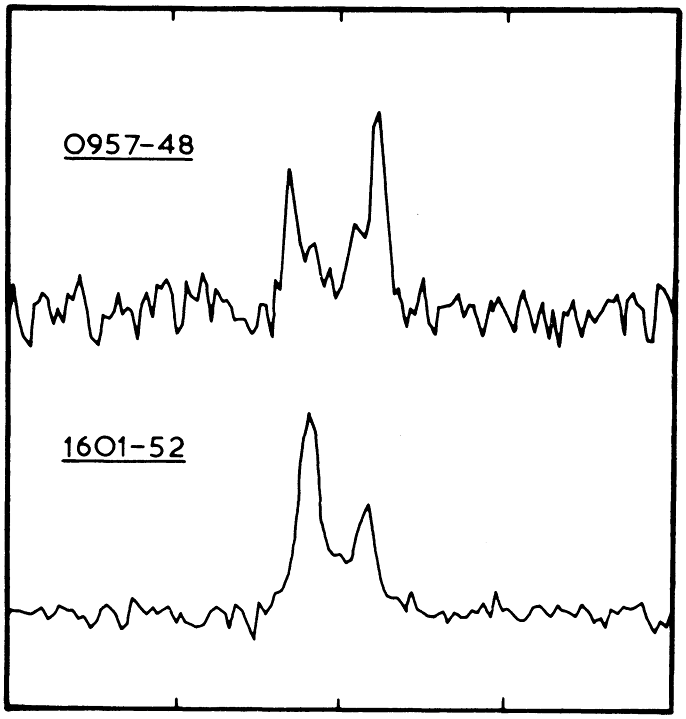

Figure 2. The pulse profiles of two pulsars with unusually wide pulses. The whole period is displayed.

pulse profiles (Fig. 2) and two have the largest dispersion measures yet recorded (524 cm^{-3} pc and 505 cm^{-3} pc).

The observations so far suggest that this experiment will approximately double the number of known pulsars and provide a uniform sample of about 250 objects.

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Davies, J. G., Lyne, A. G., and Seiradakis, J. H., *Mon. Not. R. Astron. Soc.* (in press) (1977).
Hulse, R. A., and Taylor, J. H., *Astrophys. J. (Lett.),* **201**, L55 (1975).
Large, M. I., and Vaughan, A. E., *Mon. Not. R. Astron. Soc.*, **151**, 277 (1971).

The Natural Wave Modes in a Pulsar Magnetosphere

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Our purpose in this paper is to explore the properties of the natural wave modes of a relativistically streaming electron-positron gas and to apply the results to the interpretation of the polarization characteristics of pulsar radio emission.

The Response Tensor

The response tensor for an electron-positron gas in the presence of a strong magnetic field was calculated by Svetozarova and Tsytovich (1962) and in the long-wavelength limit, assumed to be relevant here, their result was corrected by Melrose (1974). For a monoenergetic distribution of particles all in their lowest Landau orbitals (an assumption thought appropriate for a pulsar magnetosphere, cf. Blandford 1975), a fraction f of which are positrons with the remaining fraction 1-f being electrons, the polarization tensor α_{ij} in the rest frame reduces to

$$\alpha_{11} = \alpha_{22} = \frac{\omega_p^2}{\pi}\frac{m^2c^3\omega^2}{[(\hbar\omega^2-2\Omega_e mc^2)^2-4m^2c^4\omega^2]},$$

$$\alpha_{12} = -\alpha_{21} = i(1-2f)\frac{\omega_p^2}{2\pi}\frac{mc\omega(\hbar\omega^2-2\Omega_e mc^2)}{[(\hbar\omega^2-2\Omega_e mc^2)^2-4m^2c^4\omega^2]}$$

$$\alpha_{33} = \frac{\omega_p^2}{\pi}\frac{m^2c^3}{(\hbar\omega^2-4m^2c^4)},$$

$$\alpha_{13} = \alpha_{31} = \alpha_{23} = \alpha_{32} = 0, \qquad (1)$$

where ω_p is the plasma frequency and Ω_e is the gyrofrequency. Intrinsic relativistic quantum effects are included in (1), and these are unimportant for $\hbar\omega \ll mc^2$ and $\hbar\omega^2 \ll \Omega_e mc^2$. These inequalities correspond to photon energies $\ll$0.511 MeV and $\ll 7.5\times10^4 B_{12}^{1/2}$ eV respectively, where B_{12} is the magnetic induction in units of 10^{12}G.

The result (1) applies to an electron-positron gas at rest and the corresponding result for an electron-positron gas streaming with speed βc (Lorentz factor γ) along the magnetic field may be obtained from it by applying a Lorentz transformation (Melrose 1973). Writing the result in terms of the dielectric tensor $\varepsilon_{ij} = \delta_{ij} + (4\pi c/\omega^2)\alpha_{ij}$, we find

$$\varepsilon_{11} = \varepsilon_{22} = 1-\frac{\omega_p^2}{\gamma\omega^2}f_{11}, \quad \varepsilon_{12} = -\varepsilon_{21} = -\frac{\omega_p^2}{\gamma\omega^2}f_{12},$$

$$\varepsilon_{13} = \varepsilon_{31} = -\frac{\omega_p^2}{\gamma\omega^2}f_{13}, \quad \varepsilon_{23} = -\varepsilon_{32} = -\frac{\omega_p^2}{\gamma\omega^2}f_{23},$$

$$\varepsilon_{33} = 1-\frac{\omega_p^2}{\gamma\omega^2}f_{33}, \qquad (2)$$

with

$$f_{11} = \frac{\omega^2(1-\mu\beta\cos\theta)^2}{\omega^2(1-\mu\beta\cos\theta)^2-\Omega_e^2/\gamma^2},$$

$$f_{12} = -\frac{i(1-2f)\omega(1-\mu\beta\cos\theta)\Omega_e/\gamma}{\omega^2(1-\mu\beta\cos\theta)^2-\Omega_e^2/\gamma^2},$$

$$f_{13} = \frac{\mu\beta\sin\theta}{(1-\mu\beta\cos\theta)}f_{11}, \quad f_{23} = -\frac{\mu\beta\sin\theta}{(1-\mu\beta\cos\theta)}f_{12},$$

$$f_{33} = \frac{1}{\gamma^2(1-\mu\beta\cos\theta)^2} + \frac{\omega^2\mu^2\beta^2\sin^2\theta}{\omega^2(1-\mu\beta\cos\theta)^2-\Omega_e^2/\gamma^2}, \quad (3)$$

where the relativistic quantum corrections are neglected, and where ω_p is the plasma frequency in the rest frame. Also, in (3) μ denotes kc/ω and θ is the angle between the wavevector **k** and the ambient magnetic field **B**. The approximations made in deriving (2) with (3) (the long-wavelength limit and the neglect of relativistic quantum effects) require

$$\hbar\omega^2\mu^2\beta^2\sin^2\theta \ll \Omega_e mc^2, \quad \hbar\omega|\mu\cos\theta-\beta| \ll mc^2\beta,$$

$$\hbar\omega\gamma(1-\mu\beta\cos\theta) \ll mc^2, \quad \hbar\omega^2\gamma^2(1-\mu\beta\cos\theta)^2 \ll \Omega_e mc^2, \quad (4)$$

which inequalities should be well satisfied for the radio emission observed from pulsar magnetospheres. In fact, (2) with (3) could have been derived by applying the (classical) magneto-ionic theory to a streaming electron-positron gas.

The generalization of (2) from an arbitrary distribution of particles in momentum (p_z) is given by multiplying (2) by $f(p_z)$ and integrating over p_z, with $\int_{\infty}^{\infty} dp_z f(p_z) = 1$. Tsytovich and Kaplan (1972) have discussed the qualitative features of the distribution function for a relativistic non-quantum plasma in an intense magnetic field. Their work has been criticized by Suvorov and Chugunov (1973), but as these authors do not consider quantum effects in the relaxation to the ground state of gyrational motion their results may not be applicable to pulsars.

The relativistic quantum effect which is most likely to be important is the polarization of the magnetized vacuum. For $B \ll B_c (= m^2c^3/\hbar e \approx 4.4 \times 10^{13}$G) and $\hbar\omega \ll mc^2$ the additional contribution to ε_{ij} from the vacuum alone is (Melrose and Stoneham 1976)

$$\varepsilon_{11} = 1 - \frac{2\alpha L^2}{45\pi}\sin^2\theta, \quad \varepsilon_{22} = 1 + \frac{4\alpha L^2}{45\pi}\sin^2\theta,$$

$$\varepsilon_{33} = 1 + \frac{\alpha L^2}{45\pi}(5\cos^2\theta + 7\sin^2\theta)$$

$$\varepsilon_{13} = \varepsilon_{31} = -\frac{2\alpha L^2}{45\pi}\sin\theta\cos\theta, \quad \varepsilon_{12} = \varepsilon_{21} = \varepsilon_{23} = \varepsilon_{32} = 0, \quad (5)$$

where $\alpha = e^2/\hbar c \approx 1/137$ is the fine structure constant and $L = B/B_c$.

The Properties of the Wave Modes

We calculate the properties of the wave modes assuming that the refractive index is close to unity, i.e. for $|1-\mu| \ll 1$, and that the waves may be regarded as transverse. The refractive indices of the two modes are denoted $\mu_\pm$ and the polarization is described by the axial ratios $T_\pm$ of the polarization ellipse (with $T>0$ for righthand polarization and $|T|>1$ implying the major axis along the projection of **B** on the plane orthogonal to **k**). Writing $\Delta\varepsilon_{ij} = \varepsilon_{ij} - \delta_{ij}$ we find

$$\mu_\pm^2 = 1 + {}^1\!/_2\Delta\varepsilon_{11}\cos^2\theta + {}^1\!/_2\Delta\varepsilon_{22} + {}^1\!/_2\Delta\varepsilon_{33}\sin^2\theta - \varepsilon_{13}\sin\theta\cos\theta \pm$$
$$\pm\{{}^1\!/_4[\Delta\varepsilon_{11}\cos^2\theta - \Delta\varepsilon_{22} + \Delta\varepsilon_{33}\sin^2\theta - 2\varepsilon_{13}\sin\theta\cos\theta]^2 -$$
$$-(\varepsilon_{12}\cos\theta + \varepsilon_{23}\sin\theta)^2\}^{1/2}, \quad (6)$$

and

$$T_\pm = \frac{-i(\varepsilon_{12}\cos\theta + \varepsilon_{23}\sin\theta)}{1+\Delta\varepsilon_{11}+\Delta\varepsilon_{33}-\mu_\pm^2(1+\Delta\varepsilon_{11}\sin^2\theta+\Delta\varepsilon_{33}\cos^2\theta+2\varepsilon_{13}\sin\theta\cos\theta)}. \quad \ldots (7)$$

Let us summarize some implications of (6) with (7) in the context of the pulsar magnetosphere model of Ruderman and Sutherland (1975):—

1. The modes are approximately linear ($T_+ = \infty$, $T_- = 0$) for

$$|\varepsilon_{12}\cos\theta + \varepsilon_{23}\sin\theta|^2 \ll {}^1\!/_4|\Delta\varepsilon_{11}\cos^2\theta - \Delta\varepsilon_{22} + \Delta\varepsilon_{33}\sin^2\theta - 2\varepsilon_{13}\sin\theta\cos\theta|^2, \quad \ldots (8)$$

and this condition can always be satisfied except for a very small range of angles near $\theta = 0, \pi$ provided the numbers of electrons and positrons are sufficiently nearly equal. Specifically, the modes are linear for

$$|f - {}^1\!/_2| \ll \frac{\Omega_e\sin^2\theta}{4\omega\gamma^3|\cos\theta-\mu\beta|(1-\mu\beta\cos\theta)^2}. \quad (9)$$

For a pure electron gas $f = 0$. Even in this case, and with no streaming, if $\omega \ll \Omega_e$ the modes will be nearly linear for almost all θ except $\theta \cong 0$ and $\theta \cong \pi$, a result which follows from classical magneto-ionic theory. Cocke and Pacholczyk (1976) have suggested that linearly polarized modes may be explained in this way. For a pure electron gas ($f = 0$) which is streaming, condition (9) can still be satisfied. However for $\omega \gtrsim \Omega_e/\gamma^3$ it can only be satisfied for $\theta \cong \gamma^{-1}$, and it is only in this range of θ that the modes are linear.

2. For sufficiently high frequencies ($\omega^2 \gg \omega_p^2/\gamma, \omega \gg \Omega_e/\gamma$) the refractive indices of the two modes are approximately equal to $\mu = 1 - \omega_p^2/2\gamma\omega^2$.

3. For approximately equal numbers of positrons and electrons and rapid streaming the modes are approximately linear except for $0 < \theta \ll \gamma^{-1}$. The effect of this small angular range on the polarization of the radiation is negligible.

To substantiate this statement consider inequality (9) for small angles of propagation ($\cos\theta \approx 1 - {}^1\!/_2\theta^2$, $\sin\theta \approx \theta$) and rapid streaming ($\beta \approx 1 - {}^1\!/_2\gamma^{-2}$) and with $\mu = 1$. The condition for linear modes is then

$$|f - {}^1\!/_2| \ll \frac{2\Omega_e\theta^2}{\omega\gamma^3|\theta^2 - (1/\gamma^2)|(\theta^2 + (1/\gamma^2))^2}$$

For $\theta = \gamma^{-1}$ this is always satisfied (although when $\mu \neq 1$ is used in (9) this is not necessarily the case). For $\theta = \frac{1}{2}\gamma^{-1}$, for example, the modes are linear for $|f - \frac{1}{2}| \ll (32/75)\Omega_e\gamma/\omega$, which is easily satisfied for $f \approx \frac{1}{2}$. For $\theta^2 \ll \gamma^{-2}$ the modes are linear for

$$|f - \tfrac{1}{2}| \ll 2\Omega_e\theta^2\gamma^3/\omega$$

Circular polarization is significant only for $(\theta\gamma)^2 \ll (\omega/2\Omega\gamma)|f - \frac{1}{2}|$. For $f \approx \frac{1}{2}$ this is a very small range of angles near $\theta = 0$. The curvature of the field lines means that this small range of angles has negligible effect on the polarization of the radiation.

4. The contribution of the electron-positron gas to the dielectric tensor dominates that due to the vacuum.

5. For self-consistency, the refractive indices derived using (6) with (2) and (3) must satisfy $|\mu_\pm{}^2 - 1| \lesssim \gamma^{-2}$. This is not the case for the + mode for a small range of angles near $\theta = \gamma^{-1}$ where the mode becomes hydromagnetic, but the effect of this on the propagation of pulsar radio emission is expected to be negligible since θ increases rapidly due to the curvature of the field lines.

Application to Pulsar Radio Emission

Observations show that emission in orthogonal modes is a characteristic property of pulsar radio emission (Manchester *et al.* 1975, Backer *et al.* 1976). The orthogonality appears in the integrated profile as well as in individual subpulses and between consecutive subpulses. The relevance of the foregoing theory to these observations follows directly from the work of Cocke and Pacholczyk (1976) who discussed the effect of linearly polarized natural modes and concluded that emergent radiation polarized in orthogonal modes could result when the rotation depth $\Delta ks \approx \pi/2$, where $\Delta k = (\omega/c)(\mu_+ - \mu_-)$ and s is the distance traversed through the plasma. Their arguments apply in the present case where the linear natural modes are determined by the relativistically streaming electron-positron gas. However, for the two modes to propagate independently one requires $\Delta ks \gg 1$. This feature was used by Backer *et al.* (1976) to interpret the emission in orthogonal modes in terms of a geometric model with slightly different ray paths for the two modes. In addition, the condition $\Delta ks \approx \pi/2$ is unnecessarily restrictive since $\Delta ks \gg 1$ also results in emission in orthogonal modes.

For linearly polarized natural modes in a relativistically streaming electron-positron gas, (6) with (2) and (3) gives

$$\Delta k = \frac{\omega_p{}^2\Omega_e{}^2\sin^2\theta}{2\omega c\gamma^5(1-\beta\cos\theta)^2\,|\,\omega^2(1-\beta\cos\theta)^2 - \Omega_e{}^2/\gamma^2\,|}$$

$$\approx \frac{\omega_p{}^2\sin^2\theta}{2\omega c\gamma^3(1-\beta\cos\theta)^2} \quad (\text{for } \omega^2(1-\beta\cos\theta)^2 \ll \Omega_e{}^2/\gamma^2)$$

and for a pulsar s may be approximated by the radius of the velocity of light cylinder. Below a critical frequency ω_c (corresponding to $\Delta ks = 1$) a large rotation depth is possible for radio emission propagating through a pulsar magnetosphere and above this frequency strictly orthogonal modes are no longer exhibited. As the frequency increases above ω_c we predict a randomization of the polarization angle, a decrease in the degree of linear polarization and an increase in the degree of circular polarization. The first two predictions are consistent with the observations of Manchester *et al.* (1975) and Manchester *et al.* (1973), respectively, but sufficient data on the degree of circular polarization of pulsar radio emission is not yet available to test the third prediction.

Backer, D. C., Rankin, J. M. and Campbell, D. B., *Nature* **263**, 202 (1976).

Blandford, R. D., *Mon. Not. R. Aston. Soc.* **170**, 551 (1975).

Cocke, W. J. and Pacholczyk, A. G., *Astrophys. J.* **240**, L13 (1976).

Manchester, R. N., Taylor, J. H. and Huguenin, G. R., *Astrophys. J.* **179**, L7 (1973).

Manchester, R. N., Taylor, J. H. and Huguenin, G. R., *Astrophys. J.* **196**, 83 (1975).

Melrose, D. B., *Plasma Phys.* **15**, 99 (1973).

Melrose, D. B., *Plasma Phys.* **16**, 845 (1974).

Melrose, D. B. and Stoneham, R. J., *Nuov. Cim.* **32A**, 435 (1976).

Ruderman, M. A. and Sutherland, P. G., *Astrophys. J.* **196**, 51 (1975).

Suvorov, E. V. and Cugunov, Yu. V., *Astrophys. Space Sci.* **23**, 189 (1973).

Svetozarova, G. I. and Tsytovich, V. N., *Izv. Vysshikh. Uchebn. Zavedenii Radiofiz.* **5**, 658 (1962).

Tsytovich, V. N. and Kaplan, S. A., *Astrofizika* **8**, 441 (1972).

Galactic and Extragalactic

Observations of 4U 1608-52

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Tananbaum *et al.* (1976) have recently added to the Uhuru catalogue 4U1608-52, a source varying by over an order of magnitude in a time span of some months. They identified the object with MXB1608-52, a highly variable X-ray source of flares and bursts. (Kaluzienski *et al.* 1976; Belian *et al.* 1976; Grindlay and Gursky 1976a, b; Li 1976). There have been a number of earlier X-ray observations of sources very close to 4U1608-52 (Cooke and Pounds 1971; Luyendyk *et al.* 1973; Thomas *et al.* 1975; Grindlay and Gursky 1976b; Ricker *et al.* 1976), and in this paper we propose the identification of these earlier sources with 4U1608-52. We thereby obtain additional spectral information and can now rule out a globular cluster as the optical counterpart of the object.

Positional error boxes for all published observations are shown in Figure 1. The observational history is summarized in Table I in which are tabulated the dates of observations, the energy bands in which the observations were made, the recorded intensities relative to that of the Crab Nebula, measured spectral parameters and references to the literature.

If the proposed identifications are correct, a number of comments may be made with regard to 4U1608-52. Firstly, it is highly variable, displaying bursts which last from 10 to 100