Pulsar radio emission mechanism: Why no consensus?

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Abstract. Pulsar radio emission involves a coherent emission mechanism involving relativistic particles, but there is no consensus on what the specific emission mechanism is. There are arguments for and against several suggested mechanisms of long-standing, including two, relativistic plasma emission and anomalous Doppler emission, that depend explicitly on the dispersive properties of waves in a pulsar plasma. Models that do not include dispersion due to the relativistic spread in energy of the plasma particles can be misleading, and including the effects of this spread on the wave dispersion leads to severe constraints on these two suggested mechanisms. It is argued that all suggested mechanisms are implausible. A variant of relativistic plasma emission involving large-amplitude oscillations needs to be explored as a possible alternative.

1. Introduction
Five decades after the discovery of radio pulsars there is no consensus on what the pulsar radio emission mechanism is. There is wide agreement on several general features of the emission mechanism: the very high brightness temperature implies some form of coherent emission; the emission involves relativistic electrons (and positrons) [1] generated in pair cascades [2]; and the source region is on “open” field lines that define the “polar-cap” regions of the magnetosphere [3]. There are two well-established coherent emission mechanisms in astrophysical plasmas [4]: plasma emission and electron cyclotron emission. Pulsar radio emission is different from either of these, e.g. because it involves relativistic particles, and is a third form of coherent emission. Several emission mechanisms were suggested within the first decade or two after the discovery of pulsars and continue to be of interest: coherent curvature emission (CCE), relativistic plasma emission (RPE), anomalous Doppler emission (ADE), and linear acceleration emission (LAE) or free-electron maser emission (FEM). There are plausible arguments in favor of each of these mechanisms, but there are also arguments against each of them. The various arguments against the suggested mechanisms are so strong that one could conclude that all of them are untenable as the generic pulsar radio emission mechanism.

The suggested emission mechanisms are discussed critically in §2. Wave dispersion in a “pulsar plasma” is discussed in §3 and its negative implications for “beam-driven” RPE are discussed in §4. In §5 we speculate on non-beam-driven forms of RPE and summarize our conclusions in §6.
2. Critique of existing mechanisms

Historically, CCE was the first suggested emission mechanism \([1, 5, 6, 7]\), and there are ongoing arguments in favor of it \([8]\). CCE is due to one-dimensional (1D) motion of a highly relativistic particle along a curved magnetic field line. The frequency for incoherent curvature emission peaks at about \(\left(\frac{c}{R_c}\right)\gamma^3\), where \(R_c\) is the radius of curvature of the magnetic field line, and high \(\gamma\) is required for the frequency to be in the observed radio range. However, there are strong arguments against the postulated coherence due to bunches of highly relativistic particles \([3, 9, 10]\), and although coherence due to maser is possible in principle \([11, 12, 13]\), it does not provide a realistic alternative. CCE is not discussed further here.

LAE and FEM involve large-amplitude structures oscillating in time (frequency \(\omega_0\)) and space (wavevector \(k_0\)), respectively. Incoherent emission occurs at \(\omega \lesssim \omega_0\gamma^2, k_0c\gamma^2\), respectively, and maser emission is possible for a distribution that is an increasing function of \(\gamma\). Because LAE and FEM require a separate mechanism to generate the postulated large-amplitude structures, they may be interpreted as mechanisms for generating escaping radiation in the presence of such structures, and in this sense may be regarded as the second stage of a RPE mechanism.

All suggested mechanisms involve highly relativistic electrons, and three characteristic Lorentz factors are invoked in different models: a very large \(\gamma_p\) of order \(10^7\) for “primary” particles (which are not essential \([14]\)), and two characteristic Lorentz factors associated with the “secondary” pair plasma, bulk outflow at \(\gamma_s\) and an intrinsic spread \(\langle\gamma\rangle\) in the rest frame of the outflowing plasma. Models for the pair cascade \([15, 16]\) suggest values for the latter two of order \(\gamma_s \sim 10^3\) and \(\langle\gamma\rangle \sim 10^{-10^2}\). RPE and ADE depend intrinsically on the dispersive properties of the “pulsar plasma”, whereas CCE and LAE/FEM exist in vacuo, albeit modified by the dispersive properties of the plasma.

Recently it has been argued \([22]\) that the so-called nanoshots from the Crab pulsar \([23]\) favor a specific form of beam-driven RPE \([24, 25]\). RPE may be regarded as a pulsar counterpart of conventional plasma emission which involves two stages, with the first stage being an instability that generates turbulence in Langmuir-like waves, and a subsequent stage involving partial conversion of energy in this turbulence into escaping radio waves. Difficulties with both stages of RPE have long been recognized: the growth rates for various suggested instabilities in the first stage are too slow to be effective, and ways in which this might be overcome have been explored \([26, 27, 28, 29]\); the difficulty with the second stage has been referred to as a “bottle-neck” \([30]\). In §3 a further difficulty is pointed out with beam-driven RPE: there are effectively no “Langmuir-like waves” that can grow due to a beam instability. This provides a strong argument against beam-driven RPE. A similar difficulty arises with suggested beam-driven growth of Alfvén waves \([31, 32, 33, 34]\).

3. Wave dispersion in pulsar plasma

There is an extensive literature on wave dispersion in a pulsar plasma, defined here as a highly relativistic, strongly-magnetized, electron-positron pair plasma. The relativistic bulk streaming of the plasma may be included in the wave dispersion by treating it either in the plasma rest frame and Lorentz transforming to the pulsar frame \([17, 18]\) or in the pulsar frame directly \([19]\). We adopt the former approach. Although these two approaches are equivalent in principle, they lead to different results in practice due to different approximations. With the former approach, there are only three relevant wave modes, two of which depend strongly on a relativistic plasma dispersion function (RPDF), due to an assumed spread \(\langle\gamma\rangle \gg 1\) in the rest frame. With the latter approach, the RPDF is approximated by a moment of the distribution and four wave modes are identified \([20, 21]\). We note that the inconsistency involving the fourth mode is one of several unresolved problems relating to wave dispersion relevant to beam-driven RPE \([17]\), even when inhomogeneity (field-line curvature \([20]\)) and plasma turbulence \([32]\), are neglected.

It is helpful to consider the cold and nonrelativistic cases as a prelude to discussion of the
The dispersion relation for the L mode is then approximated by $\omega \approx \beta L \omega_p$ with $\beta \approx 2$. For parallel propagation, the magnetoionic approximation: the extraordinary, or X mode, has refractive index $\omega/kc = 1 + \omega_p^2/\beta A^2$, and the ordinary mode separates into a higher-frequency branch, called the O mode, and a lower-frequency branch identified as the Alfvén mode. For parallel propagation the dispersion curves for the latter two modes correspond to straight lines, cf. Figure 1: $\omega = \omega_p$ for the parallel longitudinal (L) mode and $z = \omega/kc = z_A$ for parallel Alfvén (A) mode. The two dispersion curves cross, and for nonzero angle of propagation, $\theta$, they reconnect to form the O mode and the (oblique) Alfvén mode. In the limit $\theta \to 0$, the O mode coincides with the L mode from the cutoff at $\omega/\omega_p = 1$ and for $k||c/\omega_p \in [0,1/z_A]$ and to the A mode at higher frequencies. The Alfvén mode coincides with the A mode from $\omega = 0$ to $\omega = \omega_p$ and to the L mode from $k||c/\omega = 1/z_A$ to the resonance $k||c \to \infty$. The (reconnected) O mode and Alfvén separate from each other with increasing $\theta$, the O mode to higher $\omega$ and larger $z$, and the Alfvén mode to lower $\omega$ and smaller $z$. There is a stop band between the resonance in the A mode, at $\omega \approx \omega_p |\cos \theta|$ for $\Omega_e \gg \omega_p$, and the cutoff in the O mode at $\omega = \omega_p$.

An important difference between these wave properties and those in the more familiar case $\Omega_e < \omega_p$ concerns the resonance. For $\Omega_e \ll \omega_p$ there is a resonance in the (upper-frequency branch of the magnetoionic) extraordinary mode at $\omega > \omega_p > \Omega_e$, and Langmuir waves correspond to a thermally modified form of these resonant waves. For $\Omega_e > \omega_p$ the resonance is in the (lower frequency branch of the magnetoionic) ordinary mode, at $\omega < \omega_p < \Omega_e$, interpreted as a resonance is in the Alfvén mode. As a consequence the concept of a “Langmuir-like” wave mode becomes ill-defined in a pulsar plasma with $\Omega_e \gg \omega_p$.

The inclusion of a nonrelativistic velocity spread modifies the wave properties and these modifications become extreme when the spread is relativistic. For parallel propagation, the dispersion relation for the A mode is modified by the inclusion of relativistic effects in $\beta A$, with $\beta A^2 \approx \Omega_e^2/\omega_p^2(\gamma)$. The dispersion relation for longitudinal waves becomes $\omega = \omega_L(z)$ with $\omega_L^2(z) = \omega_p^2 z^2 W(z)$ where the RPDF $z^2 W(z)$ is plotted in Figure 2 for a 1D thermal distribution for temperature $5 \times 10^9/\rho \, K$ with $\rho = 50, 10, 1$. The cold-plasma limit corresponds to $z^2 W(z) \to 1$. For nonrelativistic temperatures, $\rho \gg 1$, $z^2 W(z)$ is negative for very small $z$, crosses zero and rises to a maximum and then asymptotes to $1 + 3/\rho z^2$ for $\rho z^2 \gg 1$. The dispersion relation for the L mode is then approximated by $\omega_L(z) \approx \omega_p^2 + 3k||c^2 W(z)$, with
Figure 2. The real (solid) and negative imaginary (dashed) parts of the RPDF $z^2 W(z)$ are plotted as a function of $z$ for three 1D Jüttner distributions: leftmost $\rho = 50$, center $\rho = 10$ and rightmost $\rho = 1$. The imaginary part is identically zero for $z \geq 1$.

Figure 3. Dispersion curves for parallel propagation: left $\rho = 20$; right $\rho = 1$.

$V_e^2/c^2 = 1/\rho \ll 1$.

In the highly relativistic limit, $\langle \gamma \rangle \approx 1/\rho$, the peak in $z^2 W(z)$ becomes very large, with its maximum $\approx 2.7/\rho$ at $z = z_m \approx 1 - 0.013\rho^2$ very close to the light line, $z = 1$. The imaginary part of $z^2 W(z)$ is strictly zero for $z > 1$ and has a maximum at $z$ slightly less than $z_m$, where Landau damping is strong.

The dispersion curves for the two parallel modes intersect at a crossover frequency $\omega_{\text{co}} = \omega_L(z_A)$, as illustrated in Figure 3 for a nonrelativistic value (left: $\rho = 20$) and a relativistic value (right: $\rho = 1$). The maximum in the curve for the L mode is at the maximum of the RPDF, cf. Figure 2. In a highly relativistic plasma, the L mode dispersion curve starts at the cutoff frequency, $\omega_x$, given by $\omega_x^2 = \omega_L^2(\infty) = \omega_p^2(1/\gamma^4)$, and crosses the light line at $\omega_1$, determined by $\omega_1^2 = \omega_L^2(1) = \omega_p^2(\gamma\beta^2)$. The line $z = z_m$ through the maximum in $\omega_L(z)$ moves closer to the light line $z = 1$ with increasing $\gamma$. Reconnection between the L and A modes occurs for $z_A > z_0$, with the A mode strongly damped for $z_0 < z_A \lesssim z_m$ and weakly damped for $z \gtrsim z_m$. The condition $z_A > z_m$ is equivalent to $\beta_A^2 > \gamma_m^2 = 1/(1 - z_m^2)$, with $\gamma_m \approx 6(\gamma)$. Provided this condition is satisfied, the reconnected O and Alfvén modes move away from each other with increasing $\theta$; the maximum is then in the (reconnected) Alfvén mode.

For $\theta = 0$, the L mode is superluminal over the frequency range, $\omega_x \leq \omega \leq \omega_1 \approx \omega_x(\gamma)$. For $\theta \neq 0$ the reconnected O mode moves to the upper left, as in the cold-plasma limit Figure 1. The cutoff frequency is independent of $\theta$, and the frequency at which the dispersion curve crosses the light line increases to infinity over $0 < \theta \lesssim 1/\beta_A$. There are subluminal (L or O) waves only in this tiny range of angles.

The reconnected Alfvén mode is dominated by the peak in the RPDF, at $z = z_m$. As $\theta$ increases the maximum frequency of the Alfvén mode decreases, remaining at $z \approx z_m$. Landau
damping becomes strong at $z < z_m$, and Alfvén waves are weakly damped only on the portion of the dispersion curve at $z \gtrsim z_m$.

The application to pulsars in general and to any specific pulsar in particular requires estimates of $\Omega_e$, $\omega_p$ and $\beta_A$ as functions of height $r$ in the magnetosphere. For a given pulsar, the observed period $P$ and period derivative $\dot{P}$ imply the surface magnetic field $\propto (PP)^{1/2}$. A dipolar magnetic field falls off $\propto 1/r^3$ for $r \ll r_L$ with $r_L = Pc/2\pi$. One finds $\Omega_e/2\pi$ of order $3 \times 10^9 (P/P^3)^{1/2}(r/r_L)^{-3}$ Hz. An estimate of the plasma frequency is given by assuming that the electron number density is greater than the corotation value by a multiplicity $\kappa$, implying $\omega_p^2 \approx \kappa \Omega_e \Omega_e \propto (r/r_L)^{-3}$, with $\Omega_e = 2\pi/P$. These estimates also imply $\beta_A^2 = \Omega_e^2/\omega_p^2(\gamma) \propto (r/r_L)^{-3}$. The fiducial values adopted here are $P = 1$s, $\dot{P} = 10^{-15}$, $\kappa = 10^5$ [37] and $\langle \gamma \rangle = 10$. Evaluating the parameters at $r = 0.1r_L$ gives $\Omega_e/2\pi = 30$ GHz, $\omega_p/2\pi = 20$ kHz and $\beta_A^2 = 3 \times 10^4$. At $r = 0.1r_L$ the conditions $\omega \ll \Omega_e$ and $\beta_A > \gamma_m$ are well satisfied and plausibly satisfied, respectively, and they are more strongly satisfied at all $r < 0.1r_L$.

4. Beam-driven RPE

In a beam-driven instability waves are assumed to grow due to negative absorption resulting from resonance between the wave and a distribution function $g(u)$ with $dg(u)/du > 0$ below a peak at $u = u_b$, say, with $u_b = \gamma_b\beta_b$. The (Cerenkov) resonance condition $\omega - k_\parallel v_\parallel = 0$ corresponds to $z - \beta = 0$. Beam-driven waves have $z = \beta < \beta_b$. All velocities of relevance to the discussion here are very close to $c$, and it is convenient to refer to the corresponding Lorentz factors: $\gamma_\parallel^2 = 1/(1 - z^2)$, $\gamma_m^2 = 1/(1 - z_m^2)$, $\beta_A^2 = 1/(1 - z_A^2)$, $\gamma_k^2 = 1/(1 - \beta_k^2)$. The resonance condition becomes $\gamma_\parallel = \gamma$ with $\gamma < \gamma_b$ for the beam. The value of $\gamma_b$ depends on the assumptions made about the beam. Assuming extremely large $\gamma_b \sim 10^7$, corresponding to primary particles, gives a growth rate that is too small. To overcome this difficulty an inhomogeneous model was proposed [26, 27] in which bursts of pair creation result in localized bunches, with faster particles in a following bunch overtaking slower particles in a preceding bunch. Waves grow in such a model only if there is an identifiable beam, requiring that $\gamma_b$ be significantly greater than $\langle \gamma \rangle$. Although a detailed model is required to discuss this point quantitatively, it seems that the model is viable only if the following bunch has $\langle \gamma \rangle$ greater than that of the preceding bunch, with $\gamma_b$ of order the larger value.

A beam-driven instability is relevant only to either the O mode or the Alfvén mode: the X mode has polarization perpendicular to B and does not couple to a field-aligned beam. It follows that there are two possible versions of beam-driven RPE, involving growth of subluminal ($z < 1$) O mode waves and Alfvén-mode waves, respectively. In most treatment of these instabilities unrealistic assumptions are (implicitly) made about the dispersive properties of these waves: it is essential that the wave dispersion is calculated assuming $\langle \gamma \rangle \gg 1$ in the rest frame of the plasma.

For $\beta_A > \gamma_m$, so that the modes reconnect as discussed above, the O mode is subluminal only for $\omega > \omega_1$ and in a tiny range of angles, $\theta < 1/\beta_A$. For resonance to be possible between the beam and such O mode waves requires $\gamma_b > \beta_A$. The extremely large value of $\beta_A$ for $r/r_L \ll 1$ requires an even larger $\gamma_b$, which is possible only for primary particles and for these the growth rate is known to be unacceptably small. The condition $\gamma_b > \beta_A$ could be satisfied for $r/r_L \gtrsim 0.1$, but even in this case it seems implausible that the conditions for effective growth could be satisfied. One difficulty is that the polarization of O mode waves has a large component along B only for $\theta \ll 1/\beta_A$ and $\omega$ just above $\omega_1$, so that although growth at higher frequencies is possible in principle for sufficiently small $\theta$, the growth rate is very small. Another difficulty is that the expected frequency, $\omega \approx \omega_1$, of the resulting waves at $r/r_L \gtrsim 0.1$ is too low (even after Lorentz boosting to the pulsar frame) to provide a plausible explanation for the highest frequencies in pulsar radio emission [10]. On the other hand, a favorable aspect of beam-driven O mode waves is that the waves can escape directly: no second stage in the emission process is
required.

The other possibility for beam-driven wave growth involves Alfvén waves [31, 32, 33, 34]. Consider resonance between a beam and Alfvén waves in a pulsar plasma with \( \gamma \gg 1 \) in the limit of \( \theta \to 0 \). The dispersion curve is then \( z = z_A \) for \( \omega < \omega_L(z_A) \) and \( \omega = \omega_L(z) \) for \( z > z_A \).

In the limit \( \theta \to 0 \), the transverse polarization precludes growth along the former branch, and only waves on the latter branch can grow, and for small \( \theta \neq 0 \) this continues to be approximately the case. As \( \theta \) increases waves in the range \( z_A > z > z_m \) move to lower frequency, and it is only these waves that are of interest here. Resonance and hence growth is possible for \( z_A > \beta_b > z_m \) or \( \beta_A > \gamma_b > \gamma_m \). With \( \gamma_m \approx 6(\gamma) \), this requires \( \gamma_b > 6(\gamma) \). Thus the conditions under which these waves would be expected to grow are less extreme than for the O mode. However, unlike the O mode case, these waves cannot escape directly and a second stage in the RPE process is required.

ADE also depends intrinsically on the dispersive properties of the plasma. The anomalous Doppler resonance, \( \omega - k_v = -\Omega_e/\gamma \), may be rewritten as \( \omega = \omega_L/\gamma \) and approximated by \( \omega = \Omega_e \gamma^2_e/(\gamma^2 - \gamma^2_b) \). The frequency of ADE must satisfy both the resonance condition and the dispersion relation for a specific mode. At the expected high frequency the only relevant wave modes are the O mode and X mode. For small \( \theta \) both modes have \( \gamma_b \) of order \( \beta_A \), and resonance requires \( \gamma > \beta_A \). An order of magnitude estimate gives \( \omega \sim \beta_A \Omega_e \) for \( \gamma \sim \beta_A \). This estimate applies in the rest frame of the plasma, and the frequency in the pulsar frame is boosted by a Lorentz factor of order \( \gamma_s \). This estimate of \( \omega \) is higher than an earlier estimate based on a calculation of wave dispersion in the pulsar frame [19]. This higher frequency severely exacerbates the difficulty of explaining the observed radio frequency range with the source of ADE anywhere within the light cylinder. Due to this difficulty ADE is also at least implausible and arguably untenable.

### 5. Alternative forms of RPE

The argument in favor of RPE [22] seems strong. The foregoing criticism of RPE applies to the assumption that it requires a “beam-driven” instability to produce the assumed “coherent charge bunching and strong electromagnetic turbulence” [22]. The favorable aspects of RPE can be retained if an alternative way of generating the relevant “turbulence” can be shown to be effective. An interesting possibility is the large amplitude longitudinal oscillations (LAOs) that are driven directly by the electrodynamics [38, 39, 40]. An interpretation is that the inductive electric field associated with an obliquely rotating magnetic dipole has a component \( E_\parallel \) along the magnetic field; charges are accelerated along field lines by \( E_\parallel \) tending to set up a potential electric field that opposes \( E_\parallel \). Such screening of \( E_\parallel \) is oscillatory setting up the LAOs and triggering a pair cascade that increases the number of pairs and facilitates the screening. The source of the energy in the LAOs is the rotational energy, through the inductive \( E_\parallel \). (The idea that rotational energy may be partly converted directly into energy in LAOs has also been proposed in a different model involving a parametric instability [41, 42].) A model for LAOs [43] needs to be based on the actual wave dispersion of longitudinal waves in a pulsar plasma, that is, of the L mode waves discussed in §3. A version of RPE based on such waves has the attractive feature that the waves escape directly, by evolving into O mode waves as they propagate outwards. Alternatively the escaping radiation may result from a maser version of LAE associated with LAOs [44, 45, 46].

### 6. Conclusions

Why no consensus? One reason is that all suggested pulsar radio emission mechanisms encounter major difficulties and are arguably untenable, allowing a diversity of opinions on which is the least implausible. This diversity of opinions underlies the lack of consensus. Our opinion is that
the least implausible is a form of RPE based on LAOs driven directly by the electrodynamics. A detailed model needs to be developed for such RPE.

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