i.e. the position angle of the observed linear polarization is virtually independent of frequency (J. A. Roberts, private communication). One could argue this away by suggesting that the small linear and circular components in the emission come from different regions of the source. However, if one accepts that there is little Faraday rotation, then only case (a) remains. Assuming $r_c \propto \triangle T \triangle \Psi$, cf. (5), and with $\triangle\Psi \propto \omega^{-2}$, one expects $r_c \propto \omega^{-3}$ for case (a).

Conclusions which may be drawn from the foregoing discussion are:

(i) A boundary layer does not affect the expected degree of circular polarization due to ellipticity of the modes provided (a) that the ellipticity is small and (b) that the modes get out of phase substantially either inside the source or in the boundary layer.

(ii) The absence of Faraday depolarization implies little Faraday rotation inside the source, and the lack of variation of the position angle with frequency implies little Faraday rotation either inside the source or in the boundary layer.

(iii) Accepting that there is little Faraday rotation, the propagation-induced degree of circular polarization is very small and increases rapidly ($\propto \omega^{-3}$) with decreasing frequency. It is then not plausible to invoke this explanation of the circular polarization.

(iv) Alternatively, one can argue that the linear and circular components come from different regions of the source. One then expects $r_c \propto \omega^{-1}$ on the assumption that there is a "smooth" boundary layer.

(v) That the observed circular polarization arises as a propagation effect is at best no simpler an interpretation of the observations than that adopted by Roberts *et al.* (1975), namely that the circular polarization is due to the intrinsic circular component in synchrotron emission.

I have profitted from stimulating discussions with Dr M. M. Komesaroff and Dr J. A. Roberts on this topic, and Dr G. A. Dulk made a number of helpful criticisms of an earlier version of the manuscript. The typographical errors in the paper by Sazonov were pointed out to me by Dr P. D. Noerdlinger.

Jones, T. S. and O'Dell, S. L., *Astrophys. J.*, **214**, 522 (1977a).
Jones, T. W. and O'Dell, S. L., *Astrophys. J.*, **215**, 236 (1977b).
Pacholczyk, A. G., *Mon. Not. Roy. Astron. Soc.*, **163**, 29P (1973).
Roberts, J. A., Roger, R. S., Ribes, J.-C., Cook, D. J., Murray, J. D., Cooper, B. F. C. and Biraud, F., *Aust. J. Phys.*, **28**, 325 (1975).
Sazonov, V. N., *Zh. Eksp. Teor. Fiz.*, **56**, 1075; *Soviet Phys. JETP.*, **29**, 578.

A Gyro-Synchrotron Maser in the Solar Corona?

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Introduction

Stewart (1978) has reported four moving type IV bursts observed with the Culgoora radio heliograph at 43, 80 and 160 MHz. After an early phase, the brightness temperatures of the observed bursts decreased with increasing frequency and with time. The highest brightness temperature observed at 43 MHz was 10^{10}K, and it seems that the brightness temperature would have been still higher at even lower frequencies. Existing theoretical ideas on moving type IV bursts are based on data (at 80 MHz primarily) which included no brightness temperatures in excess of 10^9K. the accepted interpretation involved gyro-synchrotron radiation from mildly relativistic electrons (energies $\approx$ 100 keV); reabsorption by the electrons themselves restricts the brightness temperature to less than about 100 keV $\approx 10^9$K (Wild and Smerd 1972, Dulk 1973). Stewart's (1978) new data at 43 MHz require that this accepted interpretation be modified; he has suggested that higher energy electrons are involved. An alternative suggestion is explored here, namely that the absorption might be negative. In other words, the high brightness temperatures observed could be due to a gyro-synchrotron maser involving electrons with energies of about 100 keV.

The possibility of a cyclotron maser has been known for many years, e.g. Twiss (1958), Hirshfield and Wachtel (1964). In the present case it is unlikely that the observed emission at 43 MHz is at the fundamental of the gyrofrequency Ω_e, which would require $B = 14$G. For a more reasonable value of the magnetic field, say $B = 5$ to 10G, the emission at 43 MHz would be at the second or third harmonic. Here we explore the possibility of negative absorption at low harmonics of the gyrofrequency. Relativistic effects are neglected in the initial discussion, and it is later pointed out that relativistic effects become important at higher harmonics even for the mildly relativistic electrons ($\approx$100 keV) under consideration.

Negative Gyromagnetic Absorption

Gyromagnetic absorption can be negative either because the distribution function $f(p)$ of the radiation electrons is an increasing function of momentum p or because it is anisotropic favouring particles with pitch angles α close to $\pi/2$. Bekefi (1966, p.302) discussed a specific example of the first type; we think that $\partial f(p)/\partial p > 0$ is an unlikely condition in an astrophysical source and discuss this case no further. Anisotropic distribution with $\partial f(p,\alpha)/\partial \sin\alpha > 0$ can be generated simply by particles propagating in a direction of increasing magnetic field strength, as in laboratory cyclotron masers.

(i) Bi-Maxwellian Streaming Distribution.

The first distribution we discuss is a (non-relativisitic) bi-Maxwellian streaming distribution

$$f(p,\alpha) = \frac{n_1}{(2\pi)^{3/2} m_e^3 c^3 \beta_{\perp o}^2 \beta_{\parallel o}} \exp\left[-\frac{\beta^2\sin^2\alpha}{2\beta_{\perp o}^2} - \frac{(\beta\cos\alpha - \beta_s)^2}{2\beta_{\parallel o}^2}\right], \quad (1)$$

where n_1, $\beta_{\perp o}$, $\beta_{\parallel o}$ and β_s are constants, and with $p = m_e c\beta$ here. The distribution (1) has been discussed by Melrose (1973, 1976). We neglect the streaming motion here, i.e. set $\beta_s = 0$, because it is not important in the following discussion (although it is essential in leading to the current required to maintain the magnetic field of the plasmoid). At the sth harmonic, the absorption coefficient is negative only when

$$g(s,\omega) = \frac{s\Omega_e}{\omega} + \frac{\beta_{\perp o}^2}{\beta_{\parallel o}^2}\left(1 - \frac{s\Omega_e}{\omega}\right) \quad (2)$$

is negative. The line width is determined by the parallel Doppler spread (see below however) and its half-width is given roughly by

$$\left| \frac{\omega - s\Omega_e}{s\Omega_e} \right| \lesssim \beta_{\|o}|\cos\theta|, \quad (3)$$

where θ is the angle of emission. One finds that effective growth occurs only for (Melrose 1973, 1976)

$$\beta_{\perp o}^2 \gtrsim \beta_{\|o}. \quad (4)$$

The absorption coefficient is then negative for

$$\omega \lesssim s\Omega_e(1 - \beta_{\|o}), \quad (5)$$

with a maximum negative value of order

$$|\gamma(s)|_{max} \approx \frac{\omega_{P1}^2}{\omega\beta_{\|o}^2} \frac{(\frac{1}{2}s\beta_{\perp o})^{2s}}{s!}, \quad (6)$$

with $\omega_{P1}^2 = 4\pi n_1 e^2/m_e$.

(ii) Anisotropic Power-Law Distribution.

The second type of non-relativistic distribution we consider is one of the form

$$f(p,\alpha) \propto p^{-a}(\sin\alpha)^q. \quad (7)$$

Power-law distributions in momenta are usually assumed in this connection, and observational evidence favours a steep power law, e.g. $a \approx 10$ (Frost and Dennis 1971, Takakura 1972, Dulk 1973). Power-law distributions in $\sin\alpha$ have been assumed in connection with the interpretation of the Jovian synchrotron spectrum (Thorne 1963, Roberts and Komesaroff 1965), and quite steep power laws are inferred, e.g. $q \approx 45$ (Roberts 1976).

On evaluating the absorption coefficient one finds that the dominant contribution is from the lowest energy electrons for $a > 2s + 4$ and from the highest energy electrons for $a < 2s + 4$. In either case the distribution (7) must be cut off at one end or the other. For the modest values of harmonic number s of interest here, the lower energy particles dominate. If $\beta_1 c$ is the speed of the lowest energy particles, one finds growth only for $\omega < s\Omega_e$ and

$$\left| \frac{\omega - s\Omega_e}{s\Omega_e\cos\theta} \right| \lesssim \frac{\beta_1}{q^{1/2}} \left\{ \frac{2a(a - 2s + 2)}{(a - 2s)(a + 2 + \frac{2as}{q})} \right\}^{1/2} \approx \frac{\beta_1}{(2/q)^{1/2}} \quad (8)$$

where the approximation applies for sufficiently large q and for $a \gg 2s$. The maximum growth rate, analogous to (6) is then

$$|\gamma(s)|_{max} \approx \frac{\omega_{P1}^2}{\omega} q\beta_1^{2s-2}. \quad (9)$$

again for large q and $a \gg 2s$.

The optical depth

The foregoing discussion shows that gyromagnetic absorption can be negative at low harmonics. Effective maser action requires that the optical depth in the region of negative absorption be greater than unity. It follows from (6), that the optical depth is given roughly by

$$\tau \approx \frac{|\gamma(s)|_{max}L}{c} \approx \frac{\omega_{P1}^2 L}{\omega\beta_{\|o}^2 c} \frac{(\frac{1}{2}s\beta_{\perp o})^{2s}}{s!}. \quad (10)$$

Setting the path length for amplification L equal to $\beta_{\|o}L_B$ (Melrose 1976) with L_B equal to 3×10^{10}cm, and taking $n_1 = 1\text{cm}^{-3}, \beta_{\perp o}^2 = \beta_{\|o} = 0.1$, one finds $\tau \approx 10^2$, for $s = 2$ or 3.

Application to Moving Type IV Bursts

The following points are relevant to the application of the foregoing discussion to the interpretation of moving type IV bursts:-

1. The absorption coefficient can be negative at low harmonics when the distribution of radiating electrons is anisotropic in the sense that the average perpendicular energy is much greater than the average parallel energy. The required conditions could result from the directed motion of the electrons (associated with a current flow say) forcing them into a region of increasing magnetic field strength.

2. The optical depth is greater than unity for plausible parameters in the range where the non-relativistic approximation applies. The non-relativistic approximation fails above the harmonic number at which the bandwidth due to relativistic effects exceeds the separation between different harmonics, e.g. Bekefi (1966, p.202). This occurs for $s^{3/2}\beta_{\perp o}^2 \gtrsim 1$. Hence, for $\frac{1}{2}mc^2\beta_{\perp o}^2 = 100$ keV $= \frac{1}{5}mc^2$, one finds that the non-relativistic approximation is not entirely adequate even at $s = 2$. The formulae derived above might be applicable at $s = 2$ and 3 for parameters relevant here, but not at $s \geq 4$ for 100 keV electrons.

3. The maser should be saturated in the sense that the power radiated balances the power input into the perpendicular motion due to propagation into the region of increasing magnetic field. Let a fraction η of the total number N of particles take part in the maser. Their total energy increases at the rate

$$\frac{d}{dt}(\tfrac{1}{2}\eta Nmc^2\beta_\perp^2) = (\tfrac{1}{2}\eta Nmc^2\beta_\perp^2)\frac{v_D}{L_B}, \quad (11)$$

where we have used

$$\frac{d}{dt}\left(\frac{\beta_\perp^2}{B}\right) = 0, \qquad \frac{dB}{dt} = v_D\frac{B}{L_B} \quad (12)$$

with v_D a drift velocity and L_B the characteristic distance over which B changes. Maxwell's equations give

$$j = n_1 e v_D \approx \frac{c}{4\pi}\text{curl } B \approx \frac{cB}{4\pi L_B}, \quad (13)$$

i.e.

$$v_D = \frac{cB}{4\pi n_1 e L_B} \approx \tfrac{1}{2} \times 10^{19} \frac{B}{L_B n_1} \text{ cm s}^{-1}. \quad (14)$$

With $B = 10$G, $L_B = 3 \times 10^{10}$cm, $n_1 = \text{cm}^{-3}$, $N = 10^{33}$ and $\frac{1}{2}mc^2\beta_\perp^2 = 100$keV, the total power available is then $\eta 10^{25}$erg s^{-1} and the drift speed is comparable with the random speed of the particles. The total power in the moving type IV bursts reported by Stewart (1978) is less than about 10^{21}erg s^{-1}. Supposing that all the power available goes into emission at $s = 2$ or 3, a saturated maser operating in a small section of the source ($N \approx 10^{-4}$) could account for the power observed.

4. Amplified emission at the fundamental ($s = 1$) is possible, and it can escape provided the plasma frequency is less than the gyrofrequency. Most of the power radiated would then go into the fundamental. Consequently, one might expect a much higher brightness temperature at lower frequencies, e.g. $T \gtrsim 10^{12}$K at $f \lesssim 20$ MHz for a source with B = 5 to 10G in which 1% of the particles are involved in the maser at any given time.

5. The broad frequency spectrum observed requires that the maser operate over a region where the range of values of B (and hence in $\omega = s\Omega_e$ with s = 2 or 3) is roughly two to one. One formal problem arises: the maser operates only at $\omega < s\Omega_e$, with positive absorption at $\omega > s\Omega_e$. Emission in the direction of decreasing B must pass through a region where absorption occurs. Favourable escape paths, e.g. towards increasing B, should exist, but then one would expect to see bright spots, i.e. for the source to show fine structure.

6. Undesirable features of the proposed mechanism, when compared with the observed properties, include the following:- (a) one would expect to see fine spatial structure; (b) one would expect the frequency spectrum to rise steeply towards decreasing frequencies where the maser is most effective; (c) significant amplification strongly favours the extraordinary mode, and quite high degrees of polarization would be expected.

Conclusion

A gyro-synchrotron maser is a possible interpretation of the observed high brightness temperatures (10^{10}K) in some moving type IV bursts. However, unless considerably higher brightness temperatures, e.g. at lower frequencies, are found, it seems that non-amplified emission from electrons with energies (≈ 1 MeV) higher than considered in the past is a more plausible interpretation (Stewart *et al.* 1978). If the brightness temperature at lower frequencies is found to be much higher, e.g. $T \gtrsim 10^{12}$K at $f \lesssim 20$ MHz, this would provide evidence in favour of the maser mechanism.

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Bekefi, G. "Radiation Processes in Plasmas", Wiley, New York (1966).
Dulk, G. A. *Solar Phys.* **32**, 491 (1973).
Frost, K. J. and Dennis, B. R. *Astrophys. J.* **165**, 665 (1971).
Hirshfield, J. L. and Wachtel, J. M. *Phys. Rev. Lett.* **12**, 533 (1964).
Melrose, D. B. *Aust. J. Phys.* **26**, 229 (1973).
Melrose, D. B. *Astrophys. J.* **207**, 651 (1976).
Roberts, J. A. *Proc. Astron. Soc. Aust.* **3**, 53 (1976).
Roberts, J. A. and Komesaroff, M. M. *Icarus* **4**, 127 (1965).
Stewart, R. T., Duncan, R. A., Suzuki, S and Nelson, G. T. *Proc. Astron. Soc. Aust.* **3**, . . . (1978).
Takakura, T. *Solar Phys.* **26**, 151 (1972).
Thorne, K. S. *Astrophys. J. Supp.* **8**, 1 (1963).
Twiss, R. Q. *Aust. J. Phys.* **11**,564 (1958).
Wild, J. P. and Smerd, S. F. *Ann. Rev. Astron. Astrophys.* **10**, 159 (1972).

Energy Changes of Cosmic Rays

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Introduction

Recently (Gleeson (1972), Quenby (1973), Gleeson and Webb (1974,1978)) it has been shown that the mean rate of change of momentum of cosmic rays reckoned for a volume fixed in the solar system is

$$<p> = \frac{1}{3} p\mathbf{V}.\mathbf{G}, \quad (1)$$

where $\mathbf{G} = (1/U_p)(\partial U_p/\partial \mathbf{r})$ si the cosmic-ray density gradient with U_p the differential number density with respect to momentum p at position $\mathbf{r}$. (cf also the integral form of (1) by Jokipii and Parker 1967).

In this paper we demonstrate, with an example of cosmic-rays moving in a uniform tube in which there is a convective background with constant speed V, that the formula (1) correctly gives the cosmic ray momentum changes. In this example the divergence of the solar wind velocity is zero and the adiabatic deceleration process with change $-\frac{1}{3}p\nabla.\mathbf{V}$ is inoperative. A full version of the theory and a complementary example are given in Gleeson and Webb (1978 a, b) with complete references.

The Example.

The physical situation is depicted in Figure 1. Cosmic-rays from the large reservoir of particles enter the tube at $x = 0$, and are subsequently scattered by the 'heavy' scattering centres embedded in the uniform convective background moving with velocity $V\mathbf{e}_x$ down the tube. The particles are scattered without change of energy in a frame of reference moving with the scatterers. Distance down the tube is denoted by x. A_0 is the uniform cross-sectional area of the tube; and as is usual in cosmic-ray studies the particle speed $v \gg V$. We assume that the particle distribution function in momentum position space $F(\mathbf{x,p})$ is of the form

$$F(\mathbf{x,p}) = F_0(x,p) + F_1(x,p) \cos\theta, \quad (2)$$

where θ is the angle between the x axis and the particle momentum $\mathbf{p}$, and a steady-state situation prevails. With this

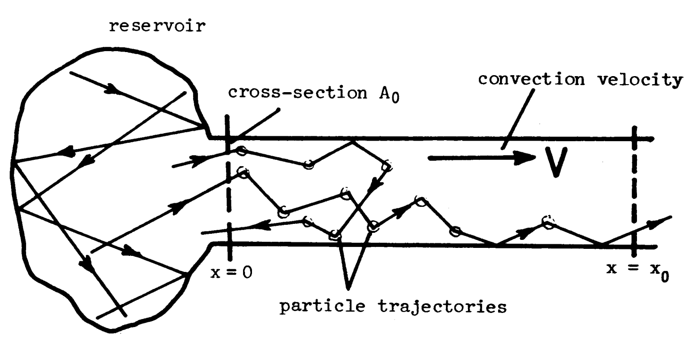

Figure 1. Illustrating the arrangement used.