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The Mechanisms for the Generation of Pulsar Radio and High-energy Emission

George Z. Machabeli

Center for Plasma Astrophysics of Abastumani Astrophysical Observatory,
Georgian Academy of Sciences, A. Kazbegi Avenue 2a, GE-380060 Tbilisi, Georgia

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Abstract: The recent theory of propagation and generation of waves in a pulsar's magnetosphere is discussed. In particular we consider models of pulsar radio emission due to plasma instabilities. Wave–particle interactions can lead to quasilinear diffusion increasing a particle's pitch-angle. The recent model of γ-ray emission from synchrotron radiation as a result of quasilinear diffusion is discussed.

Keywords: pulsars: general — radiation mechanisms: non-thermal

1 Introduction

At present there are about 12 competing theories which differ both in the physical effects responsible for the pulsar radio emission and in the location where the radiation is generated. To date, the most widely discussed theory attributes the emission to coherent curvature emission by bunches of particles. Although this theory can explain a broad range of observed pulsar properties by the careful arrangement of the magnetic field geometry and the form and size of bunches, 30 years of theoretical efforts have failed to explain the origin of these bunches (Melrose 1995). This theory can also be ruled out on observational grounds (Lesch, Kramer, & Kunzl 1998). In addition to the work of Lesch et al. (1998) we note that this theory also fails to explain the observed correlations of the conal peaks (Kazbegi et al. 1991a) and the large size of the emitting region (Gwinn et al. 1997).

We propose that the pulsar radio emission is generated by plasma instabilities developing in the outflowing plasma on the open field lines of the pulsar magnetosphere. Plasma can be considered as an active medium that can amplify its normal modes. The wave amplification can be the result of the resonant wave–particle interaction, i.e. in the rest frame of the particle the frequency of the resonant wave is zero or a multiple of the gyrational frequency. The plasma instabilities that we argue operate in the pulsar magnetosphere may be described by the (somewhat contradictory) term 'incoherent broad-band maser'. Each single emission by a charged particle is a result of the stimulated, as opposed to spontaneous, emission process (hereby the term maser). Unlike conventional lasers in which basically one single frequency gets amplified, in this case charged particles can resonate with many mutually incoherent waves with different frequencies.

We discuss the linear theory of wave excitation in the pulsar magnetosphere developed by Lominadze, Machabeli, & Mikhailovskii (1979), Machabeli & Usov (1979), Kazbegi et al. (1991a), Kazbegi, Machabeli, & Melikidze (1991b), Lyutikov, Machabeli, & Blandford (1999), and Lyutikov, Blandford, & Machabeli (1999).

We assume that pulsar radiation is generated in the pulsar magnetosphere. A spinning magnetised neutron star generates the electric field which extracts electrons from the star surface and accelerates them. As a result a low density ($n_b = 7 \cdot 10^{-2}\, B_0 P^{-1}$, where P is the pulsar period and B_0 is the star magnetic field at the surface) and energetic (the particle Lorentz factor is $\gamma_b = 10^7$ for typical pulsars) primary beam is formed. In a weakly curved magnetic field, electrons generate γ-quanta which in turn produce electron–positron pairs. The pitch angle of the particles which are produced is non-zero, so that secondary particles generate synchrotron radiation. This radiation in turn produces more pairs, and so on until the plasma becomes dense and screens the electric field (Goldreich & Julian 1969; Sturrock 1971). As a result a multicomponent plasma is formed, containing:

(i) electrons and positrons with $\gamma \sim \gamma_p$ and $n \sim n_p$
(ii) a high-energy plasma 'tail' (the tail of the distribution function) with $\gamma \sim \gamma_t$ and $n \sim n_t$
(iii) the primary beam with $\gamma \sim \gamma_b$ and $n \sim n_b$

The equipartition of energy among the plasma components is assumed to be

$$n_p\gamma_p \approx n_t\gamma_t \approx \frac{n_b\gamma_b}{2}. \qquad (1)$$

The total energy outflow of the particles can reach a value

$$L_\pm \approx 10^{30}\left(\frac{B_0}{10^{12}G}\right)^{6/7}\left(\frac{P}{1S}\right)^{-15/7} \mathrm{erg\,s^{-1}}. \qquad (2)$$

Let us introduce cylindrical coordinates x, r, φ. The x axis is transverse to the plane where the curved magnetic field line lies, r is the radial and φ the azimuthal coordinate. The latter describes the curvature of the field line (torsion is neglected and $\partial R_B/\partial r = 0$; R_B is the curvature radius of the field line).

The pulsar magnetic field is dipolar:

$$B = B_0\left(\frac{r_0}{r}\right)^3, \qquad (B_0 = 10^{12}G), \qquad (3)$$

where r_0 is the neutron star radius.

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We have drift caused by the magnetic field inhomogeneity. The velocity of this drift

$$\frac{u_x}{c} = \frac{p_\varphi c}{\omega_B R_B} \tag{4}$$

(here p_φ is the normalized φ component (along magnetic field lines) of particle momentum and $\omega_B = eB/mc$) is directed along the positive direction of x-axis for the beam particles.

Among all pulsars discovered so far, only a small fraction emit γ-rays (e.g. Ulmer 1994). These high-energy pulsars are usually young and have strong magnetic fields of 10^{12} G. There are currently two main types of model for pulsar high-energy emission (e.g. Arons 1996): polar cap and outer gap models.

In polar cap models both acceleration of particles and the emission of the γ-rays occur near polar caps (e.g. Harding, Ozernoy, & Usov 1993; Usov 1994; Daugherty & Harding 1996; Usov & Melrose 1996). Secondary and tertiary particles produced from cascades have non-zero pitch angles and can emit synchrotron radiation, but their curvature radiation is not important. Thus, in the polar cap model, the overall spectra should consist of curvature radiation from primary particles, and synchrotron radiation from secondary and tertiary pairs, together with inverse Compton radiation.

In outer gap models (e.g. Cheng, Ho, & Ruderman 1986), the particles are accelerated and radiate near the light cylinder. The dominant pair cascade process is pair production from photon–photon collision with photons produced through inverse Compton scattering. As in the polar cap model, secondary and tertiary $e^\pm$ contribute to γ-ray spectra through synchrotron radiation. In both types of model the emission arises predominantly from the cascade region itself as a result of synchrotron radiation and inverse Compton scattering. Apart from the standard models, γ-rays can also be produced by plasma processes, as recently proposed by Machabeli et al. (2000a). Plasma instabilities can result in quasilinear diffusion, leading to pitch-angle increase. Particles that are initially at the ground state can acquire a pitch-angle, emitting synchrotron radiation.

2 Wave Generation

We hypothesize that pulsar radiation is generated by the instabilities developing in the outflowing plasma on the open field lines in the outer regions of the pulsar magnetosphere. Radiation is generated by two kinds of electromagnetic plasma instabilities — cyclotron-Cherenkov and Cherenkov-drift instabilities. The cyclotron-Cherenkov instability is responsible for the generation of the core-type emission and the Cherenkov-drift instability is responsible for the generation of the cone-type emission (Rankin 1986). The wave generated by these instabilities are vacuum-like electromagnetic waves so they may leave the magnetosphere directly.

In contrast to most modern theories of pulsar radio emission, cyclotron-Cherenkov and Cherenkov-drift instabilities occur in the outer parts of the magnetosphere. The location of the emission region is determined by the corresponding resonant condition for the cyclotron-Cherenkov and Cherenkov-drift instabilities. Instabilities develop in a limited region on the open field lines. The size of the emission region is determined by the curvature of the magnetic field lines, which limits the length of the resonant wave–particle interaction. The location of the cyclotron-Cherenkov instability is restricted to those field lines with large radii of curvature, while the Cherenkov-drift instability occurs on field lines with curvature bounded both from above and from below. Thus, both instabilities produce narrow pulses, although they operate at radii where the opening angle of the open field lines is large.

Cyclotron-Cherenkov generation of a wave by fast particles is not new in astrophysics. For example, cosmic rays in the interstellar medium and in supernova shocks generate Alfvén waves by a similar mechanism. In the case of Alfvén waves in the non-relativistic electron–ion plasma, the frequency of the waves ω can be much smaller than the kv term and can be neglected in the resonance condition. The important difference between these applications and cyclotron-Cherenkov instability in the pulsar magnetosphere is that the generated waves belong not to the hydrodynamic Alfvén waves that cannot leave the plasma, but to near-vacuum electromagnetic waves.

The cyclotron-Cherenkov instability develops at the anomalous cyclotron resonance

$$\omega(\vec{k}) - k_\parallel v_\parallel + \frac{\omega_B}{\gamma} = 0, \tag{5}$$

where ω is the frequency of the normal mode, $\vec{k}$ is a wave vector, v is the velocity of the resonant particle, $\omega_B = |e|B/mc$ is the non-relativistic gyrofrequency, γ is the Lorentz-factor in the pulsar frame, e is the charge of the resonant particle, m is the mass, and c is the speed of light. Note a plus sign before the ω_B term.

The cyclotron-Cherenkov instability may be considered as a maser using the induced cyclotron-Cherenkov emission. The free energy for the growth of the instability comes from the non-equilibrium anisotropic distribution of fast particles. The condition that the emission dominates the absorption requires population inversion in the distribution function of fast particles (maser action).

There is a possibility for the development of the Cherenkov-drift instability, which occurs at the resonance

$$\omega(\vec{k}) - k_\parallel v_\parallel - k_\perp u_d = 0, \tag{6}$$

where $u_d = \gamma v_\parallel c/\omega_B R_B$ is the relativistic drift velocity. A weak inhomogeneity of the magnetic field results in a curvature drift motion of the particle perpendicular to the local plane of the magnetic field line. A gradient drift [proportional to $(\vec{B} \cdot \vec{\nabla})$] is much smaller than the curvature drift and will be neglected. When the motion of the particle parallel to the magnetic field is ultrarelativistic, the drift motion can become weakly relativistic even in a weakly inhomogeneous field resulting in the generation of electromagnetic, vacuum-like waves. The presence of three ingredients (a strong but finite magnetic field, inhomogeneity of the field, and a medium with the index of refraction larger than unity) is essential for this type of emission. We will call this mechanism Cherenkov-drift emission stressing the fact that microscopically it is virtually a Cherenkov-type emission process.

It must be mentioned that the physics of Cherenkov-drift emission is new (Kazbegi et al. 1991b; Lyutikov et al. 1999).

If $u_d = 0$ we have ordinary Cherenkov resonance and the electric field $\vec{E}$ of generated waves is polarised along the magnetic field.

But, if $u_d \neq 0$, particle drift motion is perpendicular to the local plane of the curved magnetic field line. The emitted electromagnetic waves are polarised along u_d. It is the principal difference between Cherenkov and Cherenkov-drift mechanism.

Now let us discuss electromagnetic processes associated with a charged particle moving in a strong circular magnetic field. Attacks on this problem are made by Blandford (1975), Melrose (1978), Zhelezniakov & Shaposhnikov (1979), Kazbegi, Machabeli, & Melikidze (1987), and Luo & Melrose (1992). These works follow the approach which emphasises the analogy between curvature emission and conventional cyclotron emission. This approach, although formally correct, has limited applicability and ignores two important features of the emission mechanism. The first is that, in adopting a plane-wave formalism, the interaction length for an individual electron, $\approx R_c/\gamma_b$, was essentially coextensive with the region over which the waves could interact with an electron. This approach precludes a strong amplification under all circumstances because the wave would have grown substantially during a single interaction. The second problem was that a dispersion of the waves was neglected. The first shortcoming can be addressed by expanding the electromagnetic field in cylindrical waves centered on $r = 0$ and the second explicitly by considering general plasma modes.

It is very important that the above mentioned authors do not take into account the drift motion of the particles.

3 Waves in the Pulsar Magnetosphere

The linear collective properties of an electron–positron plasma are now well established (Benford & Bushauer 1977; Volokitin, Krasnosel'skikh, & Machabeli 1985; Melrose 1986; Arons & Barnard 1986). According to the theory, for oblique propagation with respect to the external magnetic field (directed along the z axis), there are three normal modes. One is a purely transverse extraordinary mode (X-mode) with dispersion in the laboratory frame

$$\omega_x = kc\left(1 - \frac{1}{8}\frac{\omega_p^2}{\omega_B^2}\frac{1}{\gamma_0^3}\right), \tag{7}$$

where $\omega_p = [4\pi(n_{e^-} + n_{e^+})e^2/m]^{1/2} = [8\pi ne^2/m]^{1/2}$ is the combined plasma frequency (i.e. taking into account contributions of electrons and positrons), ω_B is the cyclotron frequency and γ_0 is the Lorentz-factor of plasma particles moving along the field lines. Below, we consider strongly magnetised plasma, $\omega_p \ll \omega_B/\gamma_0$. The second and third modes are of mixed longitudinal–transverse character. The lower frequency mode is analogous to the Alfvén wave, and the higher frequency mode to the fast (superluminal), $v_{ph}c$, ordinary mode (O-mode). Analytical expressions for dispersion of these modes are available in some limits. We consider the case $kc \ll \sqrt{2}\omega_p$, for waves propagating almost parallel to the magnetic field $|\vec{k}_\perp| \ll k_z$. For the O-mode we have

$$\omega_0 \simeq \frac{\omega_p^2}{\gamma_0^3} + 3k_z^2c^2 + |\vec{k}_\perp|^2c^2. \tag{8}$$

Here we do not consider the Alfvén mode, and so do not specify its dispersion. However, we note that for parallel propagation there is a coupling point $\omega_p \approx \omega_0 = k_0c$, where all three modes are indistinguishable (in a cold plasma), and proper consideration of their non-linear properties must take this into account. The electric field of the X-mode is perpendicular to the plane of vectors $\vec{k}$ and $\vec{B}$; the electric fields of the O-mode and Alfvén mode are in the plane. The oblique subluminal Alfvén mode is strongly suppressed due to Landau damping if its phase speed, effectively $v_A/(1+v_A^2/c^2)^{1/2}$, is less than the speed of the bulk of the particles; in the opposite limit Alfvén waves are weakly damped.

Low frequency modes analogous to the ion–acoustic wave in an electron–ion plasma are absent in an electron–positron plasma. Thus when considering non-linear effects in the wave propagation, the only possibility for amplitude modulations of the O-mode is due to resonant excitation of a beat wave. This was pointed out in Pataraya & Melikidze (1980), see also Pelletier, Sol, & Asseo (1988), and Asseo & Pelletier (1990). In the cited papers, the possibility of amplitude modulation of the purely parallel O-mode (which is often called the Langmuir mode) was considered for $kc \ll \omega_p$; the beat wave is generated as a result of interaction of two Langmuir waves with close frequencies $|\omega_L - \omega_L'| \ll \omega_p$. However, in an electron–positron plasma, with equal densities of electrons and positrons (as in Pataraya & Melikidze 1980; Pelletier et al. 1988; Asseo & Pelletier 1990), this second-order process cancels because of the equal masses and opposite charges of the plasma particles. This is because the second-order non-linear current is proportional to the charge cubed, and the electron and positron contributions are equal and opposite. Furthermore, when considering interaction of waves under the condition $|\omega_L - \omega_L'| \ll \omega_p$, the beat wave cannot be generated in the superluminal $|\omega_L - \omega_L'| \gg |k_z - k_z'|c$ range of phase velocities because of the wave dispersion. And, finally, when considering waves near the coupling point $\omega_L \approx \omega_0 = k_0c$, one needs to invoke the non-linear interactions with the X-mode and the Alfvén mode. This

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possibility was studied in Amiranashvili & Ignatov (1995), where self-similar unstable solutions satisfying the non-linear Schrödinger equation were found. An analogous problem was considered in (9) where it was demonstrated that small transverse perturbations lead to unstable solutions.

4 Modulation Instability

Consider the possibility of modulations of the fast O-mode by transverse waves. Note that by 'transverse' we imply not only the X-mode, but also the high frequency (compared with ω_p) O-mode, where its dispersion is close to the vacuum case. The difference between these modes in this case is only in their polarisation. In the interaction, we are interested in the longitudinal superluminal component of the perturbation appearing as a result of the interaction of two transverse waves:

$$\left|\frac{\omega^t - \omega'^t}{k_z - k'_z}\right| c. \tag{9}$$

Using $|\vec{k}| \simeq k_z(1 - \vec{k}_\perp^2/2k_z^2)$, this inequality implies

$$1 - \frac{1}{2}\left[\frac{\vec{k}_\perp^2}{k_z(k_z - k'_z)} - \frac{\vec{k}'^2_\perp}{k'_z(k_z - k'_z)}\right] 1. \tag{10}$$

To satisfy the latter for $k_z k'_z$, we require $|\vec{k}'_\perp||\vec{k}_\perp|$ as well as

$$\frac{|\vec{k}'_\perp|}{k_z}\,\frac{|\vec{k}'_\perp| - |\vec{k}_\perp|}{k_z - k'_z} 1. \tag{11}$$

We note that there is an important qualitative difference between a pair plasma and the more familiar (and more studied) electron–ion plasma when wave–wave and wave–particle interactions are considered. In an electron–ion plasma, the density fluctuations associated with Debye shielding can produce electric dipole radiation when forced to oscillate. This effect has no counterpart in a pair plasma because the electrons and positrons oscillate out of phase. As a consequence, the non-linear shielding (Melrose 1986), which tends to dominate wave–wave interactions in an electron–ion plasma, is absent, and Thomson scattering, which is the same for electrons and positrons, becomes the dominant non-linear effect. Thus the process of reradiation of a wave by a particle in such a plasma can be considered on an isolated electron (positron), similar to Thomson scattering when an electromagnetic wave t forces oscillations of the particle generating a wave t′.

Let us investigate the possibility of a non-linear (modulational) instability of an ordinary mode (O-mode) having phase velocity in the broad range of velocities exceeding the speed of light (Machabeli, Vladimirov, & Melrose 1999). We propose the mechanism of the non-resonant generation of the O-mode by two high frequency transverse waves t and t′, propagating along the external magnetic field in the opposite directions. The electric field vectors $\vec{E}$ and $\vec{E}'$ of these waves are perpendicular to the magnetic field, and intensive interaction with parallel (longitudinal) perturbations is possible because of non-linear drift motions of plasma particles in the fields $\vec{E}$, $\vec{E}'$, and $\vec{B}_0$. This formulation is entirely different from that of Pataraya & Melikidze (1980) and Amiranashvili & Ignatov (1995).

We consider a strongly magnetised electron–positron plasma moving with relativistic velocity along the magnetic field lines. In our study, we use (unless the opposite is specified) the reference frame connected with the moving plasma. This assumption does not mean that the perpendicular motions of plasma particles (appearing as a result of interactions) are non-relativistic. Furthermore, we calculate the non-linear current using the assumptions of small amplitudes of the interacting waves as well as small ratios of plasma kinetic pressure to the magnetic pressure and plasma electron frequency to the frequencies of the two electromagnetic pump waves (the beat wave appearing as a result of their interaction modulates the considered longitudinal O-mode). We also assume that the amplitudes of the electromagnetic pump waves considerably exceed the amplitude of the longitudinal mode that is justified by our assumption that mostly transverse waves are generated in the pulsar magnetosphere through (anomalous) cyclotron resonance (Lominadze et al. 1979; Machabeli & Usov 1979). Using the above small parameters, we obtain a system of non-linear three-dimensional equations, which is solved analytically neglecting back reaction of the modulations on the pump waves. Generally, the non-linear system is unstable.

When the modulation-type instability occurs the non-resonant process pumps wave energy from the high frequency to the low frequency waves. As a result of the development of modulational instability not only $E_\parallel$, but also $E_\perp$ is generated (Machabeli, Vladimirov, Melrose, & Luo 2000). Non-resonant interaction between these fields with particles can be described by non-resonant quasilinear diffusion (for a detailed discussion, see e.g. Machabeli, Luo, Melrose, & Vladimirov 2000).

5 Non-resonant Quasilinear Diffusion

Non-resonant quasilinear diffusion (NQD) is driven by a non-resonant instability in the plasma, which is due to the back reaction on the distribution of particles to the transfer of free energy from the particles to the waves in the instability. That is, NQD is similar to the more familiar resonant quasilinear diffusion (RQD), i.e. the back reaction on the particles of emission or absorption of waves owing to a resonant wave–particle interaction. In a non-relativistic plasma, a familiar non-resonant instability is the 'garden-hose' instability considered by several authors (e.g. Chandrasekhar, Kaufman, & Watson 1958; Lominadze et al. 1979). NQD in the non-relativistic case has been studied in detail (e.g. Davidson 1972; Shapiro & Shevchenko 1968). Volokitin et al. (1985) and Machabeli et al. (2000a,b) considered the relativistic generalisation of NQD theory. NQD tends to stabilise a non-resonant

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instability by modifying the distribution function of the particles in such way as to reduce the growth of the instability. For any instability, such as the garden-hose instability, driven by an excess (compare to an isotropic distribution) of parallel over perpendicular energy, the redistribution results in an increase in the perpendicular energy. The perpendicular energy transferred to the particles is radiated away as synchrotron radiation, producing the observed γ-rays.

The main ingredient of the model proposed by Machabeli et al. (2000b) is that a non-resonant instability is driven by the anisotropy in the particle distribution. There is a strong argument that this must occur. As a result of synchrotron losses the escaping pairs are expected to be all in their lowest Landau orbital as they leave the inner magnetosphere. It is known that such a one-dimensional distribution is garden-hose unstable in a pulsar magnetosphere, but this instability is ineffective inside the light cylinder (e.g. Gedalin, Melrose, & Gruman 1998). Provided some other non-resonant instability driven by the anisotropy does develop inside the light cylinder, the theory developed by Machabeli et al. (2000b) is applicable. A specific non-resonant instability, owing to a three-wave interaction, that satisfies this requirement was discussed recently by Machabeli et al. (2000a,b), and it is this instability that is invoked in the model. It results in superluminal Langmuir-like waves, which cannot be resonantly damped by the particles. The free energy is in the parallel motion of the particles, and this is transferred partly to these waves and partly to the perpendicular energy through NQD.

6 Conclusion

To conclude, we argue that t-modes generated in a pulsar magnetosphere can create beat density modulations along the magnetic field. When the modulation frequency is much less than the frequency of the generated field perturbations, the growth of the parallel potential field is accompanied by the growth of the transverse electromagnetic field. The energy of the modulations can then be converted into perpendicular momentum of the non-resonant particles and this in turn leads to the synchrotron radiation at high photon energies. Direct application of this mechanism to observational data includes case by case analysis of concrete results and is the subject of further investigations.

The mechanism of γ-radiation is synchrotron. Synchrotron radiation is characterised by a power diagram

$$P_\nu = P_{\nu_0}\frac{1}{1-\frac{v}{c}\cos\alpha}. \qquad (12)$$

P_{ν_0} is the synchrotron radiation power in the rest particle's frame. As for the frequency

$$\nu = \nu_0 = \frac{\left(1-\frac{v^2}{c^2}\right)^{1/2}}{1-\frac{v}{c}\cos\alpha}, \qquad (13)$$

$\nu_0 = \frac{3}{2}\omega_B\gamma_\perp^2$ is the synchrotron radiation frequency in the rest frame, ν is the frequency in the observer frame, and α is the angle between radiation propagation direction and pulsar magnetic field.

When $\alpha = 0$, $\nu = 2\nu_0\gamma$ and $P = P_{\nu_0}\gamma^2$. But when $\alpha = \pi/2$, $\nu = \nu_0/2$ and $P = P_{\nu_0}$. So, for young pulsars, which have hard γ-radiation, there exists the principal possibility of observing radio emission (if the latter exists at all) between the main pulse and interpulse.

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