrals (single-particle momentum-space wave functions) for $A\rightarrow\alpha+B$ and $C\rightarrow\varphi+B$. The recipe for the construction of ρ from these ingredients will be given elsewhere.

Figure 3 illustrates qualitatively how the observed double peak in the excitation function is predicted by Eq. (5). As the incident beam energy increases, the Breit-Wigner peaks move from the left through the ρ peak, which moves more slowly. When both resonances are under the ρ peak, there is almost complete cancellation (the relative minus sign between the Breit-Wigner terms is quite general and results from isospin

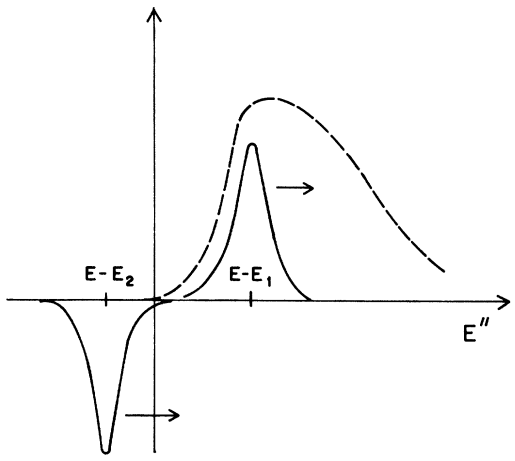

FIG. 3. Behavior of the integrand of Eq. (5). The dashed line represents $\rho(E'',E)$ and the solid lines, the two-resonance Breit-Wigner factor.

nonconservation), which explains the sharpness of the dip between the bumps in Fig. 1(b). The observed spacing and widths of the maxima of $|F(E)|^2$ are not identical with those of the ^{6}Li resonances and are somewhat different also for the ^{12}C and ^{16}O reactions. Part of these rather small discrepancies results from kinematics, and the rest can easily be explained in terms of the behavior of $\rho(E'',E)$, i.e., in terms of the A and C overlap integrals.

A reliable theoretical estimate of the absolute magnitude of the cross sections is beyond the scope of this note and is probably beyond our present competence in nuclear reaction theory. One can, however, assume some simple forms for the various wave functions and calculate $F(E)$ as defined in the strong-absorption model by Eq. (4) in order to get a rough idea of the magnitude of the cross section predicted by the mechanism shown in Fig. 2(b). I have done this and was able to reproduce the observed peak value of 100 μb/sr, assuming 5% isospin mixing in the ^{6}Li states and a slightly reduced value of R (4.1 fm); this indicates that the theory is in principle able to explain the observations.

In summary, we have been led, fairly inevitably, to a new model of direct isospin-nonconserving reactions, obtaining qualitative agreement with the data. It is hoped that similar methods will find application in the study of other classes of second-order direct transitions.

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SPINS AND ORBITAL ANGULAR MOMENTUM CONFIGURATIONS OF LEVELS IN ^{90}Y

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Polarization and differential cross-section measurements have been made for ^{89}Y(p, p)^{89}Y in the isobaric analog regions of the 2^- and 1^- states at 2.48 and 2.63 MeV in ^{90}Y. An optical-model–plus–resonance analysis of the data indicates that these states can be described by a $d_{3/2}$ state coupled to the $\frac{1}{2}^-$ target spin. The admixture of $d_{5/2}$ and $s_{1/2}$ configurations in the 2^- and 1^- states, respectively, is found to be small. The spin sequence $2^-, 1^-$ with the higher spin lower in energy is strongly favored.

A number of spin and parity assignments for isobaric analog resonances observed in elastic proton scattering have been made.[1] In all these experiments the target nucleus had spin zero. The l values for the resonances could be determined from differential cross-section measurements, and the values of the total spin could then be obtained from polarization measurements.[1-3] The analysis of proton resonances produced in elastic scattering from targets with nonzero spin