

This item is likely protected under Title 17 of the U.S. Copyright Law. Unless on a Creative Commons license, for uses protected by Copyright Law, contact the copyright holder or the author.

Access to this work was provided by the University of Maryland, Baltimore County (UMBC) ScholarWorks@UMBC digital repository on the Maryland Shared Open Access (MD-SOAR) platform.

Please provide feedback

Please support the ScholarWorks@UMBC repository by emailing scholarworks-group@umbc.edu and telling us what having access to this work means to you and why it's important to you. Thank you.

THE THEORY OF INVERSE COMPTON SCATTERING IN A COLD PLASMA*

MELVYN L. GOLDSTEIN† AND ALLEN M. LENCHEK

Department of Physics and Astronomy, University of Maryland, College Park, Maryland

Received 1970 May 22

ABSTRACT

We have derived a theory for inverse Compton scattering in a cold collisionless plasma. In the absence of a magnetic field, the general shape of the scattered-photon spectrum is qualitatively similar to the vacuum spectrum. The average rate of electron-energy loss is three-quarters of the vacuum value if the frequency of the incident photons is close to the plasma frequency. In order for the plasma to have a large effect on the scattered-photon spectrum, the incident radiation must propagate below the plasma frequency, which is possible if a magnetic field is present. As an example, we use gyrosynchrotron photons propagating in the ordinary mode below the plasma frequency as the source of incident photons. When these photons are scattered, the spectrum is greatly suppressed compared with the vacuum results. The total scattering cross-section is also sharply decreased in a magnetoactive plasma.

I. INTRODUCTION

Inverse Compton scattering of low-energy photons by high-energy electrons has been recognized in astrophysics both as an important source of high-energy photons and as an energy sink for the electrons. The theory of inverse Compton scattering is closely related to the theory of synchrotron radiation. When synchrotron radiation is produced by electrons moving in a cold collisionless plasma, the spectrum of emitted radiation is suppressed. This suppression phenomenon has been referred to as the Razin effect (for example, see Razin 1957, 1960; Ramaty 1968). The spectrum of inverse Compton photons is also modified by the presence of a cold plasma. The modifications become important when the electron energy is small ($\gamma mc^2 \sim 50$ MeV) and the incident photons are propagating with frequencies well below the plasma frequency. In § II, we develop the theory which describes the modifications to the vacuum spectrum when no magnetic field is present. In § III we extend our results to include a nonzero magnetic field. In this situation, one has incident electromagnetic radiation below the plasma frequency, and the effects of the plasma are important.

II. SCATTERING IN A PLASMA WITH NO MAGNETIC FIELD

When electrons and photons propagate in a medium with an index of refraction $n < 1$, the correct Compton cross-section differs from the one used in the vacuum theory. We found it easiest to derive the appropriate cross-section in the laboratory frame of reference, where the bulk velocity of the plasma is zero. (Of course, the derivation can also be done in the rest frame of the electron, as is usual in the vacuum theory.) It can be shown that it is sufficient to derive the Thomson differential cross-section describing the laboratory interaction in which the electron and photon are incident head-on.

In the laboratory frame, let the plasma be isotropic and homogeneous and have zero bulk velocity. The index of refraction is $n = (1 - \alpha_p^2/\alpha^2)^{1/2}$, where

$$\alpha_p = \frac{e}{mc^2} (N_e/m\pi)^{1/2}$$

* Research supported by NASA grant NGL 21-002-033

† Present address: Dept. of Environmental Sciences, Tel-Aviv University, Ramat-Aviv, Israel.

is the electron plasma frequency in units of electron rest energy, N_e is the electron number density in the plasma, and α is the energy of the photon. (All energies will be measured in units of electron rest energy.) The differential cross-section is (Goldstein 1969; also see Pechacek 1966; Pechacek and Trivelpiece 1967)

$$\frac{d\sigma}{d(\cos \theta)} = \frac{2\pi r_0^2 n}{\gamma^2(1 + \beta n \cos \theta)^4} [(1 + \beta n \cos \theta)^2 - \frac{1}{2}(1 - n^2\beta^2) \sin^2 \theta] \quad (1)$$

where

$$r_0^2 = e^2/mc^2,$$

β is the electron's velocity in units of c , and θ is the angle between $-\beta$ and the direction of propagation of the radiated photon. With suitable assumptions, discussed in § III below, it will be possible to use equation (1) even when the magnetic field is not zero.

When equation (1) is integrated over solid angle with $n = 1$, the result is σ_T , the total Thomson cross-section, as expected. However, integrating equation (1) with $n \neq 1$ (now numerically) will not give σ_T . This is most easily understood by setting β equal to 0 so that the electron and plasma are at rest. As the incident frequency approaches the plasma frequency, the photon group velocity ($cn_1 = cn$) approaches zero, and there can be no scattering. Consequently, one does not expect the total scattering cross-section to be σ_T independent of plasma parameters and electron velocity.

We can now derive the scattered photon spectrum in the presence of a cold plasma, following the notation and general procedure of Jones (1968). We assume that the electron has an energy γmc^2 , and the photons are all monoenergetic and isotropically distributed in the laboratory system with energies α_1 and α before and after scattering, respectively. Let primes denote quantities measured in the electron rest frame. Figure 1 shows the geometry of the scattering event in the electron rest frame. The photon is incident at an angle θ'_1 measured with respect to $-\beta$. The scattering angle is χ' ; the azimuth is ϕ' (measured so that the $[\alpha', \beta]$ -plane is the $[\phi' = 0]$ -plane); and θ' is the angle made by the scattered photon with respect to $-\beta$. We will make the approximation that $\theta'_1 = 0$. This is the usual approximation made in the literature in the vacuum theory (see, e.g., Felten and Morrison 1966; Ginzburg and Syrovatskii 1964). Jones (1968) has derived the scattered-photon spectrum in a vacuum without setting θ'_1 equal to 0. His results indicate that the approximation is very good for $\alpha < \alpha_1$ when $\gamma \geq 2$. However, the low-frequency end of the spectrum, $\alpha/\alpha_1 \approx 1/(4\gamma^2)$, is exaggerated when this approximation is made. When a plasma is present, we show below that our results in the frequency range $\alpha < \alpha_1$ are not meaningful because of this approximation, but the fraction of photons scattered into this region is small, and the spectrum above α_1 is unaffected by this approximation when $\gamma \geq 2$. One further approximation is to choose α_1

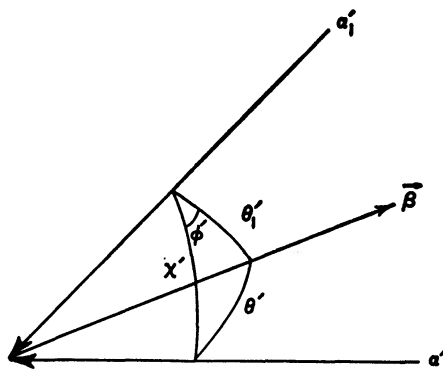


FIG. 1.—The scattering process in the electron rest frame

and γ so that $4\alpha_1\gamma \ll 1$, which is equivalent to assuming that in the electron rest frame the collision can be described as classical Thomson scattering.

Because the incident photons are monoenergetic and isotropic in the laboratory system, their distribution function is given by

$$N_1(\alpha_1, x_1) = \frac{1}{4\pi} \delta\left[\alpha_1 - \frac{\alpha'_1}{\gamma(1 + \beta n_1 x_1)}\right], \quad (2)$$

where n_1 is the index of refraction seen by the incident photon and $x_1 = \cos \theta_1$. The incident-photon energy distribution in the electron rest frame can be shown to be

$$N'_1(\alpha'_1) = \frac{\alpha'_1}{2n_1\gamma\beta\alpha_1^2} S[\alpha'_1; \alpha_1\gamma(1 - \beta n_1), \alpha_1\gamma(1 + \beta n_1)], \quad (3)$$

where $S(x; a, b)$ is defined by

$$S(x; a, b) = 1 \quad \text{for } a \leq x \leq b,$$

$$S(x; a, b) = 0 \quad \text{otherwise.}$$

The number of collisions per unit time t' is $N'c\sigma'$, so that by using equations (1) and (3), the Doppler relation $\alpha' = \gamma\alpha(1 + \beta n \cos \theta)$, and the condition for Thomson scattering, $\alpha'_1 = \alpha'$, one can show that the number of photons scattered per unit volume, time, and energy is

$$\frac{d^2N}{d\alpha dt} = \frac{\pi c r_0^2 \alpha}{\gamma^2 \beta^3 n^2 n_1 \alpha_1^2} \left[\frac{1}{2} \xi (1 + \beta^2 n^2) - (1 - \beta^2 n^2) \ln(\xi) - \frac{(1 - \beta^2 n^2)}{2\xi} \right]_{\xi_L}^{\xi_U} \quad (4)$$

where $\xi = (1 + \beta n \cos \theta)$. If we define $\xi_1 = \alpha_1/\alpha (1 - \beta n_1)$ and $\xi_2 = \alpha_1/\alpha (1 + \beta n_1)$, then the limits ξ_U and ξ_L are determined by what part, if any, of the interval $[\xi_1, \xi_2]$ lies within the bounds $[1 - \beta n, 1 + \beta n]$. (For a detailed discussion of the limits ξ_U and ξ_L , see Goldstein 1969.)

It is easiest to evaluate equation (4) numerically and then plot the resulting spectrum. This is done in Figure 2 for several values of γ and α_p . The figure is a plot of $d^2N/d\alpha dt$

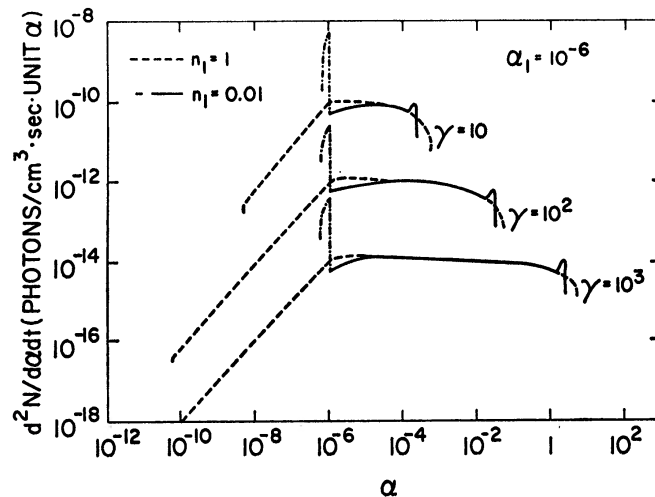


FIG. 2.—Spectra of photons scattered by monoenergetic electrons with three different energies, γ . For each γ , spectra are shown for scattering in a vacuum ($n_1 = 1$) and in a plasma ($n_1 = 0.01$). The energy of the incident photon is $\alpha_1 = 10^{-6}$. All energies are in units of electron rest mass. The sharp spike near $\alpha = \alpha_1$ is not physically meaningful and results from the approximation that $\theta'_1 = 0$.

versus α with $\alpha_1 = 10^{-6}$. The dotted line is the spectrum when $n_1 = 1$ (vacuum), and the solid line for $n_1 = 0.01$. The vacuum results duplicate those of Jones (1968). The cutoff at low frequencies when $n_1 = 0.01$ occurs because the photons cannot propagate with frequencies below the plasma frequency. At high frequencies ($n_1 = 0.01$) the cutoff occurs because the incident-photon momentum, $\alpha_1 n_1 c$, goes to zero as $\alpha_p \rightarrow \alpha_1$. Consequently, the maximum energy that the scattered photon can have is halved. But when $\beta \approx 1$, the total cross-section, regardless of the value of n_1 , must be σ_T . The rise at the high-frequency end of the spectrum, near $\alpha \approx \gamma^2 \alpha_1$, compensates for the faster cutoff. By integrating numerically, one can show that

$$\frac{1}{c} \int_{\alpha_{\min}}^{\alpha_{\max}} \frac{d^2 N}{d\alpha dt} d\alpha = \sigma_T, \quad (5)$$

when $\gamma \gg 1$, for all values of n . Changing the value of α_1 does not affect the shape of the spectrum if the ratio α_p/α_1 is held constant.

We can now investigate the validity of our assumption that $\theta'_1 = 0$. If we take the photon energy distribution in the electron rest frame, equation (3), and transform it back into the laboratory system with *no* scattering [$\alpha'_1 = \alpha' = \gamma\alpha(1 + \beta n)$], we would have (cf. Jones 1968)

$$N(\alpha)d\alpha = \frac{\alpha(1 + \beta n)}{2n_1\beta\alpha_1^2} S\left[\alpha; \alpha_1 \frac{(1 - \beta n_1)}{(1 + \beta n)}, \alpha_1 \frac{(1 + \beta n_1)}{(1 + \beta n)}\right]. \quad (6)$$

The approximation that $\theta'_1 = 0$ produces photons at all energies $\alpha < \alpha_1$. Therefore, this region of the spectrum will be overpopulated solely because of this assumption. As n_1 becomes much less than unity, the low-frequency ($\alpha < \alpha_1$) end of the spectrum (eq. [4]) becomes an increasingly poor approximation of the exact result. The discontinuity seen in Figure 2 at $\alpha = \alpha_1$ (*dot-dashed line*) reflects the breakdown of the assumption $\theta'_1 = 0$, and is not physically meaningful. The spectrum for $\alpha > \alpha_1$ is not affected by this assumption. It would be very difficult to derive the exact spectrum (4) with a plasma present, and one would not expect the results to differ from our approximate derivation when $\alpha > \alpha_1$.

By using equation (4), the rate of electron-energy loss can be calculated from

$$-\frac{dE}{dt} = \int_{\alpha_{\min}}^{\alpha_{\max}} \frac{d^2 N}{d\alpha dt} (\alpha - \alpha_1) d\alpha. \quad (7)$$

When $n_1 = 1$, the result in a vacuum is (Jones 1965)

$$-\frac{dE}{dt} = \frac{4}{3} \sigma_T c w_{\text{ph}} (p/mc^2)^2, \quad (8)$$

where p is the momentum of the electron, and w_{ph} is the energy density of the ambient photon field, defined by

$$w_{\text{ph}} = \int_0^\infty d\alpha_1 N_1(\alpha_1) \alpha_1 d\alpha_1. \quad (9)$$

For $n_1 \neq 0$, the integration must be done numerically. The result is that as $n_1 \rightarrow 0$, dE/dt approaches

$$-\frac{dE}{dt} \approx \sigma_T c w_{\text{ph}} (p/mc^2)^2. \quad (10)$$

The dependence on electron energy is maintained; only the value of the numerical coefficient is changed when dielectric effects are important.

Unless the plasma is so dense that the incident photons cannot propagate ($\alpha_p > \alpha_1$), the effect of the plasma on the inverse Compton photon spectrum and rate of energy loss is not very pronounced. This is analogous to the situation in synchrotron radiation when the plasma frequency is less than the cyclotron frequency. Then the Razin effect is not large, and becomes large only when $\alpha_p \gg \alpha_B = eB/2\pi m^2 c^3$.

For the plasma to have a large effect on the physics of inverse Compton scattering, we need a situation in which the incident photons propagate with frequencies less than the plasma frequency. This is the subject of § III.

III. SCATTERING WITH A MAGNETIC FIELD

Because we are concerned with inverse Compton scattering of real photons, we expect that plasma effects will become important when the plasma frequency exceeds the frequency of the incident photon. Electromagnetic radiation can propagate at frequencies less than the plasma frequency only when a magnetic field is present. However, there are few physical processes that produce electromagnetic radiation below the plasma frequency. When fast electrons move through a magnetoactive plasma, the gyrosynchrotron emission they produce is radiated both above and below the plasma frequency. In order to illustrate the important aspects of inverse Compton scattering of photons below the plasma frequency, in the discussion below we have assumed that the incident photons are produced by gyrosynchrotron radiation and propagate in the ordinary mode below the plasma frequency. To simplify the problem further, we allowed the same electrons to Compton-scatter the incident gyrosynchrotron radiation they produce. However, in astrophysical situations, scattering of gyrosynchrotron photons in particular is not likely to be important because the scattered radiation will probably be swamped by the high-frequency end of the initial photon spectrum unless those frequencies are completely suppressed by the Razin effect or some other mechanism such as a pitch-angle anisotropy of the electrons (Ramaty 1969*a, b*). If other sources of electromagnetic radiation below the plasma frequency are found in regions permeated with fast electrons, the Compton-scattered spectra may very well be observable, in which case they will be similar to those presented below.

Emission spectra of gyrosynchrotron photons have been computed by Liemohn (1965), Ramaty (1969*a*), and Goldstein (1969). Photons are radiated at frequencies and angles given by solutions to the emission equation

$$\alpha_1 = \frac{s\alpha_B}{\gamma(1 - n_1\beta_2\gamma_1)} \quad (11)$$

where s is the harmonic of the emission, $\beta_2 = \beta \cos \phi$, ϕ is the pitch angle of the electron, $\gamma_1 = \cos \theta_1$ (now θ_1 is the angle between the magnetic field \mathbf{B} and the wave normal \hat{K} , see Fig. 3), and, with a magnetic field, n_1 is given by

$$n_1^4 = 1 + 2\alpha_p^2(\alpha_p^2 - \alpha_1^2)/(\alpha_B^4 D_+) , \quad (12)$$

where

$$D_+ = \frac{1}{\alpha_B^2} \left[\alpha_1^4 \sin \theta_1 + \frac{4\alpha_1^2}{\alpha_B^2} (\alpha_p^2 - \alpha_1^2)^2 \cos^2 \theta_1 \right]^{1/2} - \frac{2\alpha_1^2}{\alpha_B^4} (\alpha_p^2 - \alpha_1^2) - \frac{\alpha_1^2}{\alpha_B^2} \sin^2 \theta_1 ,$$

and the subscript (+) refers to the ordinary (whistler) mode of propagation.

The distribution function of incident photons is no longer isotropic. One can show that the photon distribution function, normalized to one incident photon per unit volume, in the electron rest frame is given by (Goldstein 1969)

$$N'_1(a_1, \phi) = N_1(a_1, \phi) S(a'_1; a_L, a_U) , \quad (13)$$

where

$$N_1(\alpha_1, \phi) = A(\alpha_1, y_s) / \int d\alpha_1 A(\alpha_1, y_s) ,$$

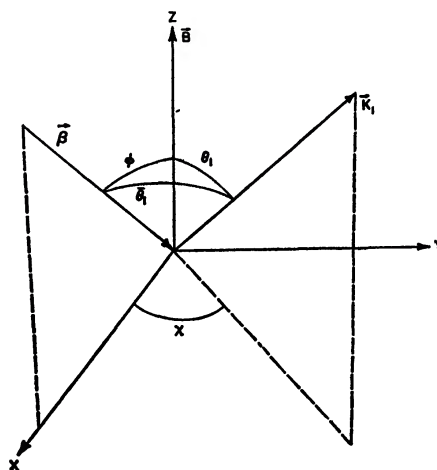


FIG. 3.—The geometry before scattering in the laboratory reference frame in a magnetoactive plasma

$$A(\alpha_1, y_s) = \sum_{s=-\infty}^{s=+\infty} \left[\frac{1}{(1 + a_y^2)(n_1 + \cos \theta_s \partial n_1 / \partial \cos \theta)|_{\theta=\theta_s}} \right. \\ \left. \times n_1 \left\{ -\beta_1 \left[J'_s(x_s) + a_y \left(\frac{\cot \theta_s}{n_1} - \frac{\beta_2}{\sin \theta_s} \right) \right] J_s(x_s) \right\}^2 \left(1 + \frac{\partial \ln n_1}{\partial \ln \alpha_1} \right) \right],$$

$\beta_1 = \beta \sin \phi$, and a_y is the polarization coefficient defined by

$$a_y = - \frac{\alpha_1 y_s}{\alpha_B} \left[\frac{\alpha_1^2}{\alpha_B^2} + \frac{\alpha_p^2}{\alpha_B^2} / (n_1^2 - 1) \right]^{-1}.$$

(We have taken the other polarization coefficient, a_k , to be zero. This is equivalent to taking the component of electric field parallel to B to be zero, and will be justified below.) In addition,

$$x_s = s n_1 \beta_1 (1 - y_1^2)^{1/2} / (1 - n_1 \beta_2 y_1),$$

θ_s is the solution to the emission equation (11), J_s, J'_s are the Bessel function of order s and its derivative, respectively, and

$$\alpha_{L,U} = \frac{S \{ 1 \mp \beta / [n_1 (1 + \partial \ln n_1 / \partial \ln \alpha_1)] \}}{(1 - n_1 \beta_2 \cos \theta_s)}.$$

As an example, $N_1(\alpha_1, \phi)$ is evaluated for $\alpha_p/\alpha_B = 3$, $\phi = 45^\circ$, and $\gamma = 10$. The photon density below α_p is shown as a function of α_1 in Figure 4, averaged over a suitable bandwidth of α_1 , in order to smooth out the rapid oscillations characteristic of gyrosynchrotron spectra that have been integrated over emission angle θ_1 (Liemohn 1965). Examples of spectra evaluated for different values of α_p/α_B , ϕ , and α can be found in Goldstein (1969). For simplicity, we assume that the photons are scattered by electrons having the same α and ϕ as those which produced the incident radiation. Within the context of this cold plasma theory, n_1 has resonances at various frequencies (Liemohn 1965). However, the magnitude of n_1 would be bounded if the thermal properties of the plasma were rigorously included. We can estimate roughly a limit on n_1 if we note that the radiation will be unable to propagate if the thermal velocities of the plasma become comparable to the phase velocity of the wave, c/n_1 . Thus we want

$$n_1 \ll \left(\frac{m c^2}{2 k T} \right)^{1/2} \quad (14)$$

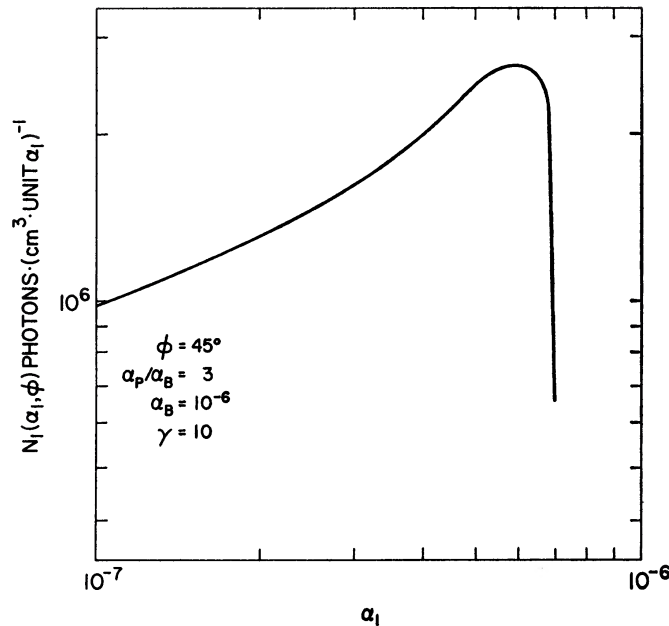


FIG. 4.—Incident-photon density $N_1(\alpha_1, \phi)$, as a function of frequency, α_1 , averaged over a suitable bandwidth of α_1 to smooth out the rapid oscillations characteristic of gyrosynchrotron spectra that have been integrated over emission angle. The photons are produced by gyrosynchrotron emission of electrons with $\gamma = 10$ and $\phi = 45^\circ$, moving through a magnetoactive plasma with $\alpha_B = 10^{-6}$ and $\alpha_p = 3 \times 10^{-6}$.

where k is Boltzmann's constant. Arbitrarily, we will restrict n_1 to be ≤ 50 , corresponding to temperatures of $T \ll 10^6$ ° K.

In a magnetoactive plasma, the electric vector of the radiation field has a nonzero longitudinal component, E_k . Therefore, Poynting's vector is not in the same direction as the propagation vector, \hat{K} . The group velocity vector, which will be in the same direction as the time-averaged Poynting vector, will not be parallel to the phase velocity (Landau and Lifshitz 1960). This property of a magnetoactive plasma would greatly complicate any derivation of a cross-section for inverse Compton scattering in such a medium. However, one can show (Goldstein 1969; Ramaty 1969a) that the emissivity associated with the longitudinal component, E_k , is small. Consequently, we have set $a_k = 0$ and can use equation (1) for our scattering cross-section.

We will not be concerned with the effects of the transformed magnetic field on the electron's motion. The acceleration produced by the transformed field causes the gyrosynchrotron emission, which we have already calculated. We can assume that the Compton scattering occurs in a time interval much smaller than the gyroperiod of the electron, so that we can approximate the instantaneous rest frame of the electron as an inertial frame, and develop the theory as in § II above.

We limit the kinematics to scattering events in which the frequency of the scattered photon is above the plasma frequency. Otherwise, the Compton-scattered radiation would be swamped by the ambient gyrosynchrotron flux. Following the derivation outlined in § I, one can show that the number of Compton-scattered, low-frequency ($\alpha_1 < \alpha_p$) gyrosynchrotron photons produced per unit volume, time and energy is

$$\frac{d^2 N}{d\alpha d\tau} = \frac{\pi c r_0^2 N'_1(\alpha_1, \phi)}{\beta^3 \gamma^2 n^2} \left[(1 + \beta^2 n^2) \ln \xi + \frac{2(1 - n^2 \beta^2)}{\xi} - \frac{1}{2} \frac{(1 - n^2 \beta^2)}{\xi^2} \right]_{\xi_L}^{\xi_U}, \quad (15)$$

where $\xi = 1 + \beta n \cos \theta$, and the limits ξ_U and ξ_L are found from the part of the interval $[\xi_1, \xi_2]$ within the bounds $[1 - \beta n, 1 + \beta n]$, where now $\xi_{1,2} = \alpha_{L,U}/\gamma\alpha$.

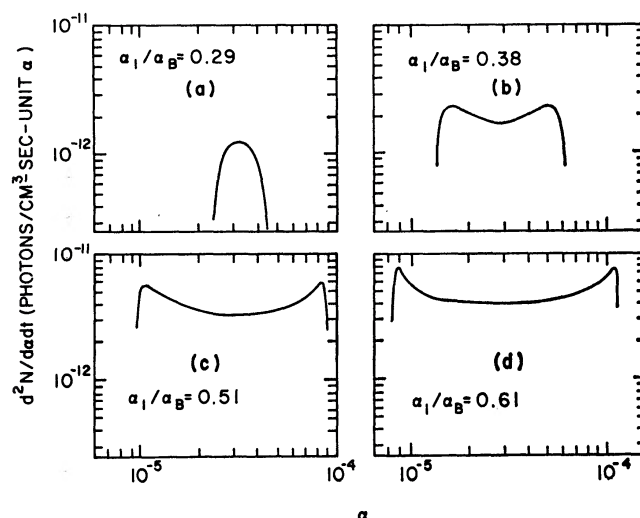


FIG. 5.—Spectra of gyrosynchrotron photons scattered by the inverse Compton process. The energy and pitch angle of the electrons is $\gamma = 10$ and $\phi = 45^\circ$. α_B and α_p are 10^{-6} and 3×10^{-6} , respectively. Spectra are shown for four values of α_1 .

The substitution of the limits and the evaluation of equation (15) were done numerically. The results are shown in Figure 5 for several values of $\alpha_1 < \alpha_p$. The inverse Compton spectrum is completely suppressed when $\alpha_1 \lesssim 0.28\alpha_p$. We have plotted spectra with $\alpha_1/\alpha_B = 0.29, 0.38, 0.51$, and 0.61 in Figure 5a, 5b, 5c, and 5d, respectively. In order to facilitate comparison with Figures 2 and 4 we chose $\alpha_B = 10^{-6}$. Again, $\alpha_p/\alpha_B = 3$, $\gamma = 10$, and $\phi = 45^\circ$. As in Figure 2, the spectra are cut off at $\alpha \simeq 2\gamma^2\alpha_1$, and the average scattered photon energy is $\sim \gamma^2\alpha_1$.

The spectra in Figure 5 are suppressed at low frequencies, reminiscent of the suppression of synchrotron radiation in the Razin effect. In addition, if we were now to compute the integral in equation (5), we would find that the total cross-section is smaller than σ_T , decreasing with decreasing α_1/α_B . In the problems of vacuum and zero magnetic field, the density of the incident photons, $N_1(\alpha_1)$, increases with decreasing α_1 , so that the total cross-section remains equal to σ_T . However, the presence of the magnetic field changes the frequency dependence of $N_1(\alpha_1)$, so that $N_1(\alpha_1, \phi)$ decreases with decreasing α_1 (see Fig. 4). Therefore, the total cross-section is suppressed in a magnetoactive plasma.

The suppression of both the total cross-section and low-frequency scattered-photon spectrum becomes large as $\gamma\alpha_1/\alpha_p$ becomes small. Similarly, the Razin effect is strongest when $\frac{3}{2}(\gamma_B/\alpha_p)$ is small (Ramaty 1968). Thus the intimate physical relationship between inverse Compton scattering and synchrotron radiation extends to the behavior of the photon spectra emitted in a dielectric medium.

IV. SUMMARY

We have presented a theory of inverse Compton scattering in the presence of a cold, collisionless plasma and have compared the results with the Razin effect. In a plasma with no magnetic field, the frequency of the incident photon is above the plasma frequency, analogous to having the cyclotron frequency be above the plasma frequency in the Razin effect—the situation in which the Razin suppression is small. In this case one expects the spectrum of Compton-scattered photons to be changed only slightly from the vacuum spectrum. As the frequency of the incident photon approaches the plasma frequency, the rate of electron energy loss becomes about three-quarters of the vacuum result.

The frequency spectrum of the scattered photons is not greatly changed by the presence of the plasma when $B = 0$. However, the maximum photon energy is only half the vacuum value when the frequency of the incident photon is close to the plasma frequency because the group velocity of the incident photon approaches zero as its frequency approaches the plasma frequency. Then the scattered-photon spectrum below the incident frequency is completely suppressed.

In a magnetoactive plasma, electromagnetic radiation can propagate below the plasma frequency. We have used the gyrosynchrotron photons as an example of an incident-photon flux below the plasma frequency to investigate the effects of a magnetoactive plasma on the inverse Compton process. We let monoenergetic electrons of a given pitch angle produce the ambient photons, and then let these same electrons scatter the photons. When the incident photons are propagating in the ordinary mode, below the plasma frequency, the spectrum of photons scattered by the inverse Compton process is greatly modified. For example, when $\gamma = 10$ and $\theta = 45^\circ$, the spectrum is completely suppressed for $\alpha_1/\alpha_B \lesssim 0.28$. For higher incident frequencies, but with $\gamma\alpha_1/\alpha_p$ small, the spectrum is cut off near the plasma frequency, as it is in the Razin effect when $\gamma\alpha_B/\alpha_p$ is small.

One of us (M. L. G.) would like to thank Drs. Reuven Ramaty and Frank C. Jones for many fruitful and stimulating discussions.

REFERENCES

- Felten, J. E., and Morrison, P. 1966, *A p. J.*, **146**, 686.
 Ginzburg, V. L., and Syrovatskii, S. I. 1964, *The Origin of Cosmic Rays* (Oxford: Pergamon Press).
 Goldstein, M. L. 1969, unpublished doctoral thesis, University of Maryland, Tech. Rept. 70-059.
 Jones, F. C. 1965, *Phys. Rev.*, **137**, B1306.
 ———. 1968, *ibid.*, **167**, 1159.
 Landau, L., and Lifshitz, E. 1960, *Electrodynamics of Continuous Media* (Reading, Mass.: Addison-Wesley Publishing Co.).
 Liemohn, H. B. 1965, *Radio Sci.*, **69D**, 741.
 Pechacek, R. E. 1966, unpublished doctoral thesis, University of California, Berkeley.
 Pechacek, R. E., and Trivelpiece, A. W. 1967, *Phys. Fluids*, **10**, 1688.
 Ramaty, R. 1968, *J. Geophys. Res.*, **73**, 3573.
 ———. 1969a, *A p. J.*, **158**, 753.
 ———. 1969b, *A p. Letters*, **4**, 43.
 Razin, V. A. 1957, unpublished dissertation, Gorkii State University.
 ———. 1960, *Radio Phys.*, **3**, 73.

