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EFFECT OF RADIATION PROCESSES ON TRANSPORT PHENOMENA IN A PLASMA IN A STRONG MAGNETIC FIELD

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It is shown that radiation and absorption of electromagnetic waves by electrons in a plasma in a strong magnetic field not only affect electron relaxation but can also have a profound effect on transport phenomena. For these effects to occur in stationary fields an additional mechanism capable of changing the photon distribution must operate in addition to the radiative collisions. If the plasma contains "trapped" radiation this mechanism is the reflection of electromagnetic waves at a frequency approximately $\omega_H = eH/mc$ (m is the electron mass and H is the magnetic field) from mirrors surrounding the plasma. If the transverse dimension of a cylindrical mirror (with axis parallel to H) is appreciably smaller than the mean free path of photons of frequency ω_H the transverse conductivity of the plasma (as well as the transverse thermal conductivity) resulting from radiation processes is independent of mirror dimensions and is determined by the effective time between electron radiative collisions τ_{eff} (14). When $\tau_{\text{eff}} \ll \tau_e^{(s)}$, where $\tau_e^{(s)}$ is the mean Coulomb relaxation time, the transverse conductivity of the plasma is determined by radiative collisions rather than Coulomb collisions. The effects become comparable at $n \sim 10^{14}\ \text{cm}^{-3}$, $H \sim 5 \times 10^7$ G and $T \sim 3 \times 10^{-3}\ mc^2$.

In alternating fields the radiation processes can affect transport phenomena in the absence of mirrors.

1. The present authors have shown [1] that radiation and absorption of electromagnetic waves by electrons in a plasma in a strong magnetic field can have an important influence on the establishment of the equilibrium electron distribution. At nonrelativistic temperatures $T \ll mc^2$ (m is the electron mass) the electrons interact primarily with photons in a frequency range $\Delta \sim \omega_H\sqrt{T/mc^2}$ about the resonance frequency $\omega_H = eH/mc$ (H is the magnetic field). For these photons, absorption and emission leads to the rapid establishment, in a time $\tau_p \sim \omega_H\Omega^{-2}\sqrt{T/mc^2}$ (where Ω is the plasma frequency), of a quasi-equilibrium Rayleigh–Jeans distribution with temperature equal to the mean energy of the transverse (with respect to **H**) electron motion.[1)] The change in the electron distribution caused by radiative processes is much slower; this change occurs in a time $\tau_e^{(r)}$, which is of the order of the ratio of mean electron energy to the mean radiation intensity $\tau_e^{(r)} \sim c/r_0\omega_H^2$ (r_0 is the classical radius of the electron).

If $\tau_e^{(r)} \ll \tau_e^{(s)}$, where $\tau_e^{(s)}$ is the mean time between electron Coulomb collisions, then $\tau_e^{(r)}$ rather than $\tau_e^{(s)}$ determines the time required for the establishment of a Maxwellian electron distribution over transverse momentum (radiative relaxation for the longitudinal momentum component requires a time of approximately $mc^2\tau_e^{(r)}/T$).[2)]

It is of interest to examine the effect of radiation processes on transport phenomena in addition to electron and photon relaxation in a plasma in a strong magnetic field.[3)] In the present work we in-

1)Thomson scattering of photons by electrons in the magnetic field can also affect the photon relaxation process. The maximum possible size of this effect has been estimated in [1]. Actually, however, Thomson scattering can be neglected because absorption of electromagnetic waves in the plasma means that the photons are not strictly monochromatic so that the Thomson scattering is no longer a resonance process.

2)When $\tau_e^{(r)} \ll \tau_e^{(s)}$, a Rayleigh-Jeans equilibrium distribution in the frequency range Δ close to the resonance frequency ω_H characterized by a constant temperature is established (in a time $\sim\tau_p$) rather than a quasi-equilibrium distribution.

3)Radiation effects only influence the electrical and thermal conductivity of the plasma and can be neglected in the viscosity since plasma viscosity is determined primarily by the ions.

vestigate this question for the nonrelativistic case.

In determining the transport coefficients in a plasma in a strong magnetic field one usually uses the kinetic equation in the particle velocity space.[2-5] However, this method is not suitable for analysis of radiative collisions because the radiation and absorption of a photon cannot be uniquely related to a given polar angle φ in electron velocity space (the ambiguity in φ in radiative processes is of order 2π).[4)] To analyze the effect of radiation processes on transport phenomena in a plasma we do not use a kinetic equation in velocity space, but rather in a space comprising those variables that are used for the quantum-mechanical description of the motion of an electron in a magnetic field. These familiar variables are the quantum number n, which describes the motion in the plane perpendicular to the field **H** (the xy plane); p_z, the projection of the momentum in the direction of **H**; and ξ, the coordinate of the guiding center of the Larmor orbit along the x axis. The same variables are used to describe the motion of an electron in crossed electric and magnetic fields. If there is an electric field **E** along the x axis the energy of an electron in a state characterized by $\kappa \equiv n, p_z, \xi$ is

$$\mathscr{E}_\varkappa = \hbar\omega_H(n + \tfrac{1}{2}) + p_z^2/2m - eE\xi.$$

2. We assume that the electric field **E** and the gradients of density n and temperature T are perpendicular to **H** (along the x axis). In this case, if one considers single radiative processes only, the kinetic equations describing the electron distribution function f_κ and photon distribution function $N_{\mathbf{k}}(x)$ are

$$\frac{\partial f_\varkappa}{\partial t} = \dot{f}_\varkappa^{(r)} \equiv \sum_{\varkappa'\mathbf{k}} W(\varkappa\mathbf{k};\varkappa')\{f_{\varkappa'}[1 + N_{\mathbf{k}}(\xi)] - f_\varkappa N_{\mathbf{k}}(\xi)\}$$

$$+ \sum_{\varkappa'\mathbf{k}} W(\varkappa'\mathbf{k};\varkappa)\{f_{\varkappa'}N_{\mathbf{k}}(\xi') - f_\varkappa[1 + N_{\mathbf{k}}(\xi')]\}, \quad (1)$$

$$\frac{\partial N_{\mathbf{k}}(x)}{\partial t} + c\frac{k_x}{k}\frac{\partial N_{\mathbf{k}}(x)}{\partial x} = \dot{N}_k^{(r)}$$

$$= \sum_{\varkappa\varkappa'} W(\varkappa\mathbf{k};\varkappa')\{f_{\varkappa'}[1 + N_{\mathbf{k}}(x)] - f_\varkappa N_{\mathbf{k}}(x)\}_{\xi=x}, \quad (2)$$

where $W(\kappa\mathbf{k};\kappa')$ is the probability for transition of an electron from state κ' to state κ with the radiation of a photon characterized by a wave vector **k** (we do not distinguish between different photon polarizations).

We note that (1) does not contain kinematic terms. The electric field only appears implicitly in (1), i.e., in the probability expression $W(\kappa\mathbf{k};\kappa')$, which differs from zero when the following relations are satisfied:

$$\mathscr{E}_\varkappa + \hbar\omega = \mathscr{E}_{\varkappa'}, \qquad p_z + \hbar k_z = p_z', \qquad \xi + \hbar c k_y/eH = \xi'.$$

It is a simple matter to find stationary solutions of (1) and (2) for small values of E, $\partial n/\partial x$ and $\partial T/\partial x$. Under these conditions we have from (2)

$$N_{\mathbf{k}}(x) = \nu_{\mathbf{k}}\tau_{\mathbf{k}} - \frac{c\tau_{\mathbf{k}}}{\hbar\omega}\frac{k_x}{k}\frac{\partial T}{\partial x}; \quad (3)$$

$$\frac{1}{\tau_{\mathbf{k}}} = \sum_{\varkappa\varkappa'} W(\varkappa\mathbf{k};\varkappa')(f_\varkappa - f_{\varkappa'})_{\xi=x},$$

$$\nu_{\mathbf{k}} = \sum_{\varkappa\varkappa'} W(\varkappa\mathbf{k};\varkappa') f_{\varkappa'}|_{\xi=x}. \quad (3')$$

Using (3) one can easily show that to quadratic terms in E, $\partial n/\partial x$ and $\partial T/\partial x$ (1) is satisfied by the Maxwellian distribution[5)]

$$f_\varkappa = (2\pi)^{3/2} n(\xi)\hbar^3[mT(\xi)]^{-3/2}e^{-\varepsilon_\varkappa/T} \quad (4)$$

with energy $\epsilon_\kappa = \mathscr{E}_\kappa + eE\xi$, which the electron would have in the absence of an electric field.

Using (3) we find the photon distribution

$$N_{\mathbf{k}}(x) = \frac{T(x)}{\hbar\omega}\left\{1 + u_0\frac{k_y}{\omega} + v_0\frac{k_x}{\omega}\right\};$$

$$u_0 = -\frac{cE}{H}\left\{1 - \frac{T}{eEn}\frac{\partial n}{\partial x} + \frac{\tau_{\mathbf{k}}^{(0)}}{eE}\frac{\partial}{\partial T}\left(\frac{T}{\tau_{\mathbf{k}}^{(0)}}\right)\frac{\partial T}{\partial x}\right\},$$

$$v_0 = -\frac{c^2\tau_{\mathbf{k}}^{(0)}}{T}\frac{\partial T}{\partial x}, \qquad \tau_{\mathbf{k}}^{(0)} = \frac{T}{\hbar\omega}\frac{1}{\nu_{\mathbf{k}}^{(0)}}, \quad (5)$$

where $\nu_{\mathbf{k}}^{(0)}$ is given by (3') with the distribution in (4) replacing f_κ. This distribution differs from the Rayleigh-Jeans equilibrium distribution: when $\partial T/\partial x = 0$ this distribution corresponds to entrainment of the photon gas as a whole by the electrons in the plasma and the photons move with drift velocity u_0 along the y axis.

It is evident that the distributions in (4) and (5) do not yield an electric current along the electric field. The current density along the x axis is, in fact, given by the expression:[6]

$$j_x = \frac{e}{V}\sum_{\varkappa\varkappa'\mathbf{k}} W(\varkappa\mathbf{k};\varkappa')(\xi - \xi')\{f_{\varkappa'}[1 + N_{\mathbf{k}}(\xi)] - f_\varkappa N_{\mathbf{k}}(\xi)\}, \quad (6)$$

[4)]For this reason the calculation of the effect of radiation phenomena on electrical conductivity of a plasma reported in [1], which uses the usual kinetic equation, can only serve an illustrative purpose. Carrying out the calculation in this way corresponds to the assumption of an equilibrium photon distribution in the presence of an electric field; as will be shown later, this assumption gives a value that differs from the proper value by a factor T/mc^2.

[5)]This solution for the distribution function of a degenerate electron was first given by Titeica[6] in connection with the electrical conductivity of metals.

where V is the normalized volume. Taking account of the definition $\dot{N}_{\mathbf{k}}^{(r)}$ and noting that $\xi - \xi' = -\hbar c k_y/eH$ we have

$$j_x = -\frac{\hbar c}{H}\frac{1}{V}\sum_{\mathbf{k}} k_y \dot{N}_{\mathbf{k}}^{(r)}. \tag{6'}$$

In the stationary state

$$\dot{N}_{\mathbf{k}}^{(r)} = -c\frac{k_x}{k}\frac{1}{\hbar\omega}\frac{\partial T}{\partial x}$$

and the current j_x vanishes.

Thus, taking account of only single radiative collisions in the stationary case in an infinite plasma produces no electric current along the electric field. The radiative collisions can, however, have an important effect on the electrical conductivity of the plasma. This will be the case if there is additional mechanism capable of changing the photon distribution, in addition to radiative collisions.

For a plasma with trapped radiation a mechanism of this kind is photon reflection (frequencies in the range $\Delta \sim \omega_H\sqrt{T/mc^2}$ about the frequency ω_H) from mirrors surrounding the plasma and parallel to the magnetic field. The effect of the mirrors is to reduce the entrainment velocity of the photon gas by the plasma electrons owing to which the photon distribution becomes more nearly a Rayleigh-Jeans distribution.

If the plasma is bounded by a cylindrical mirror (with axis along **H**) the entrainment velocity of the photons is reduced by approximately $(L + l_{\mathbf{k}})/L$ where L is the transverse mirror dimension and $l_{\mathbf{k}}$ is the mean free path of a photon with frequency approximately ω_H with respect to the radiation process. When $L \ll l_{\mathbf{k}}$ the entrainment essentially vanishes and to terms of order $L/l_{\mathbf{k}}$ the photon distribution becomes a Rayleigh-Jeans distribution. The latter [along with the Maxwellian distribution (4) for electrons], in contrast with the total entrainment distribution (5), does not cause the current j_x to vanish. If $L \ll l_{\mathbf{k}}$, we shall show that the current j_x, and consequently the electrical conductivity, do not depend explicitly on L.

In order to explain the reduction in photon entrainment velocity caused by the mirrors we assume for simplicity that there are two mirrors perpendicular to the y axis and that these scatter photons elastically in the frequency range Δ about the frequency ω_H. Under these conditions the change in the photon distribution function caused by reflection of photons from the mirrors can be written symbolically in the form

$$\dot{N}_{\mathbf{k}}^{(s)} = (c/2L)(N_{\mathbf{k}'} - N_{\mathbf{k}}),$$

where L is the distance between mirrors and $\mathbf{k}' = (k_x, -k_y, k_z)$.

The kinetic equation for determining the photon distribution function (taking account of radiative collisions and photon reflection from the mirrors) is

$$\dot{N}_{\mathbf{k}}^{(r)} + \dot{N}_{\mathbf{k}}^{(s)} = c\frac{k_x}{k}\frac{1}{\hbar\omega}\frac{\partial T}{\partial x}.$$

Solving this equation together with the kinetic equation $\dot{f}_{\kappa}^{(r)} = 0$ for the electrons, we have

$$f_\varkappa = \frac{(2\pi)^{3/2}\hbar^3 n(\xi)}{[mT(\xi)]^{3/2}} e^{-\varepsilon_\varkappa/T}, \qquad N_{\mathbf{k}}(x) = \frac{T(x)}{\hbar\omega}\left\{1 + u\frac{k_y}{\omega} + v_0\frac{k_x}{\omega}\right\},$$

where $u = u_0 L/(L + l_{\mathbf{k}})$ is the drift velocity in the presence of the mirrors.

If we locate two more mirrors perpendicularly to the x axis (separated by a distance L) v_0 in the last formula becomes $v = v_0 L/(L + l_{\mathbf{k}})$ and the photon distribution becomes

$$N_{\mathbf{k}}(x) = \frac{T(x)}{\hbar\omega}\left\{1 + u\frac{k_y}{\omega} + v\frac{k_x}{\omega}\right\}. \tag{7}$$

3. We now compute the current density j_x produced by radiative collisions of electrons, assuming that the transverse dimension of the cylindrical mirror L is much smaller than $l_{\mathbf{k}}$. In (6) we substitute the Maxwellian distribution (4) in place of f_κ and the Rayleigh-Jeans distribution $N_{\mathbf{k}}(x) = T(x)/\hbar\omega$ in place of $N_{\mathbf{k}}$, noting that

$$\{f_{\varkappa'}[1 + N_{\mathbf{k}}(\xi)] - f_\varkappa N_{\mathbf{k}}(\xi)\}_{\xi=x} = \frac{ck_y}{eH\omega}\left(-eE + \frac{T}{n}\frac{\partial n}{\partial x} + T\frac{\partial \ln f_\varkappa}{\partial T}\frac{\partial T}{\partial x}\right) f_\varkappa, \tag{8}$$

whence

$$j_x = \sigma^{(r)}\left(E - \frac{T}{ne}\frac{\partial n}{\partial x}\right) - \Upsilon\frac{\partial T}{\partial x}; \tag{9}$$

$$\sigma^{(r)} = \frac{T}{2H^2}\int\frac{d\mathbf{k}}{(2\pi)^3}\frac{\sin^2\vartheta}{\tau_{\mathbf{k}}^{(0)}}, \qquad \Upsilon = \frac{1}{2}\frac{T}{eH^2}\frac{\partial}{\partial T}\left(T\int\frac{d\mathbf{k}}{(2\pi)^3}\frac{\sin^2\vartheta}{\tau_{\mathbf{k}}^{(0)}}\right), \tag{10}$$

$$\frac{1}{\tau_{\mathbf{k}}^{(0)}} = \frac{\hbar\omega}{T}\sum_{\varkappa\varkappa'} W(\varkappa\mathbf{k}; \varkappa') f_\varkappa \tag{10'}$$

(ϑ is the angle between **k** and **H**).

We note the following relation that derives from (10)

$$\Upsilon = (T/e)\,\partial\sigma^{(r)}/\partial T. \tag{11}$$

In the nonrelativistic case [1]

$$\frac{1}{\tau_{\mathbf{k}}^{(0)}} = \frac{\pi\sqrt{2\pi}ne^2}{m\omega}\left(\frac{mc^2}{T}\right)^{1/2}\frac{1+\cos^2\vartheta}{|\cos\vartheta|}\exp\left\{-\frac{mc^2}{2T}\frac{(\omega-\omega_H)^2}{\omega_H^2\cos^2\vartheta}\right\}. \tag{12}$$

Substituting this expression in (1) we obtain finally

$$\sigma^{(r)} = \frac{3}{10}\frac{ne^2}{m\omega_H^2\tau_e^{(r)}}\frac{T}{mc^2}, \qquad \gamma = \frac{3}{10}\frac{ne}{m\omega_H^2\tau_e^{(r)}}\frac{T}{mc^2},$$
$$\tau_e^{(r)} = \frac{3}{4}\frac{c}{r_0\omega_H^2}. \tag{13}$$

We note that these expressions [just as Eqs. (10), (11), and (12)] take account of the existence of two polarization states for the photons (they are obtained under the assumption of an equilibrium photon distribution function; $W(\kappa\mathbf{k}; \kappa')$ represents the radiation probability summed over polarizations).

It is evident that the electrical conductivity $\sigma^{(r)}$ corresponds to the effective time between collisions; this quantity is of order

$$\tau_{eff} = \tau_e^{(r)} mc^2/T. \tag{14}$$

If $\tau_{eff} \ll \tau_e^{(s)}$ the transverse conductivity of the plasma is determined by radiative processes rather than Coulomb collisions.

The order of magnitude of the ratio of the radiative conductivity $\sigma^{(r)}$ to the conductivity associated with Coulomb collisions $\sigma^{(s)}$ is

$$\frac{\sigma^{(r)}}{\sigma^{(s)}} \approx \frac{12\pi}{10\mathscr{L}}\left(\frac{\omega_H}{\Omega}\right)^2\left(\frac{T}{mc^2}\right)^{5/2},$$

where $\mathscr{L}$ is the Coulomb logarithm. This ratio approaches unity when $n \sim 10^{13}\ \mathrm{cm}^{-3}$, $T \sim 10^{-2}\ mc^2$ and $H \sim 10^7$ G. Under these conditions the mean free path of a photon of frequency ω_H is $l_{\mathbf{k}} \sim 2 \times 10^2$ cm.

Equation (9) has been obtained assuming an equilibrium photon distribution. Actually, however, it follows from (7) that the photon distribution is not an equilibrium distribution. For this reason there is a small current along the y axis; the magnitude of this current is approximately $j_y \sim j_x L/l_{\mathbf{k}}$. There is also a small correction to the current j_x (of order $j_x L/l_{\mathbf{k}}$).

Up to this point we have been considering the effect of radiation processes on transverse plasma conductivity. It is evident that in the absence of mirrors oriented perpendicularly to **H** these effects will not affect the longitudinal (with respect to **H**) plasma conductivity. In the presence of such mirrors, however, there is a longitudinal radiative conductivity; when $L \ll l_{\mathbf{k}}$ (L is the distance between mirrors) the longitudinal component is determined by the effective time between radiative collisions, just as the transverse component. If the electron motion is to be unaffected by the mirrors the condition $L \gg l_e^{(r)}$ must be satisfied ($l_e^{(r)}$ is the radiative mean free path of the electrons); it is difficult to satisfy this condition simultaneously with condition $L \ll l_{\mathbf{k}}$.

We noted above that some additional mechanism is required in order for photon scattering to have an effect on conductivity; this statement holds for stationary fields. In a nonstationary field radiative processes lead to finite conductivity even in the absence of mirrors. In this case, if ν is the frequency of the electric field the effect of nonstationary is equivalent to the presence of a mirror for which $L \sim c/\nu$.

4. We now find the heat flux along the x axis due to the electrons. This quantity is given by the following expression, which is analogous to Eq. (6) for the current j_x:

$$q_x^{(e)} = \frac{1}{V}\sum_{\kappa\kappa'\mathbf{k}} W(\kappa\mathbf{k}; \kappa')(\xi - \xi')\,\varepsilon_\kappa\{f_{\kappa'}[1 + N_{\mathbf{k}}(\xi)] - f_\kappa N_{\mathbf{k}}(\xi)\}. \tag{15}$$

Substituting (8) and using (10′) we have

$$q_x^{(e)} = -\varkappa_e^{(r)}\frac{\partial T}{\partial x} + \chi\left(E - \frac{T}{ne}\frac{\partial n}{\partial x}\right);$$
$$\varkappa_e^{(r)} = \frac{T}{2(eH^2)}\frac{\partial}{\partial T}\left[T^{1/2}\frac{\partial}{\partial T}\left(T^{5/2}\int\frac{dk}{(2\pi)^3}\frac{\sin^2\vartheta}{\tau_k^{(0)}}\right)\right],$$
$$\chi = \frac{T^{1/2}}{2eH^2}\frac{\partial}{\partial T}\left[T^{5/2}\int\frac{dk}{(2\pi)^3}\frac{\sin^2\vartheta}{\tau_{\mathbf{k}}^{(0)}}\right]. \tag{16}$$

Using (10) we can relate χ and $\kappa_e^{(r)}$ to $\sigma^{(r)}$:

$$\varkappa_e^{(r)} = \frac{T}{e^2}\frac{\partial}{\partial T}\left[T^{1/2}\frac{\partial}{\partial T}(T^{3/2}\sigma^{(r)})\right] \qquad \chi = \frac{T^{1/2}}{e}\frac{\partial}{\partial T}(T^{3/2}\sigma^{(r)}). \tag{17}$$

Further, using (13) we finally obtain for the nonrelativistic case

$$\varkappa_e^{(r)} = \frac{3}{2}\frac{nT}{m\omega_H^2\tau_e^{(r)}}\frac{T}{mc^2}, \qquad \chi = \frac{3}{4}\frac{enT}{m\omega_H^2\tau_e^{(r)}}\frac{T}{mc^2}. \tag{18}$$

The thermal conductivity $\kappa_e^{(r)}$, like the electrical conductivity $\sigma^{(r)}$, evidently corresponds to the effective time between radiative collisions, approximately τ_{eff}.

We note that the relative effect of radiative phenomena on the thermal conductivity is smaller than on electrical conductivity since the thermal conductivity, which is due to Coulomb collisions, contains the additional factor $\sqrt{M/m}$ (M is the ion mass).

Equation (18) has been obtained under the assumption of an equilibrium photon distribution. The deviation of the photon distribution function (7) from the Rayleigh-Jeans distribution leads to a small electron heat flux along the y axis; this quantity is of order $q_y^{(e)} \approx q_x^{(e)} L/l_{\mathbf{k}}$.

The expressions for electric current and heat flux due to electrons can be written in more symmetric form by means of the chemical potential

$$\mu = T \ln (nT^{-3/2}) + T \ln [(2\pi/m)^{3/2} \hbar^3]$$

and the relation

$$\gamma + \tfrac{3}{2}\sigma^{(r)}/e = \chi/T, \tag{19}$$

which follows from (11) and (17):

$$j_x = \sigma^{(r)} \left(E - \frac{T}{e}\frac{\partial}{\partial x}\left(\frac{\mu}{T}\right)\right) - \chi \frac{1}{T}\frac{\partial T}{\partial x},$$
$$q_x^{(e)} = \chi\left(E - \frac{T}{e}\frac{\partial}{\partial x}\left(\frac{\mu}{T}\right)\right) - \left(\varkappa_e^{(r)} + \frac{3}{2e}\chi\right)\frac{\partial T}{\partial x}. \tag{20}$$

We note that (19) obeys the symmetry requirements for the kinetic coefficients.

5. The heat flux $q^{(e)}$ must be supplemented by the heat flux $q^{(p)}$ due to photons with frequencies lying in the range Δ about ω_H. This heat flow is obviously

$$\mathbf{q}^{(p)} = c \int \frac{d\mathbf{k}}{(2\pi)^3}\frac{\mathbf{k}}{k}\hbar\omega N_{\mathbf{k}},$$

where the integration is carried over frequencies in the range Δ.

Replacing $N_{\mathbf{k}}$ by the expression given in (7) we have

$$q_x^{(p)} = -\varkappa_p^{(r)}\, \partial T/\partial x, \qquad \varkappa_p^{(r)} \sim Lc^{-2}\omega_H^3 \sqrt{T/mc^2} \tag{21}$$

(the ratio of the fluxes $q_y^{(p)}$ and $q_x^{(p)}$, of order $q_y^{(p)}/q_x^{(p)} \sim T/mc^2(\omega_H\tau_p)^{-1}$, is much smaller than unity). We see that the photon heat flux depends on the distance between mirrors.

We also note that the ratio of photon and electron heat fluxes

$$q_x^{(p)}/q_x^{(e)} \sim L\omega_H c^{-1} (\omega_H/\Omega)^2 (mc^2/T)^{3/2},$$

can be greater than unity.

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