nucleus	ΔE, Mev	$\frac{d\sigma}{d\Omega}$, mbn/sterad	$\frac{d\sigma_{E2}}{d\Omega}$, mbn/sterad	σ_{E2}, mbn	nucleus	ΔL, Mev	$\frac{d\sigma}{d\Omega}$, mbn/sterad	$\frac{d\sigma_{E2}}{d\Omega}$, mbn/sterad	σ_{E2}, mbn
Li^7	0.476	35			Na^{23}	0.439	9	0.15	2
F^{19}	0.197	16	0.06	0.7	Mg^{24}	1.370	7	0.20	3.5
F^{19}	1.355	2			Al^{27}	0.843	1.5		
F^{19}	1.426	2			Al^{27}	1.013	2.5		
F^{19}	1.558	8							

of the levels in F^{19}, Na^{23}, and Mg^{24}. The calculations were made with formulas (B-32) and (B-38) of Ref. 9. The probabilities given in the table for the transitions from the ground state into excited states with energies ΔE are taken from Refs. 9 and 10.

The tabulated results indicate that when E_d = 4.5 Mev the contribution of σ_{E_2} to the experimental value of σ_{tot} (σ_{tot} = 30 to 50 mbn for (d, d′) reactions on Mg^{24}, Na^{23}, and F^{19}) amounts to several percent. One can therefore conclude that at E_d = 4 to 4.5 Mev and above the process of nuclear excitation by the Coulomb field of the incident deuterons cannot be the dominant process that leads to inealstic scattering of the deuterons. This conclusion contradicts the (d, d′) reaction theory developed by Mullin and Guth.[11]

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COMPTON SCATTERING OF CIRCULARLY POLARIZED PHOTONS BY ELECTRONS WITH ORIENTED SPIN

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THE further development of the quantum electrodynamics of longitudinally polarized electrons and photons assumes new significance in connection with the discovery of parity nonconservation. As we have already noted,[1,2] to take longitudinal polarization of electrons into account one must, in the calculation of matrix elements, use not the Casimir formula but formula 2.12 of Ref. 3, in which s stands for the eigenvalue of the operator $(\Delta\sigma)/i\sqrt{-\nabla^2}$. This operator gives double the electron spin projection onto its direction of motion.

It is of interest to observe that the above-mentioned formula can also be used if the electron is initially at rest. In that case the formula may be brought into the form

$$\langle\alpha_{n'}, \alpha_n\rangle = \frac{1}{16}\mathrm{Sp}\,\alpha_{n'}\left(1 + \rho_1\varepsilon' s'\frac{k'}{K'} + \rho_3\varepsilon'\frac{k_0}{K'}\right) \times \left(1 + s'\frac{\sigma k'}{k'}\right)\alpha_n(1 + \rho_3\varepsilon)(1 + \sigma s), \quad (1)$$

where $\alpha_{n'}$, α_n are Dirac matrices characterizing the electron velocity, $\epsilon = \pm 1$ stands for the sign of the energy, the primed quantities refer to the final state of the electron, and the initial spin direction is chosen to be $\mathbf{s} = s\mathbf{k}/k$ ($\hbar\mathbf{k}$ is the electron initial momentum, which tends to zero; $\hbar cK = \hbar c\sqrt{k^2 + k_0^2}$ is its proper energy).

Summing Eq. (1) over the spins of the final state we obtain:

$$\sum_{s'=\pm 1} \langle \alpha_{n'}, \alpha_n \rangle$$

$$= \frac{1}{8} \operatorname{Sp} \alpha_{n'} \left(1 + \varepsilon' \frac{\boldsymbol{\alpha} \mathbf{k}'}{K'} + \rho_3 \varepsilon' \frac{k_0}{K'}\right) \alpha_n (1 + \rho_3 \varepsilon)(1 + \boldsymbol{\sigma}\mathbf{s}). \quad (2)$$

Similarly, the circular polarization of the photons will be taken into account provided we do not average over the photon polarization states.

The appropriate matrix element is then of the form:[4]

$$S_l = {}^1/_2 \{(\bar{\alpha}^+ \bar{\alpha}) - (\bar{\alpha}^+ \varkappa^0)(\bar{\alpha} \varkappa^0) + il (\varkappa^0 [\bar{\alpha}^+ \bar{\alpha}])\}, \quad (3)$$

where $l = 1$ and $l = -1$ correspond to a right-circular (right handed screw) and left-circular photon polarization, respectively (see also Ref. 2). The quantity κ' describes the photon wave vector and $\bar{\alpha}$ stands for the matrix element

$$\bar{\alpha} = \int \psi^+ \alpha \psi d^3x.$$

The corresponding quantities referring to the scattered photon are denoted by a prime.

From (2) and (3) one obtains a generalization of the Klein-Nishina formula which includes the dependence on the polarization of the incident (l) and scattered (l') photons and on the spin orientation s of the initial electron (the final electron spin has been summed over):

$$d\sigma_{ll'} = \frac{r_0^2 \varkappa'^2 d\Omega}{4\varkappa^2} \Big\{ \frac{\varkappa}{\varkappa'} + \frac{\varkappa'}{\varkappa} - \sin^2\theta + ll' \left(\frac{\varkappa}{\varkappa'} + \frac{\varkappa'}{\varkappa}\right) \cos\theta - l \Big[\left(\frac{\varkappa}{\varkappa'} - \frac{\varkappa'}{\varkappa}\right) \times \cos\theta \cos\psi + \left(1 - \frac{\varkappa'}{\varkappa}\right) \sin\theta \sin\psi \cos\varphi \Big] - l' \Big[\left(\frac{\varkappa}{\varkappa'} - \frac{\varkappa'}{\varkappa} \cos^2\theta - \sin^2\theta\right) \cos\psi + \frac{1}{2} \left(1 - \frac{\varkappa'}{\varkappa}\right) \sin 2\theta \sin\psi \cos\varphi \Big] \Big\}, \quad (4)$$

where $\cos\psi = \boldsymbol{\kappa}^0 \cdot \mathbf{s}$, $\cos\theta = \boldsymbol{\kappa}^0 \cdot \boldsymbol{\kappa}'^0$ and φ is the angle between the $(\boldsymbol{\kappa}^0, \mathbf{s})$ and $(\boldsymbol{\kappa}^0, \boldsymbol{\kappa}'^0)$ planes. Summing in (4) over the final photon polarization (l') we obtain

$$d\sigma_l = \frac{r_0^2 \varkappa'^2 d\Omega}{2\varkappa^2} \Big\{ \frac{\varkappa}{\varkappa'} + \frac{\varkappa'}{\varkappa} - \sin^2\theta - l \Big[\left(\frac{\varkappa}{\varkappa'} - \frac{\varkappa'}{\varkappa}\right) \cos\theta \cos\psi + \left(1 - \frac{\varkappa'}{\varkappa}\right) \sin\theta \sin\psi \cos\varphi \Big] \Big\}, \quad (5)$$

which agrees with the formula of Gunst and Page,[5] found by a different method. Finally, if one averages over l, the usual Klein-Nishina formula is obtained. It follows that in the scattering of unpolarized radiation the spin orientation of the electron does not enter into the integral scattering law.

The degree of circular polarization of the scattered radiation may be calculated from the formula

$$P = \{d\sigma_{ll'}(l' = 1) - d\sigma_{ll'}(l' = -1)\} \Big/ \sum_{l'=\pm 1} d\sigma_{ll'}. \quad (6)$$

In the non-relativistic case we obtain

$$P_{\text{n.r.}} = 2l \cos\theta / (1 + \cos^2\theta). \quad (7)$$

Thus if the incident radiation is unpolarized then $P_{\text{n.r.}} = 0$. In the extreme relativistic case, for small scattering angles ($\cos\theta \gg 1 - k_0/\kappa$) we again obtain formula (7). On the other hand, for large scattering angles ($\cos\theta \ll 1 - k_0/\kappa$), we obtain

$$P_{\text{rel.}} = (l \cos\theta - \cos\psi) / (1 - l \cos\theta \cos\psi). \quad (8)$$

Hence, in the extreme relativistic approximation and for large scattering angles the scattered radiation will be partially circularly polarized even if the incident radiation is unpolarized.

As the inverse of the above problem we may consider two-photon annihilation of longitudinally polarized positrons by electrons at rest with definite spin directions. Considering the two-photon pair annihilation as a transition of an electron into a free negative energy level we may use, in the calculation of the corresponding matrix element, the expression (1) provided we put in it $\mathbf{k}' = -\mathbf{k}_+$, $\epsilon' = -1$, $s' = s_+$. We then find, after summing over the polarizations of the produced photons,

$$d\sigma = \frac{r_0^2 k_0 (K_+ + k_0) d\Omega}{4k_+ (K_+ + k_0 - k_+ \cos\theta)^2} \{J_0 + s_+ (J_1 \cos\psi - J_2 \sin\psi \cos\varphi)\}. \quad (9)$$

Here

$$J_0 = K_+ (1 - \cos\alpha) - k_0 (1 + \cos\alpha) \cos\alpha;$$

$$J_1 = K_+ (1 - \cos\alpha) \cos\alpha + k_0 (1 + \cos\alpha) - \frac{k_0}{k_+} (K_+ + k_0)(1 + \cos\alpha)(\cos\theta + \cos(\alpha - \theta));$$

$$J_2 = \frac{k_0}{k_+} (K_+ + k_0)(1 + \cos\alpha)(\sin\theta - \sin(\alpha - \theta));$$

$$\cos\alpha = \frac{(\boldsymbol{\varkappa} \cdot \boldsymbol{\varkappa}')}{\varkappa \varkappa'} = \frac{K_+ \cos^2\theta + k_0 \sin^2\theta - k_+ \cos\theta}{k_+ \cos\theta - K_+};$$

$$\cos\theta = (\boldsymbol{\varkappa} \cdot \mathbf{k}_+) / \varkappa k_+; \qquad \cos\psi = (\mathbf{s} \cdot \mathbf{k}_+) / k_+;$$

φ is the angle between the $(\mathbf{k}_+, \mathbf{s})$ and $(\mathbf{k}_+, \boldsymbol{\kappa})$ planes, $\hbar \mathbf{k}_+$, $c\hbar K_+$ is the momentum and energy of the positron, s_+ is double the positron spin projection onto its direction of motion, and $\boldsymbol{\kappa}$, $\boldsymbol{\kappa}'$

are the wave vectors of the produced photons. Let us note that in the non-relativistic case $\alpha \to \pi$ and therefore (9) yields, after integration over $d\Omega$, the well known formula

$$\sigma_{\text{n.r.}} = \pi r_0^2 \frac{c}{v_+} (1 - s_+ \cos \psi), \qquad (10)$$

with the consequence that two-photon annihilation is forbidden in this approximation if the electron and positron spins are parallel, i.e., $s_+ \cos \psi = 1$.

One sees from (9) that there is polar (ψ is replaced by $\pi - \psi$) and azimuthal (φ is replaced by $\pi - \varphi$) asymmetry. This fact allows the use of two-photon annihilation of positrons by oriented electrons for experimental determination of the degree of longitudinal polarization of the positrons.

It is known[6] that the asymmetry in the scattering of positrons by oriented electrons is small at low energies and becomes appreciable only at relatively high energies.

On the other hand, two photon annihilation of oriented positrons and electrons gives a noticeable asymmetry in the produced photons only at comparatively low energies. In the high-energy region observation of the asymmetry will apparently be difficult, owing to the decrease in the absolute value of the cross section.

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QUANTUM YIELD OF PHOTOIONIZATION IN SILICON

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IT was shown in earlier works[1,2] that, if the photon energies are high enough, the quantum yield of the internal photoeffect in germanium crystals increases to values much above unity. As was already indicated,[2] the increased quantum yield must be ascribed to impact ionization by primary electrons or holes liberated upon absorption of a photon and having the necessary excess energy.

One would expect a similar phenomenon to take place in silicon. To study the photoeffect in silicon we have used crystals with p-n junctions obtained by thermal diffusion of phosphorus into type-p silicon.[3,4]

Since the diffusion length of the nonequilibrium carriers in the silicon employed was relatively low, we prepared crystals in which the depth of the p-n junction under the illuminated surface did not exceed 2 microns, in order to obtain the necessary sensitivity in the short-wave region. To measure the short-circuit current between the p and n regions of the crystal, and to determine the light flux and the reflection coefficients, we employed a setup similar to that used in the germanium experiments.[2] The interpretation of the results obtained was more difficult in our case than in the determination of the quantum yield in germanium, owing to the following circumstance: The expression for the carrier-collection coefficient α contains the carrier mobility and the diffusion length of the nonequilibrium carriers. However, in the case of a p-n junction obtained by diffusion of impurities from the surface, the mobility changes greatly with depth, increasing inward from the surface where the concentration of the impurity (phosphorus) is high. The carrier diffusion length also undoubtedly varies with depth. In view of this, the formula for the quantum yield Q

$$Q = I_{\text{sc}} / N_{h\nu} q \alpha (1 - R),$$

derived for germanium, cannot be used here (I_{sc} is the short circuit current between the p and n regions, q the electron charge, R the reflection coefficient, and $N_{h\nu}$ the number of photons incident per second. It is therefore necessary to as-