age can be calculated,

$$\Delta m = (0.542 \pm 0.004)\times 10^{10}\ \sec^{-1}. \qquad (7)$$

We can now re-evaluate the "vacuum regeneration" experiments[2,3] using (7). Reference 2 now yields

$$\varphi_{+-} = (44.7 \pm 4.0)^\circ, \qquad (8)$$

while Ref. 3 gives $(49 \pm 12)^\circ$; the main effect is that the uncertainty due to Δm (not included in the errors just quoted!) is now reduced to $\pm 2^\circ$. Combining these values and taking into account the error due to uncertainties in Δm, we get

$$\varphi_{+-} = (45.2 \pm 4.0)^\circ.$$

The value is consistent with those theories which predict

$$\varphi_{+-} \simeq \tan^{-1} 2\Delta m \tau_s = (43.2 \pm 0.4)^\circ,$$

in particular the superweak theory, for which this prediction is exact.

We are indebted to Mr. R. Norton, Mr. T. A. Nunamaker, Mr. T. Shea, and Mr. R. Wall for much assistance with the operation of our apparatus, and to Mr. D. Cosgrove for help in modifying the K^0 facility at Argonne National Laboratory. We wish to thank Professor W. K. H. Panofsky and his staff for the hospitality extended to several of us at the Stanford Linear Accelerator Center, and for the use of the facilities of the computation center.

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PHOTON SPLITTING IN A STRONG MAGNETIC FIELD

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We determine the absorption coefficient and polarization selection rules for photon splitting in a strong magnetic field, and describe the possible application of our results to pulsars.

Recent work on pulsars suggests the presence of trapped magnetic fields within an order of magnitude in either direction of the electrodynamic critical field $B_{cr} = m^2/e = 4.41 \times 10^{13}$ G.[1] (Here m and e are, respectively, the electronic mass and charge.) In such intense fields, electrodynamic processes which are unobservable in the laboratory can become important. One such process, for photons with energy $\omega > 2m$, is photopair production, for which both the photon absorption coefficient and the corresponding vacuum dispersion have been calculated by Toll.[2] For $\omega < 2m$ the photopair process is kinematically forbidden, and the only[3] photon absorption mechanism which does not require the presence of matter is photon splitting, i.e.,

$$\gamma(k) + \text{external magnetic field} \rightarrow \gamma(k_1) + \gamma(k_2). \quad (1)$$

We present in this note the results of calculations of the absorption coefficient and the polarization selection rules for this reaction, in the case of a constant and spatially uniform external magnetic field $\bar{B}$.[4]

To begin, let us consider photon splitting when dispersive effects caused by the external field are neglected, so that the photon four-momenta satisfy the vacuum dispersion relation

$$k^2 = k_1^2 = k_2^2 = 0. \quad (2)$$

Because the external field $\bar{B}$ is constant and spatially uniform, it cannot transfer four-momentum to the photons, and so the four-vectors k, k_1, k_2 must satisfy four-momentum conservation by themselves,

$$k = \omega(1, \hat{k}) = k_1 + k_2 = \omega_1(1, \hat{k}_1) + \omega_2(1, \hat{k}_2). \quad (3)$$

It is easily seen that Eqs. (2) and (3) can be satisfied only if the three propagation directions $\hat{k}$, $\hat{k}_1$, and $\hat{k}_2$ are identical, which implies that the photon four-vectors are proportional,

$$k_1 = (\omega_1/\omega)k, \quad k_2 = (\omega_2/\omega)k. \quad (4)$$

We will use Eq. (4) to simplify considerably the matrix elements for photon splitting. To leading order in e, the matrix element involving $2n+1$ interactions with the external field comes from the ring diagrams with $2n+4$ vertices which are illustrated in Fig. 1. When all permutations of the vertices are summed over, the matrix element is gauge-invariant, and therefore must couple the three photons and the external field only through their respective field-strength tensors $F_{\mu\nu}$, $F_{\mu\nu}^1$, $F_{\mu\nu}^2$, and $\bar{F}_{\mu\nu}$. Because Eq. (4) tells us that only one four-momentum is present in the problem, the matrix element for Fig. 1 is a sum of terms of the form

$$FF^1F^2 \times \underbrace{\bar{F}\cdots\bar{F}}_{2n+1 \text{ factors}} \times \underbrace{k\cdots k}_{2l \text{ factors}} \quad (5)$$

with the Lorentz indices contracted to form a Lorentz scalar (which is why the number of factors k must be even). Since $k^\mu F_{\mu\nu} = k^\mu F_{\mu\nu}^1 = k^\mu F_{\mu\nu}^2 = k^\mu \bar{F}_{\mu\nu} k^\nu = 0$, a nonvanishing contribution is obtained only if each factor k is contracted with a different $\bar{F}$, which means that we must have $l \leqslant n$. We will now show further that when $l = n$, the contribution to the matrix element still vanishes. Writing $v_\mu = \bar{F}_{\mu\nu} k^\nu$, a term with $2n$ factors k has the form

$$FF^1F^2\bar{F} \times \underbrace{v\cdots v}_{2n \text{ factors}}, \quad (6)$$

again with Lorentz indices contracted to form a Lorentz scalar. Because of the antisymmetry of the field strength, the number of factors v which can be contracted with field strengths can only be 0, 2, or 4. An enumeration of these contractions[5] shows that they must always contain at least one factor of the following five types: $F_{\alpha\beta}{}^1F^{2\,\alpha\beta}$, $v_\alpha F^{1\,\alpha\beta} F_{\beta\gamma}{}^2 v^\gamma$, $F_{\alpha\beta} F^{1\,\beta\gamma} F_{\gamma\delta}{}^2 \bar{F}^{\delta\alpha}$, $v_\alpha F^{\alpha\beta} F_{\beta\gamma}{}^1 F^{2\,\gamma\delta} \bar{F}_{\delta\epsilon} v^\epsilon$, or $v_\alpha F^{\alpha\beta} F_{\beta\gamma}{}^1 \bar{F}^{\gamma\delta} F_{\delta\epsilon}{}^2 v^\epsilon$, or factors obtained from these by permuting the photon field strengths F, F^1, and F^2. A simple direct calculation shows that for free photons propagating along the same direction, the five factors always vanish, irrespective of the orientations of the photon polarizations. We conclude, then, that the term with $2n$ factors k vanishes, so that at most $2n-2$ factors k can be present (for $n \geqslant 1$) in the term in the photon-splitting matrix element involving $2n+1$ external field factors $\bar{F}$.

Let us now apply this result to the two smallest ring diagrams: the box diagram ($n=0$) and the hexagon diagram ($n=1$). We immediately learn that the box contribution to photon splitting vanishes identically,[6,7] so that the leading diagram which contributes is the hexagon. Further-

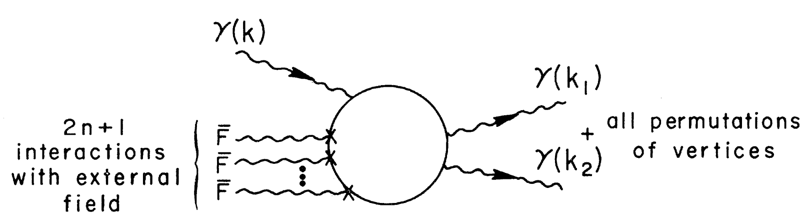

FIG. 1. Ring diagram for photon splitting involving $2n+1$ interactions with the external field.

more, the hexagon contains at most $2\times 1-2=0$ factors of k in addition to those contained in the field strengths, which means that the hexagon diagram is given exactly by its constant-field-strength limit, which can in turn be easily calculated from the Heisenberg-Euler[8] effective Lagrangian. Larger ring diagrams will, of course, also be present, but for purposes of rough order-of-magnitude estimates the leading dependence on $\overline{B}/B_{\rm cr}$ and ω/m given by the hexagon should be sufficient. Carrying out the effective-Lagrangian calculation for the hexagon gives the following formulas for the photon-spliting absorption coefficients in the various photon polarization states:

$$\kappa[(\parallel)\to(\parallel)_1+(\parallel)_2]=\frac{\alpha^3}{60\pi^2}\left(\frac{48}{315}\right)^2\left(\frac{\omega}{m}\right)^5\left(\frac{\overline{B}\sin\theta}{B_{\rm cr}}\right)^6 m=0.39\left(\frac{\omega}{m}\right)^5\left(\frac{\overline{B}\sin\theta}{B_{\rm cr}}\right)^6\ {\rm cm}^{-1},$$

$$\kappa[(\parallel)\to(\perp)_1+(\perp)_2]=\frac{\alpha^3}{60\pi^2}\left(\frac{26}{315}\right)^2\left(\frac{\omega}{m}\right)^5\left(\frac{\overline{B}\sin\theta}{B_{\rm cr}}\right)^6 m=0.12\left(\frac{\omega}{m}\right)^5\left(\frac{\overline{B}\sin\theta}{B_{\rm cr}}\right)^6\ {\rm cm}^{-1},$$

$$\kappa[(\perp)\to(\parallel)_1+(\perp)_2]+\kappa[(\perp)\to(\perp)_1+(\parallel)_2]=2\kappa[(\parallel)\to(\perp)_1+(\perp)_1]. \tag{7}$$

Here $\alpha=e^2\approx 1/137$ is the fine structure constant, θ is the angle between the photon propagation direction $\hat{k}$ and the direction $\hat{b}$ of the external magnetic field, and the linear polarization eigenmodes are labeled $\parallel$ or $\perp$ according to whether the $\vec{B}$ vector of the eigenmode lies in, or is normal to, the $\hat{k}$-$\hat{b}$ plane. Only formulas for processes involving an even number of $\perp$ photons have been given; the absorption coefficients for processes involving an odd number of $\perp$ photons vanish by a simple CP argument. To see this, we note that for each $\parallel$ photon, the matrix element will contain a CP-even factor $\vec{B}^{\rm photon}\cdot\hat{b}$ [the only other possible scalar product, $\vec{B}^{\rm photon}\cdot\vec{k}$, vanishes by transversality], while for each $\perp$ photon it will contain a CP-odd factor $\vec{E}^{\rm photon}\cdot\hat{b}$. Since the only scalar not involving the photon fields is $\vec{k}\cdot\hat{b}$, which is CP even, the matrix elements for the odd-$\perp$-photon processes are CP odd, and hence vanish.

So far we have assumed that the photons satisfy the vacuum dispersion relation of Eq. (2). Actually, because of the absorptive processes taking place in the external field, there will be dispersive effects which modify Eq. (2). A simple CP argument shows that the photon eigenmodes remain linearly polarized, with the parallel and perpendicular characters described above, but with the ratio of wave number to frequency changed from unity to

$$k/\omega=n_{\parallel,\perp}. \tag{8}$$

The indices of refraction $n_{\parallel,\perp}$ can be calculated from the total absorption coefficients $\kappa_{\parallel,\perp}$ by Kramers-Kronig (dispersion) relations, with the dominant contribution coming from photopair production [the contribution from photon splitting is smaller by a factor $\sim(\alpha/\pi)^2(\overline{B}\sin\theta/B_{\rm cr})^4$, and can be neglected]. For small $\overline{B}/B_{\rm cr}$ the calculation has been carried out numerically by Toll,[2] who gives curves for $n_{\parallel,\perp}$ as a function of frequency ω. When we have $\omega<2m$, so that the parameter $x=(\omega/2m)(\overline{B}\sin\theta/B_{\rm cr})$ is also small, Toll's results can be approximated analytically by

$$n_{\parallel,\perp}=1+\frac{\alpha}{\pi}\left(\frac{\overline{B}\sin\theta}{2B_{\rm cr}}\right)^2 N_{\parallel,\perp}(x),$$
$$N_{\parallel}(x)\sim 0.18+0.24x^2,$$
$$N_{\perp}(x)\sim 0.31+0.44x^2. \tag{9}$$

When dispersive effects are taken into account, the equations of conservation of four-momentum become

$$\omega=\omega_1+\omega_2$$
$$n(\omega)\omega\hat{k}=n(\omega_1)\omega_1\hat{k}_1+n(\omega_2)\omega_2\hat{k}_2, \tag{10}$$

with each n the refractive index appropriate to the respective photon polarization state. These conditions can be simultaneously satisfied only if

$$0\leq\Delta=n(\omega_1)\omega_1+n(\omega_2)\omega_2-(\omega_1+\omega_2)n(\omega_1+\omega_2), \tag{11}$$

in which case the photon propagation directions are not precisely parallel, but rather diverge from one another by small angles $\sim(\Delta/\omega)^{1/2}$. As a result of this nonparallelism, the box diagram is no longer precisely zero, but a careful estimate shows that it is still much smaller than the hexagon. When Δ is negative, the photon-splitting reaction is forbidden. Substituting the indices of refraction of Eq. (9) into Eq. (11) shows, indeed, that for small x the only reaction in Eq. (7) which is kinematically allowed is $(\parallel)\to(\perp)_1+(\perp)_2$, and that this reaction occurs without restriction on the photon frequencies ω_1 and ω_2.

Table I. Selection rules for photon splitting. *CP*-forbidden reactions are suppressed by a factor $\sim(\alpha/\pi)^2(\bar{B}\sin\theta/B_{cr})^4$ relative to *CP*-allowed cases.

Reaction	*CP* selection rule	Small-x kinematic selection rule
$(\Vert)\to(\Vert)_1+(\Vert)_2$	Allowed	Forbidden
$(\Vert)\to(\Vert)_1+(\perp)_2,\ (\perp)_1+(\Vert)_2$	Forbidden	Allowed
$(\Vert)\to(\perp)_1+(\perp)_2$	Allowed	Allowed
$(\perp)\to(\Vert)_1+(\Vert)_2$	Forbidden	Forbidden
$(\perp)\to(\Vert)_1+(\perp)_2,\ (\perp)_1+(\Vert)_2$	Allowed	Forbidden
$(\perp)\to(\perp)_1+(\perp)_2$	Forbidden	Forbidden

The various polarization selection rules for photon splitting are summarized in Table I. Because of the small nonparallelism of the photons in the kinematically allowed regions, the "*CP*-forbidden" reactions are not precisely forbidden, but are down by a factor $\sim(\alpha/\pi)^2(\bar{B}\sin\theta/B_{cr})^4$ relative to the "*CP*-allowed" cases. We see that for small x, all reactions by which perpendicularly polarized photons might split are kinematically forbidden, while parallel-polarized photons split predominantly into perpendicularly polarized photons. Hence photon splitting provides a mechanism for the production of linearly polarized γ rays.[9]

To conclude, let us briefly discuss the possible application of our results to pulsars. We assume that the hexagon-diagram absorption coefficients in Eq. (7) can be used for order-of-magnitude estimates even when the parameters $\bar{B}/B_{cr}$ and ω/m are of order unity.[10] Taking, for illustration, $\bar{B}/B_{cr}\sim\omega/m\sim\sin\theta\sim1$, we find $\kappa[(\Vert)\to(\perp)_1+(\perp)_2]\sim0.1\ \text{cm}^{-1}$. This gives 10^5 absorption lengths in the characteristic distance $R_{\text{pulsar}}\sim10^6$ cm over which the trapped magnetic field has its maximum strength, indicating that photon splitting can be an important absorption mechanism for γ rays emitted near the pulsar surface. Before we can apply the kinematic polarization selection rules to the pulsar problem, two questions must be dealt with. First, since Toll's curves for the indices of refraction were obtained assuming $\bar{B}/B_{cr}$ small, an extrapolation is involved in extending the selection rules forbidding perpendicular-photon decay and parallel-photon decay into parallel photons to the region where $\bar{B}/B_{cr}$ is of order unity. However, Toll's photopair production curves show that $\kappa_\perp>\kappa_\Vert$ when $\bar{B}/B_{cr}$ is unity. By combining this fact with the Kramers-Kronig relations, one easily sees that the selection rules in Table I hold as long as $\omega<2m$, and hence the extrapolation is justified. Second, one expects a plasma to be present near the pulsar surface which will contribute additional dispersive terms to the inequality of Eq. (11). For a plasma-electron density of 10^{17}-10^{-19} cm^{-3} (in rough accord with current pulsar models[10]), a detailed estimate shows that plasma-induced splitting of perpendicularly polarized photons occurs with an absorption coefficient of at most 10^{-9}-10^{-7} times the absorption coefficient for the allowed reaction $(\Vert)\to(\perp)_1+(\perp)_2$, and therefore will be completely negligible.[11] We conclude that if γ rays in the range 0.5-1 MeV are emitted near the pulsar surface, and if the surface magnetic field is as large as B_{cr}, only those gammas with perpendicular polarization will escape. A distant observer would see linearly polarized gammas, with their $\vec{B}$ vector perpendicular to the plane containing the line of sight and the traversed pulsar magnetic field.[12]

A detailed account of the calculations summarized here will be presented elsewhere.[10] The authors wish to thank P. Goldreich (who first brought this problem to our attention), E. P. Lee, and M. Rassbach for informative conversations. After this manuscript was completed, we learned that some of our results have been obtained independently by Z. Bialynicka-Birula and I. Bialynicki-Birula.[13]

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[1]For a recent view, see F. Pacini, "Neutron Stars, Pulsar Radiation and Supernova Remnants," to be published. We use unrationalized Gaussian units, with $\hbar=c=1$.

[2]J. Toll, dissertation, Princeton University, 1952 (unpublished).

[3]The phase space for a photon to split into three or more photons vanishes.

[4]In a pulsar, the field $\bar{B}$ varies over a characteristic distance of $R_{\text{pulsar}}\sim10^6$ cm, but this can be shown to have a negligible effect on our results.

[5]We need not consider contractions involving the

antisymmetric tensor $\epsilon_{\alpha\beta\gamma\delta}$, because by parity the number of such factors must be even, and they can be eliminated pairwise in terms of Kronecker deltas by means of the identity

$$\epsilon_{\alpha\beta\gamma\delta}\epsilon_{\alpha'\beta'\gamma'\delta'}=\sum_{\mathrm{Perm}(\alpha'\beta'\gamma'\delta')}(-1)^P g_{\alpha\alpha'}g_{\beta\beta'}g_{\gamma\gamma'}g_{\delta\delta'}.$$

[6]This result for the box diagram has been obtained independently by M. Rassbach.

[7]This disagrees with the conclusion of V. G. Skobov, Zh. Eksp. Teor. Fiz. 35, 1315 (1958) [Sov. Phys. JETP 8, 919 (1959)], who failed to make a properly gauge-invariant calculation of the box. Skobov's result is quoted in the review article of T. Erber, Rev. Mod. Phys. 38, 626 (1966).

[8]W. Heisenberg and H. Euler, Z. Phys. 38, 714 (1936).

[9]Toll (Ref. 2) points out that in a frequency interval $\Delta\omega \sim e\overline{B}/m$ just above the photopair threshold at $\omega = 2m$, the photopair process acts as a linear polarizer of the opposite sense, absorbing photons of perpendicular polarization, but not those of parallel polarization.

[10]For further discussion, see S. L. Adler, to be published.

[11]In a plasma, the propagation eigenmodes become elliptically polarized, but are still "almost plane $\parallel$" and "almost plane $\perp$" in nature. Faraday rotation, which arises from interference between two unattenuated propagation eigenmodes of different phase velocities, cannot occur in our case since only the "almost plane $\perp$" eigenmode propagates without attenuation. As a result of photon splitting the "almost plane $\parallel$" eigenmode is rapidly absorbed.

[12]Solid-state linear-polarization analyzers for gammas in this energy range have been described by G. T. Ewan *et al.*, Phys. Lett. 29B, 352 (1969).

[13]Z. Bialynicka-Birula and I. Bialynicki-Birula, to be published.

ONE-PION CONTRIBUTION TO NEUTRON-PROTON CHARGE-EXCHANGE SCATTERING*

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Recent measurements of neutron-proton charge-exchange scattering showed a change in the behavior of the cross section at $u=0$, near the one-pion threshold. The slope of the backward cross section also increased rather rapidly, with a maximum also near the one-pion threshold. We point out that the one-pion-exchange amplitude, interfering with a slowly varying amplitude at $u=0$, gives qualitatively the same behavior as observed in the experiments. A similar explanation has been suggested as an explanation for the n-p charge-exchange data above 2 GeV.

The vast amount of nucleon-nucleon scattering data that has accumulated over the past 30 years has shown remarkably little structure. In recent work at the Princeton-Pennsylvania Accelerator (PPA), Mischke, Shepard, and Devlin[1] have produced data on neutron-proton charge-exchange scattering which showed rather clear evidence of structure at $u=0$, near the one-pion production threshold.[2] The purpose of this Letter is to point out that this structure seems to result from interference between the one-pion exchange (OPE) amplitude and an amplitude which is slowly varying for $u \approx 0$.

The relevant data are reproduced in Fig. 1. The interesting characteristic features of this data set[3] are (1) an anomaly in the energy dependence of the backward neutron-proton elastic scattering cross section, with a dip at about 200 MeV and a small but perceptible rise between 200 and 400 MeV [Fig. 1(b)]; and (2) a peak in the quantity $\beta \equiv [(d/dt)\ln(d\sigma/dt)]_{u=0}$ centered about the pion-production threshold energy [Fig. 1(a)].

We have examined the PPA data in comparison with phenomenological NN analyses for energies up to 750 MeV. The nucleon-nucleon phase shift analyses of MacGregor, Arndt, and Wright[4] (hereafter called the Livermore phase shifts), which antedate these experimental measurements, provide good average fits to all of the available scattering data. For energies below 400 MeV, the Livermore phases are evaluated as essentially free parameters for angular momentum $l \leq 5$, and are taken as the OPE Born phase shifts[5] for $l > 5$. In Fig. 1, we plot the quantity β and the backward differential cross section for neutron kinetic energies up to 600 MeV, using the Livermore phases for $l \leq 5$ with OPE phase shifts for $6 \leq l \leq 20$ (we used $g_\pi^2/4\pi = 14.4$ for the $NN\pi$ coupling). These are the solid curves in Fig. 1. The qualitative agreement with the data is good; that is, the shelf in the backward differential cross section is reproduced, and the peak in β occurs at the same energy as the PPA data. The detailed fit to these data will be discussed later