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Plasma Physics

ABSORPTION OF INTENSE GAUSSIAN ELECTROMAGNETIC RADIATION IN A MAGNETOPLASMA VIA INVERSE BREMSSTRAHLUNG

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A self-focused Gaussian laser beam propagating in the direction of static magnetic field in a collisional magnetoplasma has been considered. The explicit expressions for the time averaged power absorbed for left and right handed circularly polarised waves have been obtained. The effect of magnetic field and focusing on the rate of powder absorbed is studied in weak field regime ($\nu_E << \nu_{Th}$) as well as strong field regime ($\nu_E >> \nu_{Th}$). It is observed that the rate of absorption of right circularly polarised radiation decreases whereas for left circularly polarised radiation increases in weak field regime whereas reverse holds good in strong field regime. It is also observed that the increase in magnetic field decreases the focussing length for right circularly polarised wave whereas defocusses the left circularly polarised wave and power absorption for both the modes shows a typical behaviour due to coupling between them.

INTRODUCTION

THE study of laser induced fusion will have greater importance in future owing to its immense practical utility in the generation of controlled thermonuclear power (Engelhardt *et al.*, 1972; Dawson, 1964; and Basov & Krophim, 1964). The laser energy is deposited in the plasma *via* processes like resonance absorption (Woo & de Groot, 1978) and plasma instabilities (Kaw & Dawson, 1969). It has also been emphasised that the absorption of energy in a real plasma also occurs *via* inverse bremsstrahlung (Johnston & Dawson, 1973). In IB process a plasma electron gains energy from the electromagnetic field of the laser beam by absorbing a laser photon during collision with a nucleus. The effect of this classical absorption in the underdense plasma is to dissipate laser energy at subcritical densities. (Claire E. Max *et al.*, 1980). Silin (1965) has

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developed a nonlinear theory of the conductivity of fully ionised plasma in a strong high frequency field. Seely (1974) has studied the laser heating of plasma electrons in a uniform static magnetic field *via* inverse bremsstrahlung process using quantum mechanical approach and has found that the multiphoton absorption coefficient decreases at laser frequencies close to electron cyclotron frequencies. A similar problem has also been studied by Ford and Connell (1980) using Drude model resulting in the dependence of the absorption rate on the relative orientation of the propagation vector of the laser beam in a magnetoplasma. The beam considered by Ford and Connell has uniform intensity distribution along the wavefront whereas in the present paper we have considered a Gaussian beam propagating parallel to the external static magnetic field.

A Gaussian beam with radially nonuniform intensity distribution modifies the plasma across the transverse cross-section which leads to the self-contraction of the lateral dimension of the laser beam in the plasma and the beam self focusses as it propagates (Akhmanov *et al.*, 1968; and Sodha *et al.*, 1976). In collisional plasma, the laser self-focusing occurs on account of the carrier redistribution caused by the nonuniform heating of electrons (Sodha *et al.*, 1976) whereas in a collisionless plasma it occurs due to the ponderomotive force (Hora, 1977). Here we consider both the modes, the right and left handed circularly polarised waves. The magnetic field affects the focusing of these modes in the collisional magnetoplasma. It is observed that the effect of the magnetic field on the absorption of right and left circularly polarised radiation is different in a different field regime. The increase in magnetic field focuses the right circularly polarised beam earlier and increases the power absorption whereas reverse is true for left circularly polarized wave.

In §2, we have set up the equations for the current density and derived an expression for the time average power absorbed per unit volume for a Gaussian beam. The self-focusing of a laser beam in a magnetoplasma is discussed in §3. The effect of self-focusing on power absorbed *via* IB in the weak as well as strong field regime has been given in §4. A brief discussion of results has been presented in §5.

2. General Formalism

We consider a Gaussian laser beam propagating along the Z-axis parallel to the direction of the static magnetic field ($\mathbf{B} = B\hat{Z}$) in a magnetoplasma. The electric vector is represented by

$$\mathbf{E} = \mathrm{Re}\, \mathbf{E}_0 e^{-i w t},$$

w being the angular frequency of incident EM wave.

The modes of propagation are given by

$$E_{\pm} = E_x \pm iE_y, \qquad \ldots(1)$$

where E_x and E_y are components of the electric field vector along x and y axes; E_+ and E_- represent the right circularly and left circularly polarised waves and expressed as

$$E_{\pm} \cdot E^{*}_{\pm} = \frac{E^{2}_{\pm 0}}{f^{2}_{\pm}(z)} \exp\left(-r^{2}/r_{0}^{2} f^{2}_{\pm}(z)\right) \qquad \ldots(2)$$

Here $E_{\pm 0}$ and r_0 are axial amplitude and initial beam width respectively and $r = \hat{i}x + \hat{j}y$. $f_{\pm}$ is dimensionless beam width parameter. The momentum balance equation is

$$m \frac{d\mathbf{v}}{dt} = -e\mathbf{E} - \frac{e}{c}(\mathbf{v} \times \mathbf{B}) - m\nu\mathbf{v} \qquad \ldots(3)$$

where m, e and $\mathbf{v}$ are mass, charge and velocity of the electron respectively. ν is the electron-ion collision frequency, and considering the ions to be fixed, it is given by

$$\nu = N \left\langle \int d\Omega v (1 - \cos \chi) \frac{dQ}{d\Omega} \right\rangle. \qquad \ldots(4)$$

Here N stands for number density of ions and $\frac{dQ}{d\Omega}$ is the differential cross-section for scattering of electrons moving with velocity v through angle χ. The cross-section for non-relativistic electrons is given by

$$\frac{dQ}{d\Omega} = \left[\frac{Ze^2}{2mv^2 \cdot \sin^2 \frac{\chi}{2}}\right]^2, \qquad \ldots(5)$$

where Z = ionic charge.

Using eqn. (5) in eqn. (4), we obtain

$$\nu(v) = 2\pi N \left(\frac{Ze^2}{m}\right)^2 \left\langle \frac{L}{v^3} \right\rangle \qquad \ldots(6)$$

where L is coulomb logarithem and given by

$$L = -2 \ln \sin \frac{\chi_{min}}{2}$$

The angle χ_{min} is the minimum scattering angle corresponding to maximum impact parameter. Taking maximum impact parameter to be the Debye length

$$\sin \frac{\chi_{min}}{2} \simeq \frac{Ze^2}{mv^2}\left[\frac{4\pi ne^2}{kT}\right]^{1/2} \qquad \ldots(7)$$

As L is a weakly dependent function of velocity, ν varies as v^{-3}. The velocity consists of two contributions, one from the random thermal motion and the other from the directed motion due to the electric field. Thus

$$\mathbf{v} = \mathbf{v}_{Th} + \mathbf{v}_{D} \qquad \ldots(8)$$

In the case of nondegenerate magnetoplasma, the thermal velocity has a root mean square value given by

$$v_{Th} = \left(\frac{3kT}{m}\right)^{1/2} \qquad ...(9)$$

and for circularly polarised radiation propagating in the direction of the static magnetic field, the eqn. (8) gives the ensamble averaged value (cf. Ford & Connel, 1980)

$$\langle v_{\pm}^2 \rangle = v_{Th}^2 + v_E^2 \left(\frac{\omega}{\omega \pm \omega_c}\right)^2, \qquad ...(10)$$

where $\omega_c = \frac{eB}{mc}$ = the electron cyclotron frequency and

$$v_E = \frac{1}{\sqrt{2}} \frac{eE}{m\omega} = \text{velocity of oscillations of electrons in the high frequency field.} \qquad ...(11)$$

In the weak field regime, $v_E \ll v_{Th}$, so that the effective electron-ion collision frequency is essentially proportional to v_{Th}^{-3} and is given by

$$\nu_T = \frac{2\pi NLZ^2e^4}{m^{1/2}(3kT)^{3/2}} \qquad ...(12)$$

whereas in strong field regime, $v_E \gg v_{Th}$, so that effective elctron-ion collision frequency is proportional to v_D^{-3} and is given by

$$\nu_{\pm s} = \frac{2^{5/2}\pi z^2 emNL(\omega \pm \omega_c)^3}{E^3}$$

$$= \nu_E \frac{(\omega \pm \omega_c)^3}{\omega^3},$$

where

$$\nu_E = \frac{2^{5/2}\pi z^2 emNL\omega^3}{E^3}. \qquad ...(13)$$

The current density is defined as

$$\mathbf{J} = -n_0 e\mathbf{v}, \qquad ...(14)$$

n_0 being the number density of electrons.

In a collisional magnetoplasma, the main source of field dependence of effective dielectric constant is the redistribution of carriers resulting from inhomogeneous heating. The energy balance equation in the presence of magnetic field results in the expression for electron temperature (T_e) (Sodha *et al.*, 1974).

$$\frac{T_e - T_0}{T_0} = \tfrac{1}{2}\,\alpha\omega^2 \left[\frac{E_+E_+^*}{(\omega + \omega_c)^2 + \nu^2} + \frac{E_-E_-^*}{(\omega - \omega_c)^2 + \nu^2}\right], \qquad ...(15)$$

where T_0 = equilibrium temperature

and $\alpha = \dfrac{e^2M}{6kT_0\omega^2m^2}$.

Here M refers to ion mass, k the Boltzmann constant. It may be noted in eqn. (15) that heating of electrons on account of two modes is additive. Due to the nonuniform distribution of electron temperature, the electron concentration is given by (Sodha, *et al.*, 1974).

$$n = n_0 \Big/ \left[1 + \alpha\omega^2\left(\frac{E_+E_+^*}{(\omega + \omega_c)^2 + \nu^2} + \frac{E_-E_-^*}{(\omega - \omega_c)^2 + \nu^2}\right)\right] \qquad ...(16)$$

Using eqns. (3) and (16) in (14), we obtain

$$J_x \pm iJ_y = \sigma_\pm\,[E_x + iE_y]$$

and
$$J_z = \sigma_{zz}E_z, \qquad ...(17)$$

where

$$\sigma_\pm = \frac{\omega_p^2\,[\nu + i(\omega \pm \omega_c)]}{4\pi(\omega \pm \omega_c)^2} \times\left[1 - \frac{\alpha\omega^2}{4}\left(\frac{E_+E_+^*}{(\omega + \omega_c)^2 + \nu^2} + \frac{E_-E_-^*}{(\omega - \omega_c)^2 + \nu^2}\right)\right] \qquad ...(18)$$

and

$$\sigma_{zz} = \frac{\omega_p^2(\nu + i\omega)}{4\pi\omega^2}\left[1 - \frac{\alpha\omega^2}{4}\left(\frac{E_+E_+^*}{(\omega + \omega_c)^2 + \nu^2} + \frac{E_-E_-^*}{(\omega - \omega_c)^2 + \nu^2}\right)\right] \qquad ...(19)$$

where $\omega_P = \left(\dfrac{4\pi ne^2}{m}\right)^{1/2}$ = Plasma frequency and σ stands for the conductivity (complex) of the plasma and is a tensor of second rank.

The time averaged power absorbed per unit volume is

$$P = \langle \mathbf{J}(t)\,.\,\mathbf{E}(t) \rangle = \tfrac{1}{2}\,\text{Real}\,(\mathbf{J}\,.\,\mathbf{E}^*) \qquad ...(20)$$

Using eqns. (2) and (17), we obtain

$$P = \frac{\omega_p^2\,\nu}{8\pi}\left[\frac{E_{+0}^2}{(\omega + \omega_c)^2 + \nu^2}\,\frac{\exp\left(-r^2/r_0^2 f_+^2\right)}{f_+^2} \times\left\{\tfrac{1}{2}\left(1 - \frac{\alpha\omega^2}{4}\,\frac{E_{+0}^2\exp\left(-r^2/r_0^2 f_+^2\right)}{(\omega + \omega_c)^2 f_+^2}\right)\right\}\right.$$

$$+\frac{E_{-0}^2}{(\omega-\omega_c)^2+\nu^2}\frac{\exp(-r^2/r_0^2 f_-^2)}{f_-^2}$$

$$\times\left\{\tfrac{1}{2}\left(1-\frac{\alpha\omega^2}{4}\frac{E_{-0}^2\exp(-r^2/r_0^2 f_-^2)}{(\omega-\omega_c)^2 f_-^2}\right)\right\}$$

$$+\frac{|E_z|^2}{\omega^2+\nu^2}\left\{1-\frac{\alpha\omega^2}{4}\left(\frac{E_{+0}^2}{(\omega+\omega_c)^2}\cdot\frac{\exp(-r^2/r_0^2 f_+^2)}{f_+^2}\right.\right.$$

$$\left.\left.+\frac{E_{-0}^2}{(\omega-\omega_c)^2}\cdot\frac{\exp(-r^2/r_0^2 f_-^2)}{f_-^2}\right)\right\}$$

$$\left.+\frac{\alpha\omega^2}{4}\left(\frac{E_{+0}^2}{(\omega+\omega_c)^2}\cdot\frac{\exp(-r^2/r_0^2 f_+^2)}{f_+^2}\right)\left(\frac{E_{-0}^2}{(\omega-\omega_c)^2}\cdot\frac{\exp(-r^2/r_0^2 f_-^2)}{f_-^2}\right)\right] \quad ...(21)$$

Under the conditions, $Z = 0$, $f_\pm = 1$, $\frac{df_\pm}{dz} = 0$ and in absence of magnetic field the above expression for power reduces to the usual expression

$$P = \frac{\omega_p^2\,\nu}{8\pi}\,\frac{E_0^2}{\omega^2}$$

SELF-FOCUSING OF WAVES

Having calculated the time average power absorption per unit volume, we now solve the equation governing the propagation of right/left circulary polarised waves in the plasma. The equation governing the propagation of any mode satisfy the equation (Vlasov *et al.*, 1971) :

$$i\frac{\partial E}{\partial z} = \frac{1}{2k_0}\left(\frac{\partial^2}{\partial x^2}+\frac{\partial^2}{\partial y^2}\right)E + P(EE^*)\,\mathbf{E}, \quad ...(22)$$

where

$$P(EE^*) = \frac{k_0}{2\epsilon_0}(\epsilon-\epsilon_0) \quad ...(23)$$

k_0 being the propagation vector; ϵ and ϵ_0 for the two modes under consideration are given by

$$\epsilon_\pm = \epsilon_{\pm 0} - i\epsilon_{i\pm} + \epsilon_{2\pm}E_\pm E_\pm^*. \quad ...(24)$$

Here

$$\epsilon_{\pm 0} = 1 - \frac{\omega_p^2/\omega^2}{1 \mp \omega_c/\omega}$$

$$\epsilon_{i\pm} = \frac{\nu}{\omega} \cdot \frac{\omega_p^2/\omega^2}{(1 \mp \omega_c/\omega)^2}$$

$$\epsilon_{2\pm} = \frac{\omega_p^2 \alpha}{4\left(1 \mp \frac{\omega_c}{\omega}\right)^3}$$

The dimensionless beam width parameter f of the mode satisfies the equation (Sodha *et al.*, 1979)

$$\frac{d^2 f}{d\eta^2} + \frac{1}{f}\left(\frac{df}{d\eta}\right)^2 = \frac{2k_0^2}{\pi E_0^2 f}[I_2 - \int Q(|E|^2)\, dx\, dy], \quad \ldots(25)$$

where $\eta = \frac{z}{k_0 r_0^2}$ is the dimensionless distance of propagation

$$I_2 = \iint \left(\frac{1}{2k_0}|\nabla_\perp E|^2 - F\right) dx\, dy \quad \ldots(26)$$

$$F(\xi) = \frac{1}{k_0} \times \int_0^\xi P(\xi)\, d\xi$$

$$Q(|E|^2) = \left(\frac{EE^* P(|E|^2)}{k_0} - 2F(|E|^2)\right) \quad \ldots(27)$$

The quasi-optic equation in absence of absorption is

$$i\frac{\partial E_\pm}{\partial z} = \frac{(1 + \epsilon_{\pm 0}/\epsilon_{0zz})}{4(\omega^2/c^2)\,\epsilon_{\pm 0}} \nabla_\perp^2 E_\pm + \frac{\epsilon_{\pm 0}}{2}\,\epsilon_{2\pm}(E_\pm E_\pm^*)\, E_\pm, \quad \ldots(28)$$

where $\epsilon_{0zz} = 1 - \omega_p^2/\omega^2$.

I_2 and Q defined by eqns. (26) and (27) respectively reduces to

$$I_2 = \frac{(1 + \epsilon_{\pm 0}/\epsilon_{0zz})\, \pi c^2 E_{\pm 0}^2}{8\omega^2 \epsilon_{\pm 0}} \left[\left(1 + \frac{\epsilon_{\pm 0}}{\epsilon_{0zz}}\right) - \frac{(\omega_p r_0/c)^2\, \alpha E_{\pm 0}^2}{16\left(1 \mp \frac{\omega_c}{\omega}\right)^3}\right] \quad \ldots(29)$$

and

$$Q(|E_\pm|^2) = 0 \quad \ldots(30)$$

Thus eqn. (25) giving the variation of f reduces to

$$\frac{d^2 f_\pm}{d\eta^2} + \frac{1}{f_\pm}\left(\frac{df_\pm}{d\eta}\right)^2 = \frac{(1 + \epsilon_{\pm 0}/\epsilon_{0zz})}{4f_\pm} \times \left[\left(1 + \frac{\epsilon_{\pm 0}}{\epsilon_{0zz}}\right) - \left(\frac{\omega_p r_0}{c}\right)^2 \frac{\alpha E_{\pm 0}^2}{32\left(1 \mp \frac{\omega_c}{\omega}\right)^5}\right] \quad \ldots(31)$$

The eqn. (31), following (Sodha *et al.*, 1979) gives on integration :

$$f_{\pm}^2 = 1 + \frac{1 + \epsilon_{\pm 0}/\epsilon_{0zz}}{4}\left[\left(1 + \frac{\epsilon_{\pm 0}}{\epsilon_{0zz}}\right) - \left(\frac{\omega p r_0}{c}\right)^2 \frac{\alpha E_{\pm 0}^2}{16\left(1 \mp \frac{\omega_c}{\omega}\right)^3}\right]\eta^2 \quad ...(32)$$

Where the boundary conditions, $f_{\pm} = 1$, $\frac{df_{\pm}}{d\eta} = 0$ at $\eta = 0$ have been used in deriving eqn. (32). Equations (31) and (32) are valid in the perturbation approximation viz. $\alpha E_{\pm 0}^2 < 1$.

Effect of Self-focusing on Power Absorption in Different Regimes

The absorption of laser light in the atmosphere of pellet produces a temperature of the order of 10^8 °K ($\approx$ 10keV) for which $\nu_{Th} \approx 7 \times 10^9$cm/sec. If $E_0 \sim 1.5 \times 10^8$ Volt/cm, $\nu_E (\approx 10^8$cm/sec) is less than ν_{Th}. In such a case, we are in weak field regime. It is, however, interesting to note that keeping all parameters to be the same, we shift to strong field regime at temperatures $\approx 10^4$ °K.

The expression for power in weakfield regime ($\nu_E \ll \nu_{Th}$), from eqn. (21) becomes :

$$P_{\pm W} = \frac{\omega_p^2 \nu_T}{8\pi}\left[\frac{E_0^2}{(\omega \pm \omega_c)^2 + \nu_T^2} \cdot \frac{\exp(-r^2/r_0^2 f_{\pm}^2)}{f_{\pm}^2}\right.$$
$$\times \left\{1 - \frac{\alpha E_0^2}{2\left(1 \pm \frac{\omega_c}{\omega}\right)^2} \cdot \frac{\exp(-r^2/r_0^2 f_{\pm}^2)}{f_{\pm}^2}\right\}$$
$$- \alpha \left\{\frac{E_0^2}{\left(1 + \frac{\omega_c}{\omega}\right)^2} \cdot \frac{\exp(-r^2/r_0^2 f_{+}^2)}{f_{+}^2}\right\}$$
$$\left.\times \left\{\frac{E_0^2}{\left(1 - \frac{\omega_c}{\omega}\right)^2} \cdot \frac{\exp(-r^2/r_0^2 f_{-}^2)}{f_{-}^2}\right\}\right] \quad ...(33)$$

Using eqn. (12) we obtain :

$$P_{\pm W} = \frac{\pi Z^2 e^6 n N L}{(3mkT)^{3/2}} \frac{E_0^2}{\omega^2\left(1 \pm \frac{\omega_c}{\omega}\right)^2} \cdot \frac{\exp(-r^2/r_0^2 f_{\pm}^2)}{f_{\pm}^2}$$

$$\times \left[1 - \alpha E_0^2 \left(\frac{\exp(-r^2/r_0^2 f_\pm^2)}{2\left(1 \pm \frac{\omega_c}{\omega}\right)^2 f_\pm^2} + \frac{\exp(-r^2/r_0^2 f_\mp^2)}{\left(1 \mp \frac{\omega_c}{\omega}\right)^2 f_\mp^2} \right) \right] \quad \text{...(34)}$$

In strongfield regime ($\nu_E \gg \nu_{Th}$), the expression for power absorbed from eqn. (21) comes out to be

$$P_{\pm s} = \frac{\omega_p^2 \nu_{\pm s}}{8\pi} \left[\frac{E_0^2}{(\omega \pm \omega_c)^2 + \nu_{\pm s}^2} \cdot \frac{\exp(-r^2/r_0^2 f_\pm^2)}{f_\pm^2} \right.$$

$$\times \left\{ 1 - \frac{\alpha E_0^2}{2\left(1 \pm \frac{\omega_c}{\omega}\right)^2} \frac{\exp(-r^2/r_0^2 f_\pm^2)}{f_\pm^2} \right\}$$

$$- \alpha \left\{ \frac{E_0^2}{\left(1 + \frac{\omega_c}{\omega}\right)^2} \cdot \frac{\exp(-r^2/r_0^2 f_+^2)}{f_+^2} \right\}$$

$$\left. \times \left\{ \frac{E_0^2}{\left(1 + \frac{\omega_c}{\omega}\right)^2} \frac{\exp(-r^2/r_0^2 f_-^2)}{f_-^2} \right\} \right] \quad \text{...(35)}$$

Using eqn. (13), the above expression for power can be expressed as

$$P_{\pm s} = \frac{2^{3/2}\pi z^2 e^3 n N L \omega (1 \pm \omega_c/\omega)}{E_0} \cdot \frac{\exp(-r^2/r_2^0 f_\pm^2)}{f_\pm^2}$$

$$\times \left[1 - \alpha E_0^2 \left(\frac{\exp(-r^2/r_0^2 f_\pm^2)}{2\left(1 \pm \frac{\omega_c}{\omega}\right)^2 f_\pm^2} + \frac{\exp(-r^2/r_0^2 f_\pm^2)}{\left(1 \mp \frac{\omega_c}{\omega}\right)^2 f_\mp^2} \right) \right] \quad \text{...(36)}$$

Equations (34) and (36) have been used in computing the power absorbed in weak and strong field regimes respectively.

Discussion of Results

The following set of parameters has been used for calculations :

$n = N = 10^{19}$

$Z = 1$ (Singly ionised plasma)

$E_0 = 1.5 \times 10^8$ volt/cm

$W = 1.8 \times 10^{15}$ rad/sec

$$\alpha E_0^2 = 0.43$$

$$\frac{\omega_c}{\omega} = 0.01 \text{ and } 0.1$$

$$T = 10^8 \text{ °K (Weak field regime)}$$

$$T = 10^4 \text{ °K (Strong field regime)}$$

Keeping all parameters the same, we have studied the absorption in weak and strong field regimes by assigning two different values of temperature. When temperature is of the order of 10^8 °K, the condition $\nu_{Th} \gg \nu_E$ is satisfied and so we are in weak field regime. For temperature of the order of 10^4 °K, $\nu_E \gg \nu_{Th}$ and we are in strong field regime. The variation of beam width parameter f of a laser beam with η in a collisional magnetoplasma for the right and left circularly polarised wave for a typical set of parameters $\left(\omega_p^2/\omega^2 = 0.1, \frac{\omega_p r_0}{c} = 11.22 \text{ and } \frac{\omega_c}{\omega} = 0.01, 0.1\right)$ is shown in Fig. 1. It is observed that the focusing of both the modes is affected by the magnetic field and enhancement of focusing for right circularly polarised wave at the higher value of static magnetic field is obvious. Fig. 2 shows the variation of power absorbed per second per unit volume with η for right and left circularly polarised waves at different values of ω_c/ω in the weak field

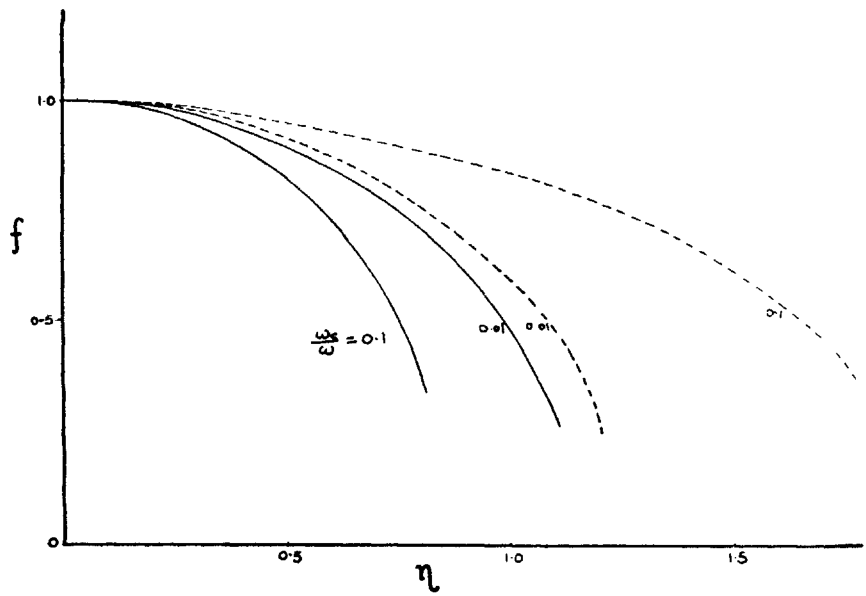

FIG. 1. Variation of beam width parameter f of a laser beam with η in a collisional magnetoplasma when $\omega_p^2/\omega^2 = 0.1$, $\frac{\omega_p r_0}{c} = 11.22$ and $\frac{\omega_c}{\omega} = 0.01, 0.1$ for the right (———) and left (– – – –) circularly polarised wave.

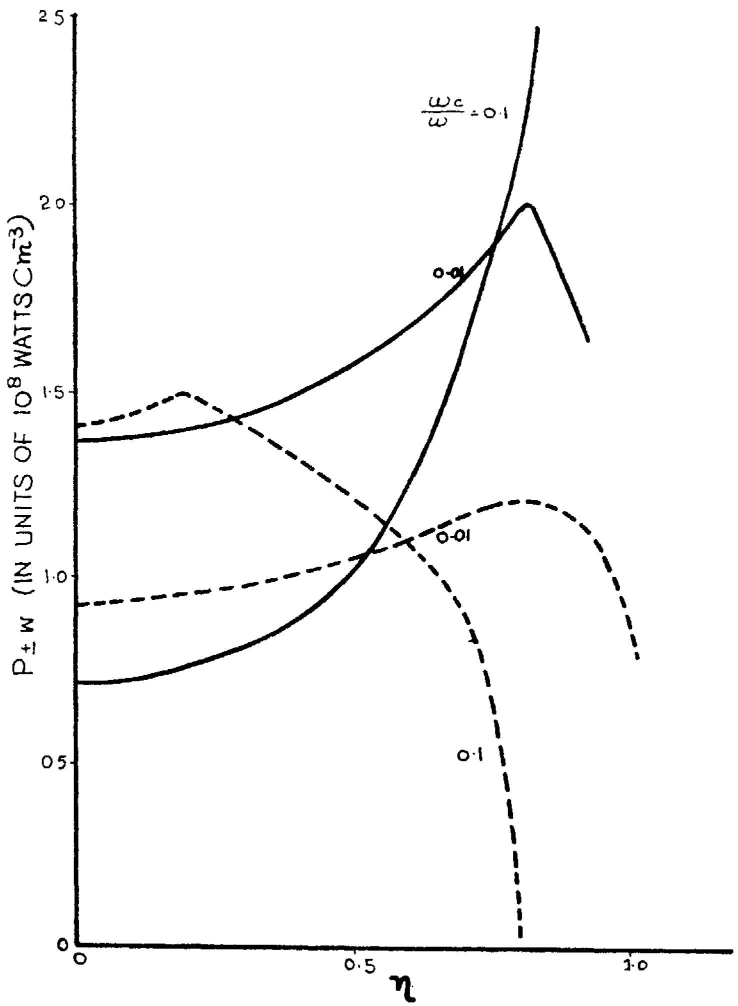

FIG. 2. Variation of power absorbed per second per unit volume $P_{\pm\omega}$ with η at $\frac{\omega_c}{\omega}$ = 0. 01 and 0.1 for the right (———) and left (- - - -) circularly polarised wave in weak field regime.

regime. In strong field regime, the variation of rate of absorption of power per unit volume with η for both the modes at different values of ω_c/ω is given in Fig. 3. In weak field regime, the absorption of right circularly polarised wave decreases whereas that of left circularly polarised wave increases with magnetic field. On the other hand, the absorption of left circularly polarised wave decreases whereas for right circularly polarised wave increases with magnetic field in strong field regime. In this regard, our results are in good agreement with the results obtained by Ford & Connel (1980). It is also observed that right circularly polarised wave focuses earlier at higher values of magnetic field, thereby increasing the power absorbed in

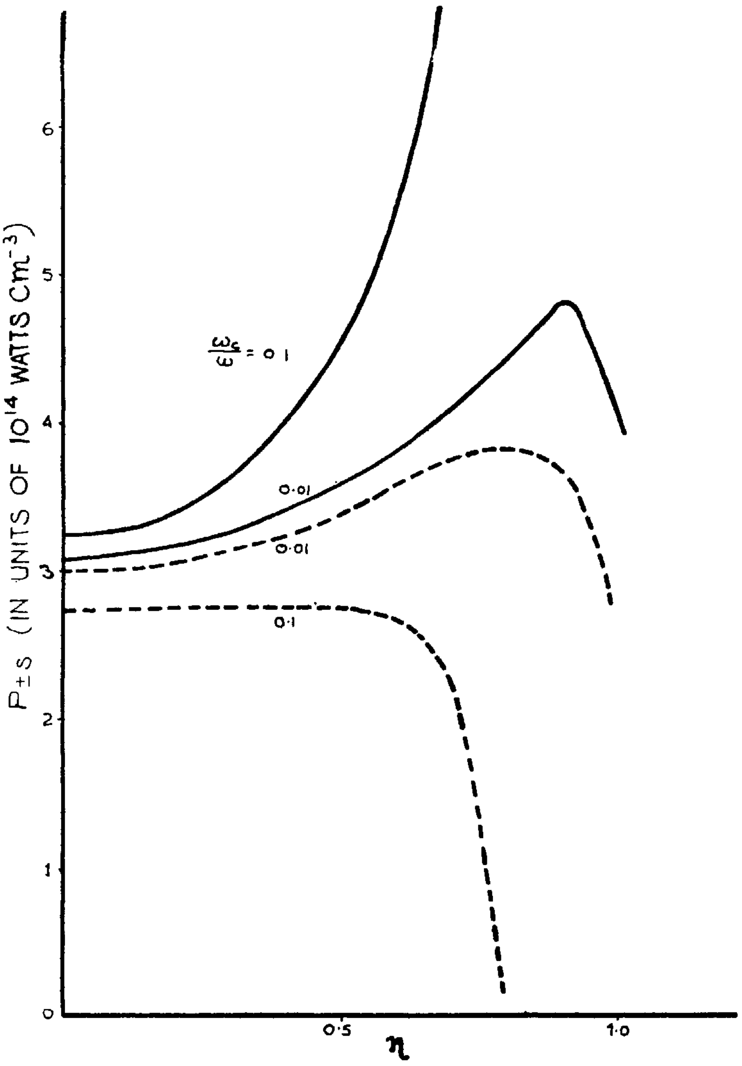

FIG. 3. Variation of power absorbed per second per unit volume $P_{\pm s}$ with η at $\frac{\omega_c}{\omega} = 0.01$ and 0.1 for right (———) and left (- - - -) circularly polarised wave in strong field regime.

both the field regime. But at low values of static magnetic field the wave first focuses and then defocuses resulting decrease in power absorbed. The rate of power absorbed for left circularly polarised wave as it propagates in plasma first increases due to focusing and then decreases due to its defocusing in weak field regime. Higher the magnetic field, more is the defocusing and consequently power absorption falls. But in a strong field regime, the defocusing may only take place at higher values of magnetic field for left circularly polarised radiation. This typical variation in power absorbed for both the modes may be understood by coupling between the right and left handed polarised modes.

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