In all these predictions it is assumed that

$$x^2 \equiv \frac{\text{sum of squares of Gamow-Teller coupling constants}}{\text{sum of squares of Fermi coupling constants}} = (1.15)^2.$$

In the predictions giving two values of the coefficient, the first corresponds to $x<0$ and the second to $x>0$. It should be noted that in the event of failure of time-reversal invariance the Gamow-Teller—Fermi interference term in the two-component neutrino theory might be reduced to make the theory consistent with the present measurement.

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Polarized Positrons from Ga^{66} and Cl^{34} †

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Several experiments[1] have shown that positrons emitted in Gamow-Teller beta transitions are polarized in their directions of motion ($\boldsymbol{\sigma}\cdot\mathbf{p}$ positive). We find this to be true also for positrons emitted by Ga^{66}, which is probably a pure Fermi transition. Preliminary results for the pure Fermi transition of Cl^{34} also indicate the same sign for $\boldsymbol{\sigma}\cdot\mathbf{p}$.

Positrons in a selected energy band were magnetically focused on a Lucite converter (Fig. 1). High-energy photons from two-quantum annihilation are almost completely circularly polarized in the direction of the positron spin.[2] This polarization was measured by absorption in an iron analyzer[3] which can be magnetized parallel or antiparallel to the incident positron momentum. For high-energy positrons the gamma rays are emitted in a narrow forward cone. This permits the use of a low-transmission magnetic spectrometer[4] without undue loss of intensity. With 1% transmission and a 12-cm thick analyzer, about 1000 counts per minute were obtained when 3-Mev positrons from a 20-mC Ga^{66} source were focused on the converter. The pulse-height spectrum from the 3×3 inch scintillation counter was remarkably flat. Uncertainties introduced by small shifts in pulse height due to stray magnetic fields or other causes were therefore small. Extensive lead shielding was provided to reduce the background due to nuclear gamma rays. Under the experimental conditions described above, scattered gamma rays contributed about 20% of the pulses above 2 Mev. For smaller pulses the background rose sharply. Repeated checks showed that effects of the analyzer field on the counter or on the positron trajectories did not contribute any measurable systematic error. A sequence of reversing lens and analyzer currents and interchanging various circuits was designed to average over possible small instrumental asymmetries.

Figure 2 shows the results obtained for the conditions described above. The ordinate is the fractional difference between counting rates for the two analyzer field directions, after correction for background. Some data were also obtained at lower positron energies and others with a thinner analyzer. These are summarized in Table I.

Results obtained with Cl^{34} are less extensive and satisfactory than for Ga^{66}, primarily because the

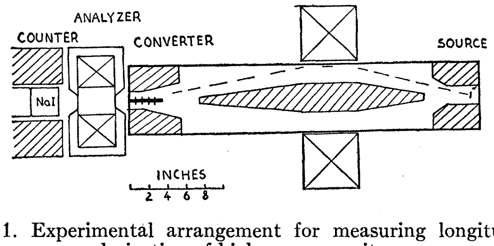

Fig. 1. Experimental arrangement for measuring longitudinal polarization of high-energy positrons.

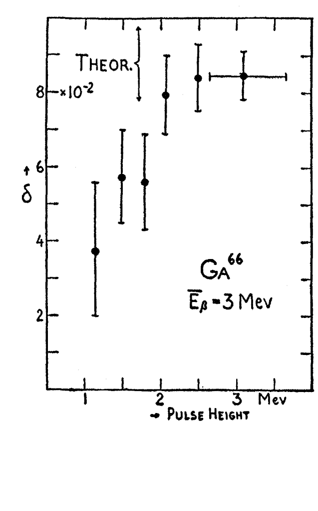

Fig. 2. Effect of field direction on transmission of annihilation radiation from Ga^{66} through analyzer.

TABLE I. Percent difference in transmission through magnetized iron for annihilation quanta above 2 Mev.

	Calc	Ga^{66} $E_\beta=3$ Mev	Ga^{66} $E_\beta=2.6$ Mev	Cl^{34} $E_\beta=3$ Mev
Thick analyzer	8.8 ± 1.0	8.4 ± 0.5	8.8 ± 1.0	5.6 ± 2.1
Thin analyzer	4.4 ± 1.5	4.9 ± 0.9	...	5.4 ± 1.3

available source strength was smaller by about a factor five. The short half-life prevented a complete cycle of check runs with each source. The gamma-ray background was relatively higher, especially for smaller pulse heights. The results are summarized in Table I, which presents the average values of δ for counter pulses above 2 Mev. The uncertainty in the theoretical value reflects our uncertainty concerning the magnetic flux distribution in the analyzer.

Although first results of a new method should be accepted with some caution, our experiment shows convincingly that the positrons from Ga^{66} are highly polarized with positive $\boldsymbol{\sigma}\cdot\mathbf{p}$. The results suggest complete polarization. They contradict a recent scattering experiment.[5] While they remove the only strong argument for a basic difference in the behavior of Gamow-Teller and Fermi transitions, no affirmative conclusions can be drawn from the Ga^{66} data until the parity of this nuclide is established with certainty. The results for Cl^{34} indicate that positrons from a pure Fermi transition are also predominantly polarized with positive $\boldsymbol{\sigma}\cdot\mathbf{p}$. The degree of polarization remains uncertain. Our data are compatible with complete polarization and make it quite improbable that the value is near zero.

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Coulomb Scattering of Polarized Electrons*

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EXPERIMENTS carried out with longitudinally polarized β beams have focused attention on the Coulomb scattering of polarized electrons. The final-state density matrix can be calculated by means of the scattering matrix M which in turn depends on two complex functions F and G (first introduced by Mott) that have been tabulated for a number of values of $\alpha=Z/137$, $\beta=v/c$, and θ (scattering angle) by Bartlett and Watson,[1] McKinley and Feshbach,[2] Sherman,[3] and others. In this note, the amount of longitudinal depolarization L and the asymmetry ratio A for a β beam are expressed exactly in terms of F and G. Approximate analytical expressions for L and A (corresponding to the second-order Born approximation) are also derived. The results are applied to a double-scattering experiment performed by de-Shalit *et al.*[4]

Denoting the unit vectors of the momenta of the initial and scattered electrons by $\mathbf{n}_1$ and $\mathbf{n}_2$ and the unit vector normal to the plane of scattering by $\mathbf{n}$, we can write the scattering matrix in the form

$$M=f-ig\mathbf{n}\cdot\boldsymbol{\sigma}=\lambda[G\sec(\tfrac{1}{2}\theta)-i\mathbf{n}\cdot\boldsymbol{\sigma}F'\csc(\tfrac{1}{2}\theta)]\times\exp(-\tfrac{1}{2}i\mathbf{n}\cdot\boldsymbol{\sigma}\theta), \quad (1)$$

where

$$F'=iq(1-\beta^2)^{\frac{1}{2}}F, \quad \text{with} \quad q=\alpha/\beta.$$

If $\boldsymbol{\zeta}_1$ and $\boldsymbol{\zeta}_2$ are the expectation values of the initial and final spin vectors of the electron in its rest frame (see Tolhoek[5] whose notations we follow closely), the final state density matrix is given by

$$\rho_2=\tfrac{1}{2}I(1+\boldsymbol{\zeta}_2\cdot\boldsymbol{\sigma})=\tfrac{1}{2}M(1+\boldsymbol{\zeta}_1\cdot\boldsymbol{\sigma})M^\dagger, \quad (2)$$

where I is the differential cross section $d\sigma/d\Omega$. In the case of an initial beam with longitudinal polarization of degree $P_1=|\boldsymbol{\zeta}_1|$, we have $\boldsymbol{\zeta}_1=P_1\mathbf{n}_1$, and, using (1) and (2),

$$\boldsymbol{\zeta}_2=S(\theta)\mathbf{n}+P_1U(\theta)\mathbf{n}_2+P_1T(\theta)\mathbf{n}\times\mathbf{n}_2, \quad (3)$$

where

$$\bar{I}(\theta)S(\theta)=-2i(F'G^*-GF'^*)\csc\theta, \quad (4a)$$

$$\bar{I}(\theta)T(\theta)=2(F'G^*+GF'^*)\csc\theta, \quad (4b)$$

$$\bar{I}(\theta)U(\theta)=|G|^2\sec^2(\tfrac{1}{2}\theta)-|F'|^2\csc^2(\tfrac{1}{2}\theta), \quad (4c)$$

$$\bar{I}(\theta)=|G|^2\sec^2(\tfrac{1}{2}\theta)+|F'|^2\csc^2(\tfrac{1}{2}\theta). \quad (4d)$$

We note the indentity $S^2+T^2+U^2=1$. For arbitrary initial polarization $\boldsymbol{\zeta}_1$, the scattered intensity is

$$I(\theta)=\mathrm{Tr}\rho_2=\lambda^2\bar{I}(\theta)[1-\boldsymbol{\zeta}_1\cdot\mathbf{n}S(\theta)]. \quad (5)$$

$S(\theta)$ is the asymmetry function tabulated by Sherman.[3] $U(\theta)$ measures the longitudinal polarization after scattering since $P_1U(\theta)=\boldsymbol{\zeta}_2\cdot\mathbf{n}_2$. The longitudinal depolarization is therefore $L=P_1(1-U)$. Since $\mathbf{n}\times\mathbf{n}_2=(\mathbf{n}_2\cos\theta-\mathbf{n}_1)\csc\theta$, we also find

$$\boldsymbol{\zeta}_2\cdot\mathbf{n}_1=P_1[U(\theta)\cos\theta-T(\theta)\sin\theta]. \quad (6)$$

for the remaining initial polarization in the direction $\mathbf{n}_1$. If the emerging beam (with polarization $\boldsymbol{\zeta}_1'=\boldsymbol{\zeta}_2$) is further scattered through an angle θ' and plane of scattering $\mathbf{n}'$, Eq. (5) gives

$$I(\theta')=\bar{I}(\theta')[1-\boldsymbol{\zeta}_1'\cdot\mathbf{n}'S(\theta')],$$