If we substitute the actual values $r_1 = 3{\cdot}7 \times 10^{10}$, $L = 1{\cdot}015 \times 10^{50}$ and $m_1 = 5{\cdot}967 \times 10^{31}$ and if we take $\alpha = 1{\cdot}25$ and $M_0/M_1 = 5$, we have

$$\frac{r_0}{r_1} = 0{\cdot}20\left[1 + \frac{4{\cdot}7 \times 10^6}{c}\right]$$

which for $c \gg 4{\cdot}7 \times 10^6$ reduces to $\dfrac{r_0}{r_1} \simeq \dfrac{10^6}{c}$

If $c \simeq 10^3$ cm s^{-1}, we have $r_0 \simeq 10^3\ r_1$. That places the secondary well outside the region to be occupied by the envelope of the primary. Equation (9), then, gives for the initial secondary mass $m_0 = 0{\cdot}87 \times 10^{30}$ g.

These values for the initial separation and mass are comparable to the values for Jupiter and the Sun. If the Sun becomes a giant star and loses a considerable fraction of its mass at very low flow velocity, part of it would be collected on Jupiter, which would then move close to the Sun. Their final configuration would probably not be very different from the actual one for the binary star SW Sge.

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Bremsstrahlung Radiation in an Intense Magnetic Field and Emission from Pulsars

Simon and Strange recently argued[1] that one of the processes (bremsstrahlung in the $0 \rightarrow 0$ transition in a magnetic field) which we proposed[2] as the chief mechanism of emission from pulsars cannot give rise to a negative absorption coefficient. They based their argument on an existing formula for the absorption coefficient valid only for synchrotron radiation. Further, their expression for the emissivity was found by us to be applicable only to the high quantum number case. When the correct formula for the absorption coefficient was used in conjunction with the emissivity valid for the $0 \rightarrow 0$ transition, we found that the absorption coefficient can become negative and amplification of electromagnetic waves is possible from the $0 \rightarrow 0$ bremsstrahlung process.

First we shall show that the result of Simon and Strange, equation (3) of their paper,

$$\alpha_\omega = \frac{8\pi^3 c}{\omega^2}\ \frac{n(\theta)\cos\theta}{n_r^2} \int_{-\infty}^{\infty} f(p_z)\frac{dQ_\omega}{dp_z}\,dp_z$$

cannot be correct as it gives a zero absorption coefficient for a gas in equilibrium, independent of any absorption process. (α_ω = absorption coefficient, ω = frequency, p_z = electron momentum parallel to the field, $f(p_z)$ = electron distribution function, θ = angle of emission with respect to the field, Q_ω = rate at which energy is emitted spontaneously by the electron with momentum between p_z and $p_z + dp_z$ at the frequency ω, per $d\omega$, per unit solid angle, in the direction θ, and $n(\theta)$ and n_r are refractive indices.) Q_ω depends on the absolute value of p_z and is therefore an even function of p_z. Consequently dQ_ω/dp_z is an odd function of p_z. In equilibrium $f(p_z)$ is an even function and hence α_ω in equation (3) of ref. 1 vanishes identically, independent of any absorption process. This clearly violates basic laws of thermodynamics.

The origin of this inconsistency can be traced to Simon and Strange's improper application of equation (2.54) of ref. 3 of their paper[3]. This result is valid only for classical synchrotron radiation, derived from the master equation (2.36) after using equations (2.46) and (2.50) which are applicable only to synchrotron radiation and not to the bremsstrahlung process. Equation (2.50) describes the momentum conservation condition which is absent in the bremsstrahlung case, for the nucleus can absorb momentum without absorbing energy. The improper application of equation (2.54) leads to the result of zero absorption coefficient.

In order to derive the correct equation, we refer to the master equation (equation (2.36) of ref. 3) modified for a one-dimensional gas

$$\alpha_\omega = \frac{8\pi^3 c^2}{n_r^2 \hbar\omega^2} \int_{-\infty}^{\infty} Q_\omega(p')[f(p) - f(p')]\,dp' \qquad (1)$$

where p and p' are electron momenta before and after absorbing the photon $\hbar\omega$. The energy conservation law requires that

$$\hbar\omega + \frac{p_z^2}{2m} = \frac{p_z'^2}{2m} \qquad (2)$$

Because $\hbar\omega \ll p_z^2/2m$, we can expand $(f(p_z) - f(p_z'))$ and write

$$f(p_z) - f(p_z') = \frac{\partial f(p_z')}{\partial p_z'}\ \Delta p_z \qquad (3)$$

$$\Delta p_z = p_z - p_z' = -\ \frac{2m\hbar\omega}{p_z + p_z'}$$

Therefore

$$\alpha_\omega = \frac{8\pi^3 c^2}{n_r^2\,\omega^2}\cdot 2m \int_{-\infty}^{\infty} \frac{Q_\omega(p_z')}{p_z + p_z'}\ \frac{\partial f(p_z')}{\partial p_z'}\,dp_z' \qquad (4)$$

By partial integration and rearranging terms, we find the correct form of equation (3) of Simon and Strange's paper

$$\alpha_\omega = \frac{8\pi^3 c^2}{n_r^2 \omega^2}\ 2m \int_0^{\infty}\left[f(p_z')\frac{\partial}{\partial p_z'}\left(\frac{Q_\omega(p_z')}{p_z' - p_z}\right) + f(-p_z')\frac{\partial}{\partial p_z'}\left(\frac{Q_\omega(p_z')}{p_z' - p_z}\right)\right]dp_z' \qquad (5)$$

which no longer vanishes when $f(p_z) = f(-p_z)$. The criterion for a negative absorption coefficient in the case $f(-p_z) \ll f(p_z)$ (which corresponds to a population inversion in terms of a coherent streaming motion with respect to the medium in the $+p_z$ direction) is

$$\frac{\partial}{\partial p_z'}\left(\frac{Q_\omega(p_z')}{p_z' + p_z}\right) < 0 \qquad (6)$$

or $Q_\omega(p_z)$ must not rise faster than the first power of p_z.

The absorption coefficient quoted by Simon and Strange has been found by us recently to be correct only in the high quantum number regime and is totally inapplicable to the $0 \rightarrow 0$ transition. The correct expression for the transition probability in the $0 \rightarrow 0$ transition is:

$$w(p_{z_1}\omega) \simeq C\,\frac{n_r}{\omega|p_z|}\ g(\omega) \qquad (7)$$

where $g(\omega)$ is a slowly varying function and C is a constant. The presence of p_z in the denominator is purely the effect of the density of the final state of a one-dimensional particle. As is well known, the density of states of a one-dimensional particle, $\rho(\varepsilon)$, is

$$\rho(\varepsilon) = \frac{dp}{d\varepsilon} \propto \frac{1}{p} \qquad (8)$$

Accordingly

$$\frac{\mathrm{d}}{\mathrm{d}p_z}\left(\frac{Q_\omega(p_z)}{p_z}\right)$$

is negative and the absorption coefficient can be negative. Laser amplification of radiation can therefore take place.

We have applied equation (7) to electromagnetic wave amplification in our model and found that the spectrum of the resultant coherent radiation is proportional to $\nu^{-\alpha}$, where α is of the order unity but α is a function of various parameters associated with the neutron star. This type of spectrum agrees roughly with that observed. A full account of this theory is being prepared for publication.

One of us (V. C.) is an NAS–NRC senior postdoctoral resident research associate.

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Possibility of a Terrestrial Component in the Doppler Shifted Zodiacal Light

LIGHT scattered off particles in circum-terrestrial space may make an appreciable contribution to the zodiacal glow. Just how significant this contribution might be unfortunately is still unknown. The information we have comes directly from the infrared observations by Petersen[1] and McQueen[2]. These show that an appreciable portion of the dust must be in circum-solar orbits, some of it quite near the Sun. Direct observations on the circum-terrestrial components, however, are still ambiguous, despite extensive particle detection experiments from space probes[3]. Such experiments have been largely sensitive to particles larger than those thought responsible for the zodiacal glow.

For this reason, the recent observations of Reay and Ring[4] are welcome, and though they are not specific enough, as yet, to permit conclusions about the circum-terrestrial particles, they indicate that the spectroscopic method is intrinsically sensitive enough to permit an early decisive test.

Our aim in this letter is first to show that the use of Mie theory does not appreciably change the conclusions already drawn by Bandermann and Wolstencroft[5] and by Reay[6], so that an appreciably better fit to the experimental data could not be expected even if quite unusual particle properties or orbits were taken seriously. This leads to our second point, which assumes a certain circum-terrestrial dust component and then predicts that the low velocity Doppler shifts observed at high elongations will be confirmed by more detailed observations and will, in fact, be present both in the morning and evening zodiacal glow.

The difference between a circum-terrestrial dust containing model and the pure interplanetary dust model, for which James[7] has computed characteristic Doppler shift curves, is large enough so that present techniques seem capable of making a distinction.

If an appreciable fraction of the dust is concentrated near the Earth, one would expect the value of the wavelength shift to be below the theoretical curve at elongations of the order of 70°, while, as James has shown, the wavelength shift for an entirely circum solar cloud exhibits no such drop.

Like the previous authors[5,6] we assume an effective Doppler shift

$$<\delta\lambda> = \frac{\int \delta\lambda(B)\,\mathrm{d}B}{\int \mathrm{d}B}$$

where $\delta\lambda(B)$ is the wavelength shift and $\mathrm{d}B$ the surface brightness due to individual volume elements along a line of sight through the zodiacal dust cloud. $<\delta\lambda>$ is the expected averaged shift obtained from direct observations. Ingham[8] as well as Reay[6] considered that the phase function which determines $\mathrm{d}B(\theta)$ for phase angle θ is the sum of phase functions for reflected and diffracted light. Such a function, however, is not identical with that obtained from Mie theory, and we therefore consider whether appreciable changes in results are to be expected on that basis. We take a phase function

$$F(\theta, x, m) = \tfrac{1}{2}\, i_1\,(\theta, x, m) + i_2\,(\theta, x, m)$$

where $x = 2\pi a/\lambda$ is the size parameter, dependent on particle radius a, and m is the refractive index. If $C(x, m)$ is chosen as the scattering cross-section and $\varphi(x)$ represents the particle size distribution function, the phase function for the polydisperse, optically thin cloud becomes

$$\Psi(\theta) = \frac{\lambda^2 \int_{x_1}^{x_2} F(\theta, x, m)\,\varphi(x)\,\mathrm{d}x}{\pi \int_{x_2}^{x_1} C(x, m)\,\varphi(x)\,\mathrm{d}x}$$

The phase function $\Psi(\theta)$ has a characteristic minimum at $\theta \sim 90°$, which is very pronounced for dielectric particles. At elongations ε of about 90°, particles relatively distant from the Earth therefore contribute appreciably to the Doppler shift, because $\theta \gg 90°$ for these grains. This makes the value of $<\delta\lambda>$ practically independent of the physical properties assumed for the scatterers.

In Fig. 1 we show the Doppler shift as a function of elongation, computed for circular orbits about the Sun.

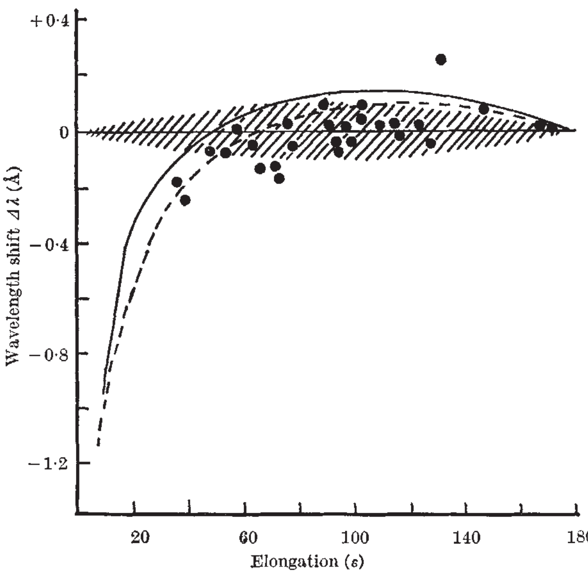

Fig. 1. Wavelength shift $\Delta\lambda$ at the evening zodiacal light as a function of elongation ε in degrees. The dashed curve represents James's results for isotropically scattered particles. The full line represents our results for Mie scatterers. The hatched area represents circum-terrestrial grains with peak elongations somewhat arbitrarily chosen as ± 6 km s^{-1}. The full points are data gathered by Reay and Ring.