when the sample is inside, by (a) a shift to the right of the position of the 176-kev peak as a result of its coincidence with the 27-kev x-rays; and (b) a marked increase in intensity for the peak at 622 kev due to the coincidence between the 176-kev and the 425-kev γ rays. Such indications, however, are definitely lacking here. If the good energy fitting of this decay scheme is to be preserved, one must assume that the energy level at 470 kev is metastable with a half-life longer than several microseconds. However, γ rays originating in this level have been reported to have low e/γ factors. An alternative solution is to place the 176-kev transition elsewhere. This would require a fourth β transition from Sb^{125}, which has not been reported.

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A β-Decay Matrix Element for a Deformed Core Model*

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The ft value for an allowed unfavored β transition, calculated on a deformed core+single-particle model, is ≤ 3 times the single-particle value and about 4 percent of the observed value. The deformations of initial and final states were based on their quadrupole moments. The calculation indicates that for this model core orthogonality generally does not account for the difference between allowed favored and unfavored ft values.

I. INTRODUCTION

THERE exist many nuclei whose quadrupole moments are much larger than can be expected from the shell model. For example, the quadrupole moments of two In ($Z=49$) isotopes and two Sb ($Z=51$) isotopes, which, according to that model, have single-hole and single-particle proton configurations, are approximately three times the theoretical values. It is noteworthy, however, that the signs of quadrupole moments are quite generally (as for In and Sb) those predicted by shell theory, and further, that no anomalous quadrupole moments appear for those nuclei which have both closed neutron and closed proton shells±one nucleon. In order to explain the high quadrupole moments observed in some regions, deformed core models have been introduced.[1]

Another discrepancy between shell theory and experiment appears in the ft values for allowed β transitions. The shell model calculations indicate that all allowed transitions should have ft values of about the same order of magnitude.[2] Actually, nearly all ft values for allowed transitions fall into two groups: Favored, $\log ft=2.9$ to 3.6, and unfavored, $\log ft=4.5$ to 6.0. Empirically, it appears that only transitions between states which, according to the supermultiplet theory, should belong to the same supermultiplet, are favored. Nearly all of these transitions appear for light nuclei (with mass number $A<40$). On the other hand, nearly all allowed transitions for heavier nuclei, as well as many for the light nuclei, have unfavored ft values. It is for the heavier nuclei that anomalously high quadrupole moments appear, and one might suspect that a deformed core model which accounts for these would also account for the high unfavored ft values. There is no indication at all, however, that substantial core deformations exist for light nuclei, so that the unfavored ft values which appear in that region would remain unexplained. It will be seen below that the most obvious interpretation of the deformed core model, in terms of a wave function in configuration space, fails to account for the whole difference between favored and unfavored transitions even in medium heavy nuclei. This need not be considered as a conflict between the deformed core models and experiment because the deformed core model's wave function in ordinary configuration space has not hitherto been specified closely enough to permit definite conclusions to be drawn. On the contrary, it may be hoped that the flexibility of the model is sufficient to avoid the apparent difficulty to which we are drawing attention.

In a deformed core+one particle model (for odd-A

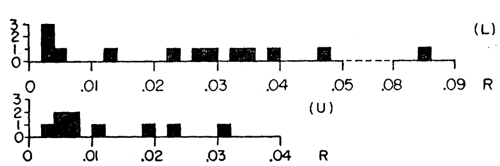

FIG. 1. Histograms of the ratio $R=[(ft)_{sp}]/[(ft)_{exp}]$ for allowed unfavored transitions of the like-core (L) and unlike-core (U) types.

* This work was supported in part by the U. S. Atomic Energy Commission and the Higgins Scientific Trust Fund.

† Hercules Fellow, 1952–1953.

[1] J. Rainwater, Phys. Rev. **79**, 432 (1950); D. L. Hill and J. A. Wheeler, Phys. Rev. **89**, 1102 (1953); A. Bohr and B. R. Mottelson, Kgl. Danske Videnskab. Selskab, Mat.-fys. Medd. **27**, No. 16 (1953).

[2] E. P. Wigner, Proceedings of the Harwell Nuclear Physics Conference (Ministry of Supply, Harwell, Berks., 1950). Also, I. Talmi, Phys. Rev. **91**, 122 (1953).

nuclei) β transitions fall into two groups:[3] The like-core transitions, for which the neutron and proton numbers, N_c, Z_c, of the initial core equal the ones of the final core, and the others, unlike-core transitions. An example of the like-core type is the following transition:[4]

$$\begin{array}{llll} \mathrm{Ge}^{75} & \beta^-\longrightarrow & \mathrm{As}^{75} & ft=160\,000 \\ N,Z=43,32 & & 42,33 & \\ J=\frac{1}{2}\ \text{(from shell theory)} & & J'=\frac{3}{2}\ \text{(measured)} & \\ & & Q=(0.3\pm0.2)\times10^{-24}\ \mathrm{cm}^2. & \end{array} \quad (1)$$

The core for both nuclei has $N_c,Z_c=42,32$. Both the ft value and Q are typical for this region. The assumption of a single particle β transition ($p_{1/2}\rightarrow p_{3/2}$) leads to $\Sigma\mathbf{G}^*\cdot\mathbf{G}=8/3$, and $(ft)_{\mathrm{sp}}=1990$. $\mathbf{G}$ is the Gamow-Teller matrix element, and the summation is over final states with different components J_z' of total angular momentum. $(ft)_{\mathrm{sp}}$ is based upon constants[5] obtained for transitions between nuclei which, according to the shell model, have single-particle and -hole configurations. The calculated ft value is too small by a factor of 80. Ge^{75} has $Q=0$ if its spin is 1/2. In the present calculation, its core will be assumed to be spherically symmetric. Spin 1/2, corresponding to a $p_{1/2}$ single-particle state, is predicted by the simple coupling rules for 43 odd nucleons. Spins 5/2 and 3/2, corresponding to $f_{5/2}$ and $p_{3/2}$, have not been observed for any nucleus in this shell with more than 37 odd nucleons.

As an illustration of an unlike-core transition, let us consider $\mathrm{Ga}^{73}(\beta^-)\mathrm{Ge}^{73}$. For the initial nucleus, N,Zs $=42,31$; for the final one, 41,32. The initial core thus has $N_c,Z_c=42,30$; the final one has 40,32. One might expect that the core factors in β-decay matrix elements (see Sec. II) would differ for the two types of cores.

All data[4] for the set of ground state to ground state transitions between odd-A nuclei in the regions where, according to the shell model, allowed transitions may appear were examined. There are 22 such transitions with definitely known ft values and $4.5<\log ft<6.0$. Two additional transitions, which may have $\Delta J=0$ or 1 and no change in parity, probably are of the l-forbidden type; they are $\mathrm{Ni}^{63}(\beta^-)\mathrm{Cu}^{63}$ and $\mathrm{Ni}^{65}(\beta^-)\mathrm{Cu}^{65}$. Their ft values both equal 3.6×10^6; the initial spins are not known. The remaining transitions in this set are of the favored type with $\log ft<4.0$.

The ratio $R=(ft)_{\mathrm{sp}}/(ft)_{\mathrm{exp}}$ of the theoretical single-particle ft value divided by the experimental one, was calculated for the 22 allowed unfavored transitions. Histograms of R of these transitions, divided into like-core (L) and unlike-core (U) groups, are given in Fig. 1. R is proportional to the square of the matrix element. It is plain that the matrix element is, on the average, larger for the L transitions; the average value of R for these is 2.3 times that for the U transitions. However, the difference seems to be much smaller than we would have expected on the basis of a deformed core model with only one particle outside the core. The ratio R was also calculated for the favored transitions in this set. This time there are ten transitions each of the L and U types. The average of R for the L transitions is 1.02 times that for the U transitions. This is consistent with the generally accepted hypothesis that the core plays at most a very slight role in the favored transitions.

It may seem surprising that there are, in fact, four transitions among the like-core ones with $0.002\leq R\leq0.006$. (See Fig. 1.) Two of these transitions are $\mathrm{O}^{19}(\beta^-)\mathrm{F}^{19}$ and $\mathrm{Si}^{31}(\beta^-)\mathrm{P}^{31}$. The single particle has been assigned an $s_{1/2}$ state for both transitions; this leads to an especially low R. Probably a several-particle configuration would describe the situation better. No reason for the low R's of the other two transitions, $\mathrm{Ca}^{45}(\beta^-)\mathrm{Sc}^{45}$ (with single-particle transition $f_{7/2}\rightarrow f_{7/2}$) and $\mathrm{Te}^{127}(\beta^-)\mathrm{I}^{127}$ (with $d_{3/2}\rightarrow d_{5/2}$), seems apparent.

II. MATRIX ELEMENTS FOR THE DEFORMED CORE MODEL

Let us calculate the matrix element for a like-core transition on the following model: $A-1$ particles are contained in a box, which is a sphere for the undeformed core, and an ellipsoid for the deformed one. The last nucleon is in a state with angular momenta j and l in a system fixed in the core (strong coupling). If x_k denotes the space and spin coordinates of the kth particle in the laboratory system, and ξ is the set $(x_1,\cdots,x_{A-1})$ and further, $x_A=x$, then the wave function is

$$\Psi_{Mn}{}^{Jj}(\xi,x)=a(J,j,n)\int\mathfrak{D}^J(R)_{nM}\psi_n{}^j(Rx)\Phi(R\xi)dR. \quad (2)$$

Φ and $\psi_n{}^j$ are core and single-particle wave functions; $\mathfrak{D}^J(R)_{nM}$ is a representation coefficient;[6] R is an element of the rotation group; a is a normalization factor. It can readily be checked that $\Psi_{Mn}{}^{Jj}(\xi,x)$ belongs to M'th line of the representation $\mathfrak{D}^J(R)$. Ψ has definite parity equal to that of ψ. The parity of Φ is even, since there are two particles to each space state in the core. The integration is taken over the parameters of all elements of the rotation group. The core state may be a superposition of states with several angular momenta, $J^{(c)}$, but it is assumed to have $J_{z_c}{}^{(c)}=0$. z_c is the z axis in the core system, and $J_{z_c}{}^{(c)}$ is the component of core angular momentum along this axis.

The matrix element of the ν component of an irreducible tensor operator T_ν^k of order k which operates only on the coordinates and spin x of the outer particle

[3] L. W. Nordheim, Report on the Indiana Conference on Nuclear Spectroscopy and the Shell Model (University of Indiana, Bloomington, 1953).

[4] Data are taken from: King, Dismuke, and Way, Oak Ridge National Laboratory Report No. 1450, 1952 (unpublished), and P. F. A. Klinkenberg, Revs. Modern Phys. **24**, 63 (1952).

[5] G. L. Trigg, Phys. Rev. **86**, 506 (1952). Also, A. Winther and O. Kofoed-Hansen, Kgl. Danske Videnskab. Selskab, Mat.-fys. Medd. **27**, No. 14 (1953).

[6] See, e.g., E. P. Wigner, *Gruppentheorie und ihre Anwendung auf die Quantenmechanik der Atomspektren* (F. Vieweg und Sohn, Braunschweig, 1931).

can readily be calculated. It is

$$\langle J'M'j'n'|T_\nu^k|JMjn\rangle = \int\int \Psi'_{M'n'}{}^{J'j'}(\xi,x)^*\, T_\nu^k\, \Psi_{Mn}{}^{Jj}(\xi,x)d\xi dx, \quad (3)$$

where the integration includes a summation over spin components. Then, with Φ and Φ' wave functions of the spherical and deformed cores,

$$\frac{\langle J'M'j'n'|T_\nu^k|JMjn\rangle}{a'(J',j',n')^*a(J,j,n)} = \int dR\int dS\mathfrak{D}^{J'}(R)_{n'M'}{}^*\mathfrak{D}^J(S)_{nM} \cdot\int \psi_{n'}{}^{j'}(Rx)^*T_\nu^k\psi_n{}^j(Sx)dx \cdot\int \Phi'(R\xi)^*\Phi(S\xi)d\xi. \quad (4)$$

$\psi_n{}^j$ belongs to the nth line of $\mathfrak{D}^j$:

$$\psi_n{}^j(Sx)=\sum_r \mathfrak{D}^j(S)_{nr}{}^*\psi_r{}^j(x). \quad (5)$$

Substitution of (5) into (4) and the transformation $R\xi=\eta$ lead to

$$\int dR\mathfrak{D}^{J'}(R)_{n'M'}{}^*\sum_{r'}\mathfrak{D}^{j'}(R)_{n'r'} \cdot\int dS\mathfrak{D}^J(S)_{nM}\sum_r \mathfrak{D}^j(S)_{nr}{}^*\cdot\int\psi_{r'}{}^{j'}(x)^*T_\nu^k\psi_r{}^j(x)dx \cdot\int\Phi'(\eta)^*\Phi(SR^{-1}\eta)d\eta. \quad (6)$$

The orthogonality relation for the representation coefficients is, with $h=\int dR$,

$$\int dR\mathfrak{D}^k(R)_{ab}{}^*\mathfrak{D}^{k'}(R)_{a'b'} = \frac{h}{2k+1}\delta(k,k')\delta(a,a')\delta(b,b'). \quad (7)$$

Φ is spherically symmetric, so that (6) and (7) lead to

$$\langle J'M'j'n'|T_\nu^k|JMjn\rangle = a'(J',j',n')^*a(J,j,n) \cdot\frac{h^2}{(2j'+1)(2j+1)}\int\psi_{M'}{}^{j'}(x)^*T_\nu^k\psi_M{}^j(x)dx \cdot\int\Phi'(\eta)^*\Phi(\eta)d\eta\cdot\delta(J',j')\delta(J,j). \quad (8)$$

The matrix element equals 0, unless $J=j$ and $J'=j'$. The first condition is obvious: Since Φ is spherically symmetric, it describes a core state with angular momentum $J^{(c)}=0$, and the nucleon's j coupled to it can give only $J=j$. The second condition follows from an expansion of Φ' in terms of core functions with definite angular momenta, $\chi[J^{(c)}]$. The core integral equals 0 unless $\chi(0)$ appears in this expansion. This will occur only if $J'=j'$.

$|a(j,j,n)|^2$ may be calculated from (4) and (8) by substitution of $J'=J=j$, $T_\nu^k=T_0^0=1$ and $\Phi'=\Phi$. Then $|a(j,j,n)|^2=(2j+1)^2h^{-2}$. The expression for $|a'(j',j',n')|^2$ is more complicated. From (4) and (6), with $V=SR^{-1}$,

$$\frac{1}{|a'(j',j',n')|^2}=\frac{h}{2j'+1}\int dV|\mathfrak{D}^{j'}(V)_{n'n'}|^2 \cdot\int\Phi'(\eta)^*\Phi'(V\eta)d\eta. \quad (9)$$

An upper limit to the integral in η is 1 for all V. Then $|a'(j',j',n')|^2\geq(2j'+1)^2h^{-2}$. It should be noted that only this lower limit, not $a'(j',j',n')$ itself, is independent of n'. Substitution in (8) gives:

$$|\langle J'M'j'n'|T_\nu^k|JMjn\rangle|^2\geq\left|\int\psi_{M'}{}^{j'}(x)^*T_\nu^k\psi_M{}^j(x)dx\right|^2 \cdot\left|\int\Phi'(\eta)^*\Phi(\eta)d\eta\right|^2. \quad (10)$$

The lower limit of the factor in the matrix element due to the core thus has absolute value $|\int\Phi'(\eta)^*\Phi(\eta)d\eta|$.

III. A CALCULATION OF THE CORE FACTOR

In order to obtain an orientation concerning the magnitude of the core matrix element, we have chosen the previously mentioned Ge^{75} transition as an example. Cube-shaped and parallelopiped-shaped boxes of equal volume were substituted for the spherical and ellipsoidal ones. It is most likely that the matrix element between the last two states does not differ substantially from the one to be calculated here. The edge

TABLE I. States for a cube-shaped box.

$n_x^2+n_y^2+n_z^2$	n_x	n_y	n_z	Number of states for one particle type (spin ½)	Cumulative number
18	1	1	4	6	46
17	2	2	3	6	40
14	1	2	3	12	34
12	2	2	2	2	22
11	1	1	3	6	20
9	1	2	2	6	14
6	1	1	2	6	8
3	1	1	1	2	2

TABLE II. $I(n,m;\xi)$ for $m\neq n$. ($\epsilon=1.0275$).

(n,m)	(1,3)	(3,1)	(2,4)	(4,2)
$I(n,m;\epsilon)$	−0.020	−0.021	−0.035	−0.038
$I(n,m;\epsilon^2)$	−0.038	−0.043	−0.066	−0.078

of the cube (state α) will be taken as $d=2\epsilon a$. The parallelopiped (state β) has a square base with edges $2a$ and height $2\epsilon^3 a$.

The wave functions are products of those for a one-dimensional square well with infinitely high walls at $x=d$ and $x=-d$:

$$u_n(x)=\frac{1}{\sqrt{d}}t\left(\frac{n\pi x}{2d}\right),\quad |x|\leq d, \tag{11}$$

$$=0,\qquad |x|>d,$$

where

$$t\left(\frac{n\pi x}{2d}\right)=\begin{cases}\cos(n\pi x/2d), & \text{if } n \text{ is odd}\\ \sin(n\pi x/2d), & \text{if } n \text{ is even.}\end{cases} \tag{12}$$

The states for the three-dimensional cube-shaped box are given in Table I. Three quantum numbers n_x, n_y, n_z determine each state, and its energy is given by $(n_x^2+n_y^2+n_z^2)\cdot\hbar^2\pi^2/8md^2$.

The deformation, ϵ, can be calculated from the contribution to Q by the core of As^{75}. This contribution will be assumed equal to 0.67×10^{-24} cm^2, a value about twice the largest observed in the region around As. The operator is

$$Q=\sum_{j=1}^{Z_c}(2z_j^2-x_j^2-y_j^2), \tag{13}$$

where the summation is over protons only. The total wave function is the antisymmetrized product of single-particle wave functions, and application of a well-known formula[7] leads to

$$\langle\beta\{n\}|Q|\beta\{n\}\rangle=\sum_{j=1}^{Z_c}\langle\beta n_j|2z_j^2-x_j^2-y_j^2|\beta n_j\rangle, \tag{14}$$

where the j's denote proton coordinates and $Z_c=32$. $n_j=(n_{jx}, n_{jy}, n_{jz}, s_{jz}, t_{j\zeta})$; it denotes the quantum numbers of a single-particle state. s_{jz} and $t_{j\zeta}$ are spin and isotopic spin components. $\{n\}$ is the set $n_1,\cdots,n_{A-1}$, with $A-1=74$. β denotes the deformed state. From the volume of the cube-shaped state,

$$8d^3=(4\pi/3)\cdot74\cdot1.4^3\times10^{-39}\text{ cm}^3, \tag{15}$$

d is calculated as 4.74×10^{-13} cm. Substitution of the assumed Q and of d into (14) leads to $\epsilon=1.0275$.

The integral of the wave functions of the two core states will be calculated next:

$$\int\Phi'(\eta)^*\Phi(\eta)d\eta\approx\langle\alpha\{n\}|\beta\{n\}\rangle=\prod_{i=1}^{A-1}\langle\alpha n_i|\beta n_i\rangle$$

$$+\text{terms containing at least one factor}$$

$$\langle\alpha n_i|\beta n_k\rangle\langle\alpha n_k|\beta n_i\rangle,\text{ with } i\neq k. \tag{16}$$

Each scalar product for one particle is a product of three integrals:

$$\langle\alpha n_i|\beta n_k\rangle=I(n_{ix},n_{kx};\epsilon)I(n_{iy},n_{ky};\epsilon)\times I(n_{kz},n_{iz};\epsilon^2)\delta(s_{iz},s_{kz})\delta(t_{i\zeta},t_{k\zeta}). \tag{17}$$

The I's are defined as follows. Let $2d$ = a particular edge for one state, and $2\xi d$ the corresponding edge for the other state. Then

$$I(n,m;\xi)=\frac{1}{d\sqrt{\xi}}\int_{-d}^{d}t\left(\frac{n\pi x}{2\xi d}\right)t\left(\frac{m\pi x}{2d}\right)dx, \tag{18}$$

which depends only upon the ratio ξ of the two edges and on n and m. From parity considerations, $I(n,m;\xi)=0$ if $n-m$ is odd, and direct integration for $n-m\neq0$ and even gives

$$I(n,m;\xi)=\frac{4}{m\pi\sqrt{\xi}}\cdot\frac{(-1)^{\nu/2}}{(n/m\xi)^2-1}t\left(\frac{n\pi}{2\xi}\right), \tag{19}$$

where $\nu=m$, if m is even, and $\nu=m+1$, if m is odd. Numerical results are given in Table II. $\xi=\epsilon$ for the x and y directions and $\xi=\epsilon^2$ for the z direction.

For $n=m$, the integral can be calculated as a series in $\delta=\xi^{-1}-1$, which is a small quantity. One can expand, with $v=n\pi x/2d$,

$$\sin\left(\frac{n\pi x}{2\xi d}\right)=\sin(1+\delta)v=\sin v+\delta v\cdot\cos v$$

$$-\frac{(\delta v)^2}{2}\sin v-\frac{(\delta v)^3}{6}\cos v+\cdots. \tag{20}$$

A similar expansion can be made for $\cos(n\pi x/2\xi d)$. Simple integrations lead for both even and odd n to

$$I(n,n;\xi)=1-\tfrac{1}{8}\delta^2(1+\tfrac{1}{3}n^2\pi^2)+\tfrac{1}{8}\delta^3-\cdots. \tag{21}$$

Table III contains numerical values of $1-I(n,n;\xi)$.

From the results of Tables I to III it follows that the first term in (16) equals 0.584, and the sum of terms containing just two integrals $I(n,m;\xi)$ with $n\neq m$ equals -0.025. The remaining terms contain more factors of these integrals and their sum is much smaller. Therefore,

$$|\langle\alpha\{n\}|\beta\{n\}\rangle|^2=(0.559)^2=0.31. \tag{22}$$

TABLE III. $[1-I(n,n;\xi)]\times10^4$. ($\epsilon=1.0275$).

$n=$	1	2	3	4
$[1-I(n,n;\epsilon)]\times10^4$	4	13	27	48
$[1-I(n,n;\epsilon^2)]\times10^4$	15	49	107	187

[7] See, for example, 6[6](9), of E. U. Condon and G. H. Shortley, *The Theory of Atomic Spectra* (Cambridge University Press, Cambridge, 1951), p. 171.

IV. CONCLUSION

From (10) and (22), by using the constants of reference 5,

$$(ft)_{\text{core+sp}} \leq 1990/0.31 = 6420.$$

This calculation yields an upper limit only ~ 3 times the single-particle ft value, despite the fact that it was based upon a quadrupole moment equal to about twice the observed one. The experimental ft value for Ge^{75} is 25 times larger than this upper limit. It is evident from the nature of the calculation that the result for other like-core nuclei in this mass region will be similar. For lighter nuclei, the core matrix element $\langle \alpha\{n\} | \beta\{n\} \rangle$ will be still larger, even though the measured ft values for unfavored transitions are not substantially lower.

Formula (8) indicates that, in fact, the matrix element equals 0 if $J' \neq j'$. A modification of the model in which at least one of the core states contains a large admixture of states with $j' \neq J'$ could therefore yield an arbitrarily small matrix element. A regularity found by Kopfermann,[8] however, seems to indicate that, for several isotope pairs at least, the single particle appears in just one state: On graphs of quadrupole moment *vs* magnetic moment, the two experimental points for odd-A isotopes and the theoretical point for the pure single-particle model are approximately collinear. Another possibility of explaining the large ft values within the framework of the deformed core model would be by assuming that both the initial and final cores are deformed. The core functions would have to be substituted in formula (6). It does not appear, however, that this would change our result drastically. It should be mentioned, finally, that the results of calculations[9] based upon the collective model in the form given by Bohr and Mottelson, a form which assumes that both cores are deformed, are of the same magnitude as the one obtained here.

An alternate explanation of the unfavored ft values for allowed transitions in terms of the original shell model would postulate instead the predominance of different configurations in the initial and final states. The fact that all the allowed transitions between nuclei with double closed shells$\pm$one nucleon fall into the favored group is consistent with this explanation. It is not clear, however, how to explain the grouping of the allowed ft values which is observed. (There are very few with $3.6 < \log ft < 4.5$.)

[8] H. Kopfermann, Naturwiss. **38**, 29 (1951).

[9] S. Suekane, Progr. Theoret. Phys. Japan **10**, 480 (1953).

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Energy Distributions of Fragments from Fission of U^{235}, U^{238}, and Pu^{239} by Fast Neutrons*

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The energies of fragments from fission of U^{235}, U^{238}, and Pu^{239} by 14-Mev neutrons, from fission of U^{238} by 2.5-Mev neutrons, and from fission of U^{235} and Pu^{239} by thermal neutrons have been measured in a single Frisch grid ionization chamber. The energy distributions for fast neutrons are similar to those previously obtained for fission by thermal neutrons. The most probable energies of the light and heavy fragments for fission by 14-Mev neutrons do not change significantly from their values for slow-neutron induced fission. The valley between peaks is higher for fission induced by 14-Mev neutrons than for low-energy neutron-induced fission.

A double Frisch grid chamber has been used to measure simultaneously the energies of both fragments from fission of U^{235} by 14-Mev neutrons. The main change in the distribution of fission modes from that for thermal-neutron-induced fission is the increased probability for symmetrical fission.

INTRODUCTION

SINCE the first work of Frisch[1] many measurements of the energy distributions of fission fragments have been made, on a number of fissionable isotopes, and at various neutron energies.[2] The purpose of the present work was to make measurements at a neutron energy of 14 Mev for comparison with measurements at other energies.

The experiments reported fall into two classes: (a) measurements of the energy of one of the fragments

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† Submitted in partial fulfillment of the requirements for the degree of Doctor of Philosophy in the Graduate College of the State University of Iowa.

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[2] See, for example, (a) J. L. Fowler and L. Rosen, Phys. Rev. **72**, 926 (1947); (b) D. C. Brunton and G. C. Hanna, Can. J. Research **A28**, 190 (1950) and Phys. Rev. **75**, 990 (1949); (c) D. C. Brunton and W. B. Thompson, Can. J. Research **A28**, 498 (1950) and Phys. Rev. **76**, 848 (1949); (d) J. Jungerman and S. C. Wright, Phys. Rev. **76**, 1112 (1949).