Letters to the Editor

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The Decay of Ta^{183}†

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IN a previous paper[1] the presence of a γ line decaying with a half-life of the order of 10 days was reported in connection with the decay of 112-day Ta^{112}. Using the curved crystal spectrometer, a photographic γ-ray spectrometer, and the β-ray spectrometer, we have investigated a strong source produced by irradiation of Ta in the Arco (Idaho) reactor.[2] Our results show a number of γ lines associated with a 5.2 ± 0.1 day half-life. This half-life can be ascribed to the decay of Ta^{183}, formed by two successive neutron captures in stable Ta^{181}. Mihelich[3] recently gave a value for the thermal neutron cross section of Ta^{182}, using the intensity ratio of two lines, one in Ta^{182}, the other in Ta^{183}. Our preliminary estimation is based on the ratio of β-spectra intensities of Ta^{182} and Ta^{183}, and gives a value

$$\sigma_c \gtrsim 10\,000 \text{ barns}$$

for this cross section. Table I lists the energy, wavelength, and relative intensity obtained for 21 lines. The associated β^- spectrum has an end point at 558 ± 10 kev, but we have no evidence that it is simple. In spite of the fact that several energy combinations are possible among the γ-lines ($A+B=D$, $O+I=U$, $R+E=U$, $Q+F=U$, $N+E=T$, $D+M=S$, $F+B=L$, $B+E=J\cdots$) we do not believe that a coherent decay scheme can be constructed using these data only. Further investigations are planned.

TABLE I. γ rays emitted by a Ta^{183} source.

Line	Wavelength (in A)	Energy (kev)	γ-ray relative intensity (internal conversion *not* taken into account)
A	266.726 ± 0.025	46.481 ± 0.005	···
B	235.729 ± 0.022	52.593 ± 0.005	···
C	···	$83\ \pm1$	Weak
D	125.141 ± 0.013	99.070 ± 0.010	1.4
E	114.865 ± 0.013	107.933 ± 0.012	2.4
F	112.985 ± 0.013	109.729 ± 0.013	0.4
G	112.303 ± 0.150	110.395 ± 0.150	0.1
H	87.134 ± 0.025	143.773 ± 0.050	0.4
I	86.005 ± 0.025	144.151 ± 0.050	2.7
J	77.229 ± 0.015	160.532 ± 0.030	1.1
K	76.834 ± 0.015	161.458 ± 0.025	3.3
L	76.371 ± 0.015	162.335 ± 0.030	1.8
M	64.339 ± 0.022	192.692 ± 0.065	0.6
N	60.458 ± 0.015	205.062 ± 0.050	1.8
O	59.086 ± 0.100	209.823 ± 0.035	1.9
P	···	$236\ \pm1$	Weak
Q	50.757 ± 0.010	244.259 ± 0.040	3.0
R	50.386 ± 0.006	246.051 ± 0.030	10
S	42.501 ± 0.010	291.708 ± 0.060	2.5
T	39.599 ± 0.010	313.078 ± 0.065	2.6
U	35.018 ± 0.010	354.038 ± 0.090	3.9

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[2] We gratefully acknowledge the efforts of Dr. W. B. Lewis in arranging the irradiation in the material-testing reactor.
[3] J. W. Mihelich, Phys. Rev. **91**, 427 (1953).

Magnetic Inelastic Scattering of Slow Neutrons*

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IN studies of crystal structure by neutron and by x-ray diffraction, the intensities of the Bragg maxima are reduced by thermal diffuse scattering. This reduction is given, both for neutrons and x-rays, by the well-known Debye-Waller factor. The thermal diffuse scattering differs somewhat for neutrons as compared with x-rays, however, because the large mass of the former implies a substantial energy change of the diffusely scattered neutrons. The energy changes of the neutrons, which need not be considered in diffraction work, are of great importance when extremely slow neutrons are scattered. For incident neutrons in the wavelength range 5–20A, the thermal or inelastic scattering consists almost entirely of energy gain by the neutrons. Under these conditions, the inelastic scattering cross section is proportional to λ; hence it is easily separated experimentally from elastic scattering, which is independent of λ. In addition, the coherent part of the elastic scattering disappears for wavelengths longer than the crystal cut off (twice the longest spacing of lattice planes).

Although measurement of the actual energy gain of the scattered neutrons would give the most direct information on the lattice vibrations in the scattering crystal, results so far in the long-wavelength region have consisted mainly of total cross sections. The total cross sections can be calculated from the Debye spectrum of lattice vibrations but the calculation is exceedingly complicated for any but the simplest crystals. If it is assumed, however, that the inelastically scattered waves from different atoms do not interfere (the incoherent approximation), the calculations can readily be carried out.[1,2] The calculated inelastic scattering using the incoherent approximation is in reasonable agreement with experiments[3,4] for beryllium, aluminum, and copper, for example. Although quantitative discrepancies occur for lead and bismuth, the qualitative variation of the cross section with temperature and wavelength is correct.

The scattering of iron exhibits a behavior that is not in agreement with the theoretical calculations, however. Early experiments[5,6] at room temperature indicated that the inelastic scattering of long-wavelength neutrons is slightly in excess of the theoretical value, but the excess was extremely difficult to determine because of the large capture cross section, which is also proportional to λ. Measurements of the iron cross section at long wavelengths have now been extended to sample temperatures well above the Curie point, using the Brookhaven slow chopper. At high temperature the observed inelastic scattering, identified by proportionality to λ, rises rapidly above the theoretical lattice-vibration value. The results are shown in Figs. 1 and 2 as functions of wavelength and sample temperature, compared with theory.[1,2] The inelastic scattering is obtained from the measured total cross section by subtracting the $1/v$ capture as well as a slight nuclear incoherent contribution (0.4 b).

The excess scattering in iron is certainly inelastic because of

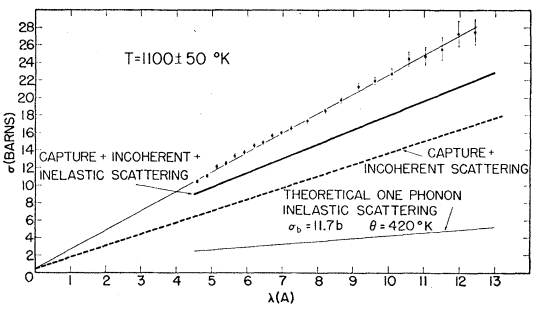

FIG. 1. The measured total cross section of iron at 1100°K as a function of wavelength.

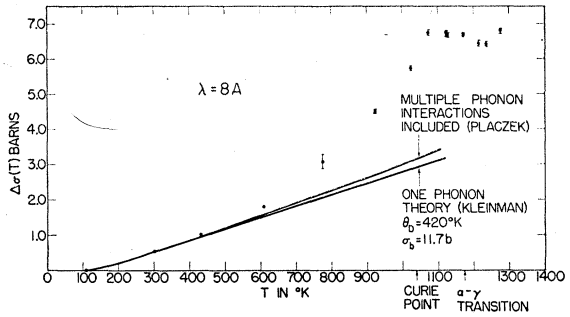

FIG. 2. Variation with temperature of the inelastic cross section of iron at 8A (cross section at T minus cross section at 110°K).

its proportionality to λ. The fact that the excess scattering shows a change in slope at the Curie point indicates that it also is related to the magnetism of iron. It thus seems certain that the additional scattering found in iron is magnetic inelastic or "spin-wave" scattering, of the type calculated by Moorhouse[7] for low temperature. Unfortunately, the calculation of magnetic inelastic scattering at low temperature, based on the spin-wave theory, is difficult to extend to the region of the Curie point. Recent calculations by Van Hove,[8] however, indicate that the shape of the experimental temperature curve near the Curie point is consistent with theory.

The scattering above the Curie point is interesting because of its relationship to the structure of iron in the paramagnetic region. For example, if iron were an ideal paramagnet above the Curie point in the sense that the unaligned spins were uncoupled, then typical elastic paramagnetic scattering would be expected. A form-factor variation with wavelength would be observed, with a leveling off of the scattering at a value of 5 barns for long wavelength (greater than 20A). As there is no evidence for this form-factor behavior, the spins are evidently coupled strongly enough so that an energy gain always occurs at scattering, a gain that is large compared to the neutron energy. The similarity of the observed scattering above and below the Curie point, and the lack of paramagnetic scattering at high temperature, is in line with a suggestion of Van Hove[8] that the magnitude of the energy change is fixed by the exchange integral, regardless of the long-range order that determines the ferromagnetic or paramagnetic nature of iron. The spatial arrangement of the magnetically active electrons in the atoms must also be closely similar both above and below the Curie point. This conclusion is reached because the atomic-form factor, which affects the inelastic scattering in a sensitive way, evidently does not change markedly at the Curie point.

We have benefited greatly in this work by discussions of scattering and magnetic theory with G. Placzek, L. Van Hove, and J. Goldman. The assistance of R. R. Smith in taking some of the data is gratefully acknowledged.

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The Sequence of γ Emission in Triple Correlations

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BIEDENHARN, Arfken, and Rose[1] have noted that angular correlation data in certain cases can fix the nuclear decay scheme uniquely. However, they also discuss an example of a triple correlation where the sequence of the last two γ's is not determined by the correlation function. This is the reaction $B^{11}(p,\gamma_1\gamma_2)C^{12}$ in which the last two γ's are emitted according to the scheme (2 D 1 D 0).

We shall prove the theorem that the same ambiguity is present in the more general γ cascade $\{j_0\ (2^{L_0}\text{ pole})\ j_a\ (2^L\text{ pole})\ L\ (2^L\text{ pole})\ 0\}$. The spins of the levels are $(j_0, j_a, L, 0)$, and the symbol 2^L denotes the multiple order of the emitted radiation.

The same conclusion holds if, in the place of the 2^{L_0} pole emission, one has to do with a particle capture followed by the emission of two gammas as in the triple gamma cascade.

There is some practical interest in the cascade of three gammas for the spins 6, 4, 2, 0. Kurath[2] has described a discrepancy in the decay scheme of Sc^{48} that is resolved only by the assumption that the product Ti^{48} decays by a triple gamma cascade. Experimental verification of the existence of the three gammas has been reported.[3] Unfortunately, the remaining uncertainty in the decay concerning the order of the last two gammas, according to the theorem to be proved, cannot be determined by angular correlation. An example wherein the ambiguity concerns the order of the first two gammas in a triple cascade is furnished by Cr^{52}. Here, in principle, angular correlation can establish the correct order of the transitions.

For the proof of the theorem we start with the general expression for the triple cascade as given by Coester and Jauch.[4] The formula given there is valid for particle or photon emission and also includes general mixture of multipoles. If we specialize to the case of pure multipole emission, we obtain the triple angular correlation function in the form:

$$W(\Omega_0\Omega_1\Omega_2)=\text{const}\sum_{\xi_1\xi_2\xi_3}\sum_{ia\,jb}\sum_{k_0k_1k_2,n_1n_2}Ck_0(\xi_0)Ck_1(\xi_1)Ck_2(\xi_2)$$
$$\times Dn_10^{k_0*}(\phi_0\theta_0)Dn_10^{k_1}(\phi_1\theta_1)Dn_20^{k_2}(\phi_2\theta_2)0)$$
$$\times(j_aj_a|\Gamma(j_0j_0L_0L_0)|0k_0)(j_aj_a|\Gamma(j_bj_bL_1L_1k_0)|k_2k_1)$$
$$\times(k_2k_1n_2n_1|k_0n_0)(j_bj_b|\Gamma(j_cj_cL_2L_2k_2)|0k_2)|Rj_a(\xi_0\xi_1\xi_2j_b)|^2. \quad (1)$$

The spins of the four nuclear levels involved are denoted by j_0, j_a, j_b, and j_c. The multipole order of the three transitions is L_0, L_1, and L_2. The Γ-coefficients are certain sums over Racah coefficients and are defined in Eqs. (B5) and (B6) of their paper. The variable ξ denotes the set of quantum numbers needed to specify the emitted radiation. Thus for photons ξ is the multipole order L and the parity.

Equation (1) may be written in the form

$$W(\Omega_0\Omega_1\Omega_2)=\sum_{k_1k_2,n_1n_2}X_{n_1n_2}{}^{k_1k_2}(\phi_0\theta_0)Dn_10^{k_1}(\phi_1\theta_1)Dn_20^{k_2}(\phi_2\theta_2)0), \quad (2)$$

from which it follows that the necessary and sufficient condition that

$$W(\Omega_0\Omega_1\Omega_2)=W(\Omega_0\Omega_2\Omega_1) \quad (3)$$

is that the coefficients $X_{n_1n_2}{}^{k_1k_2}$ satisfy the symmetry relation:

$$X_{n_1n_2}{}^{k_1k_2}=X_{n_2n_1}{}^{k_2k_1}. \quad (4)$$

This relation implies, first of all, on account of the coefficients $C_k(\xi)$ in (1), that

$$L_1=L_2. \quad (5)$$

Next we see that the last of the Γ coefficients in (1) must be independent of k_2. This coefficient is given by

$$[j_bj_b|\Gamma(j_cj_cL_2L_2k_2)|0k_2]=(-1)^{L_2+j_b-j_c-k_2}$$
$$\times(2j_b+1)(2j_c+1)^{-\frac{1}{2}}W(j_bL_2j_bL_2;j_ck_2), \quad (6)$$

and the condition that it be independent of k_2 implies $j_c=0$ and consequently $j_b=L_2=L_1$. The rest of the coefficients are then symmetrical, and the theorem stated at the beginning is proven.

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