Pinch instability: an alternative mechanism for pulsar coherent radio radiation

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Abstract

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Full Text

Preamble

Pinch Instability: An Alternative Mechanism for Pulsar Coherent Radio Radiation

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Abstract

The origin of pulsed radio emissions from pulsars has long remained a subject of enduring and intricate perplexity. A common viewpoint posits that the production of this coherent radio radiation necessitates robust electric currents. Noting that pinch instability is likely to occur concomitantly with the emergence of intense electric currents, this manuscript delves into an exploration of the conditions under which pinch instability develops within pulsar magnetospheres. We find that, under typical pulsar parameters, pinch instability can develop effectively, and the charged bunches necessary for pulsed coherent radio radiation would subsequently be generated as a result of this instability.

Subject headings: pulsars, radio pulsars, radio astronomy

Introduction

Due to the effect of unipolar induction and the excellent electrical conductivity of stars, pulsar magnetospheres should be replete with charged particles [?]. As indicated by the existence of pulsar wind nebulae [?] and the requirement to produce pulsed electromagnetic radiation [?], pulsars must constantly inject a large amount of charged particles into their surroundings. Therefore, to maintain the plasma environment, pulsars and their magnetospheres must be able to continuously generate charged particles. Such a dynamic magnetosphere may allow the development of various plasma instabilities, for example, the widely discussed two-stream instability.

For pulsar physics, the two-stream instability has been invoked to explain pulsar pulsed radio radiation [?]. However, it remains uncertain whether this type of instability can be effectively developed within pulsar magnetospheres [?]. A newly-developing and compelling scenario proposes that the fluctuating process involving intermittent sparks of electron-positron production can directly drive coherent radio radiation from pulsars \cite{Philippovetal2020, Mestel_{etal1985}, Beloborodov2008, Timokhin&Arons2013}. While this charge-fluctuation scenario has not been fully fleshed out[1], the usually implicit assumption that coherent curvature radiation emanating from charged bunches formed through mechanisms such as the two-stream instability [?] no longer appears to be the only viable choice. Regardless of whether it is the aforementioned charge-fluctuation scenario or the bunch scenario, the fundamental requirements remain consistent: namely, electron-positron pair cascades should be initiated in a strong electric field. Hence, whether from observational or theoretical perspectives, strong electric currents flowing along pulsar magnetic fields ought to emerge within pulsar magnetospheres. Given that such a current will induce a toroidal magnetic component, it is natural to ponder: can pinch instabilities effectively develop within a pulsar magnetosphere? Furthermore, can this kind of instability segment these electric currents into the bunches necessary for coherent radio radiation?

In the following section, by employing the inner gap model [?], we demonstrate that the current bunches necessary for coherent curvature radiation in the radio band could be produced through pinch instabilities within pulsar magnetospheres, coincidentally, under typical pulsar parameters.

2.1 The Basic Physical Picture

Within the framework of the inner gap model, the electromotive force provided by unipolar induction must surpass the voltage required by sparks. Consequently, the thickness of the gap will exceed the mean free path for electron-positron pair conversions of the curvature photons emitted by initial charged particles. Therefore, secondary electron-positron pairs would be spatially separated by the electric field of the gap. As shown in Figure 1 [Figure 1: see

[1] See the comments by Alice K. Harding at https://physics.aps.org/articles/v13/96#c10.

original paper], supposing that electron-positron pair cascades in a pulsar magnetosphere result in an electric current with density $\mathbf{j}$ flowing along the pulsar magnetic field with strength B_z, the current induces a toroidal magnetic component, B_θ. According to the principle of magnetic flux conservation, there is a rough critical condition for the shrinkage of the current that $B_\theta^2(r_0) = B_z^2/2$, where r_0 is the radius of the current. Since the current is gradually increasing during the cascading process, at least before a certain point in time, the two components of the magnetic field satisfy $B_\theta^2(r_0) < B_z^2/2$ at the boundary of the current. Therefore, the pinch instability is initially restrained. However, as the pulsar magnetic field may decay as $B_z \propto B_s(z = R_*)^{-3}$, where B_s is the magnetic field strength on the pulsar surface and R_* is the radius of the pulsar, the condition $B_\theta^2(r_0) > B_z^2/2$ could finally be satisfied in the pulsar magnetosphere, and then the pinch instability is able to develop.

When the current passes through a point where $B_\theta^2(r_0) = B_z^2/2$, the current begins to shrink. A disturbance of the current radius, $\delta(r_0; t)$, perpendicular to the pulsar magnetic field begins to grow, where t is the time. If the pinch instability can develop effectively, the value of $\delta(r_0; t)$ should achieve δr_0 during the crossing time L/c, where L is the length of the current and c is the speed of light. One of the standard methods to discuss the development of an instability is the normal mode analysis. For simplicity, we are only interested in the solution of $\delta(r_0; t)$ with the formalism:

$$\delta(r_0; t) \propto \exp(-i\omega_k t + ikz)$$

where ω_k is the angular frequency and k is the wave number. Considering a bulk of charged particles moving along the magnetic field approximately uniformly distributed in the direction perpendicular to the z-axis, the dispersion relation corresponding to equation (1) is given by [**?**]:

$$\omega_k^2 = \frac{2B_{0z}^2}{\mu_0 \rho r_0^2} - \frac{B_{0\theta}^2 k^2}{\mu_0 \rho} \frac{I_1(kr_0)}{I_0(kr_0)}$$

where μ_0 is the permeability of vacuum, ρ is the mass density of the current, B_{0z} is the pulsar magnetic field before the perturbation, $B_{0\theta}$ is the toroidal magnetic component induced by the current before the perturbation, and I_0 and I_1 are the modified Bessel functions with the subscripts being their orders.

2.2 Parameter Estimation

If the perturbation can be effectively developed, so that the current can be segmented into bunches, the relation $|\omega_k|L/c \gg 1$ ought to be satisfied. To ensure radiation from these bunches is coherent in the radio band, the bunch length should be, at least, $L < 0.1\,\mathrm{m}$. Consequently, the wave number of the perturbation is $k > 2\pi/L \approx 60\,\mathrm{m}^{-1}$.

Since $\mu_0\rho$ is a large value, according to equation (2), the condition $|\omega_k|L/c \gg 1$ will be soon satisfied after the current passes through the point where $B_\theta^2(r_0) = B_z^2/2$:

$$|\omega_k| \approx \frac{B_{0\theta}}{\sqrt{\mu_0\rho}} k \sqrt{\frac{I_1(kr_0)}{I_0(kr_0)}} \gg \frac{c}{L}$$

As $B_{0z} \approx B_s(R_*/z)^3$ and $B_{0\theta} \approx \frac{1}{2}\mu_0 j r_0$, with $j = \eta\rho_{GJ}c$ where ρ_{GJ} is the Goldreich-Julian charge density, equation (3) can be rewritten as:

$$\eta \approx \frac{2B_s R_*^3}{\mu_0\rho_{GJ}cr_0z^3}\frac{kI_0(kr_0)}{I_1(kr_0)}$$

where η is the charge density of the bunch. Taking typical pulsar parameters, where the rotation period $P = 1\,\mathrm{s}$, the magnetic field strength on the surface $B_s = 10^{12}\,\mathrm{G}$, and the radius $R_* = 10^6\,\mathrm{cm}$, the numerical dependence between η and r_0 is depicted in Figure 2 [Figure 2: see original paper].

As shown in Figure 2, if we adopt $\eta = 10^5$, the same as the usually optimistic expectation of the cascade pair multiplicity \cite{Timokhin&Harding2015, deJager_{etal1996}, deJager2007, Bucciantinietal2011}, the radius of the current should be $> 10^{-2}\,\mathrm{cm}$ under $z = 5 \times 10^7\,\mathrm{cm}$. The width of the current can be contributed by two ways: (i) the beam width of the curvature radiation from charge particles; (ii) the broadening associated with the conversion of photons emitted via synchrotron radiation of secondary charged particles.

The beam width of curvature radiation can be estimated as $\approx l_{cr}/\Gamma = 10^{-2}l_{cr,4}\Gamma_6\,\mathrm{cm}$, where $l_{cr} = l_{cr,4}10^4\,\mathrm{cm}$ is the mean free path of the electron-positron pair conversion of the initial curvature photons, and $\Gamma = \Gamma_6 10^6$ is the Lorentz factor of the initial charged particles.

For secondary charged particles, the broadening of the current caused by the gyration radius of these particles is so minuscule that it is disregarded in the subsequent estimation. The primary contribution to the broadening is the mean free path of the conversion of photons emitted via synchrotron radiation, which is given by [**?**, **?**]:

$$l_{sr} \approx \frac{1}{\alpha B_{12}}\frac{\hbar}{m_{ec}}\exp\left(\frac{4}{3\chi}\right) = 10^{-6}B_{12}^{-1}\exp\left(\frac{4}{3\chi}\right)\,\mathrm{cm}$$

where α is the fine-structure constant, $\hbar$ is the reduced Planck constant, m_e is the rest mass of electrons, $B_q = 4.4 \times 10^{13}\,\mathrm{G}$, ϵ_{ph} is the energy of the synchrophotons, $\chi = \epsilon_{ph}/(m_{ec}^2)$, and $B = B_{12}10^{12}\,\mathrm{G}$ is the effective magnetic field strength for the conversion. According to equation (5), to keep $r_0 \approx 10^{-2}\,\mathrm{cm}$, there should be $\epsilon_{ph} > 8m_{ec}^2$. Since the peak energy of synchrotron radiation is

$0.3\hbar\omega_c$, where $\omega_c = 3eB\gamma_\pm^2/m_e$ is the peak frequency with e being the elementary charge, the effective Lorentz factor for synchrotron radiation of secondary positron-electron pairs should be, at least:

$$\gamma_\pm > 17.8\frac{eB\hbar}{m_{ec}} \approx 750$$

For $B \approx B_s = 10^{12}\,\mathrm{G}$, there is $\gamma_\pm > 750$. By considering the radiation reaction of curvature radiation, the acceleration of electrons and positrons in the gap is approximately:

$$m_{ec}^2\frac{d\Gamma}{dz} = eE_\parallel - \frac{2e^2\Gamma^4}{3R_c^2}$$

where $E_\parallel$ is the electric field intensity of the gap, and R_c is the radius of curvature of the magnetic field lines on the pulsar surface. For a typical gap with potential $\approx 10^{12}\,\mathrm{V}$, electrons and positrons cannot achieve the saturation Lorentz factor, Γ_{sat}, determined by $eE_\parallel = 2e^2\Gamma_{sat}^4/3R_c^2$. Therefore, the Lorentz factor of the primary electrons and positrons can be estimated as $\approx eE_\parallel l_{cr}/(m_{ec}^2) \approx 10^7$, and the corresponding energy of the typical curvature photons can be up to $\approx 10^5 R_{c,6}^{-1}\,\mathrm{MeV}$, where $R_{c,6} = R_c/10^6\,\mathrm{cm}$. The effective Lorentz factor for synchrotron radiation of the electron-positron pairs converted from the initial curvature photons is $\approx 10^5(l_{cr}/R_c) = 10^3 l_{cr,4}R_{c,6}$. Hence, the condition $\gamma_\pm > 750$ can be achieved by most secondary particles. According to the exponential term in equation (5), the distribution of charged particles converted from synchrophotons emitted by secondary positron-electron pairs with momentum perpendicular to the pulsar magnetic field $< 750m_{ec}$ is unconsolidated, so the contribution of these particles to the current is negligible. Therefore, it is possible to achieve a dense current with radius $\approx 10^{-2}\,\mathrm{cm}$. Consequently, provided that the spatial distribution of the initial particles for sparks and the sparks themselves are inhomogeneous, it is reasonable to expect the generation of a series of snatchy currents in the magnetosphere.

2.3 The Stretch of Bunches

In the preceding subsection, we have shown that, under some typical values of pulsar parameters ($B_s = 10^{12}\,\mathrm{G}$, $P = 1\,\mathrm{s}$, $R_* = 10^6\,\mathrm{cm}$, $\eta = 10^5$, and $r_0 > 10^{-2}\,\mathrm{cm}$), the bunches necessary for coherent radio radiation (at $z \approx 5\times10^7\,\mathrm{cm}$) can be formed due to pinch instability. However, velocity dispersion, Δv, of the charged particles in a bunch will stretch the bunch, so that the superposition of phase positions of radio waves emitted by the bunch may become incoherent. According to the relation between velocity v and Lorentz factor γ, where $\gamma = 1/\sqrt{1-(v/c)^2}$, we simply have $\Delta v \approx \Delta\gamma c/\gamma^3$, where $\Delta\gamma$ is the dispersion of Lorentz factors of the particles in the bunch. Considering a maximum radiation altitude $z_M \approx 10^8\,\mathrm{cm}$, the stretch $z_M\Delta\gamma/\gamma^3 < 10\,\mathrm{cm}$ gives $\Delta\gamma < 10$ under

$\gamma \approx 800$. Even though the distribution function of particles in the bunch is unknown, the nearly monochromatic requirement seems impractical.

Two elements may relax this difficulty. (i) Although the Lorentz factors of the charged particles within a bunch can span a broad range, the bunch can maintain a dense core, provided that the Lorentz factors of the majority of particles are confined within a relatively narrow range. (ii) As the strength of the pulsar' s magnetic field diminishes, pinch instability has the potential to recur, as long as the dense core of the bunch retains sufficient density.

As shown in Figure 2, if the maximum radiation altitude is $\approx 10^8$ cm, when a bunch with $\eta \approx 10^4$ reaches this point, the shrinkage of the bunch can still grow effectively. Therefore, even in the presence of velocity dispersion and a notable decline in the number density of particles within the bunch over time (from $\eta \approx 10^5$ to $\eta \approx 10^4$), pinch instability can still occur due to the effective decay of the pulsar magnetic field. Comparing $\eta \approx 10^5$ at $z = 5\times10^7$ cm and $\eta \approx 10^4$ at $z = 10^8$ cm, even if particles are evenly distributed within a range of $\Delta\gamma \approx 100$, it remains feasible to form appropriate bunches. This significantly eases the aforementioned difficulty. From an observational perspective, the brightness temperature of the radio radiation is expected to undergo a significant reduction at high altitudes.

Summary and Discussions

In this manuscript, we discussed the effect of pinch instability on pulsar pulsed radio radiation. We found that under typical values of pulsar-related parameters, pinch instability can result in the formation of current bunches which may be the source of coherent radio radiation.

The above discussion is presented within the framework of the inner gap model [**?**]. Under the slot gap model [**?**], the pinch instability is still able to develop provided that the current is sufficiently intense. This could be the difference between the mechanism discussed in this manuscript and the work of Philippov et al. [**?**]. Because, if coherent radio radiation is triggered by charge fluctuations, under the inner gap model, the radio radiation ought to occur near the pulsar surface, which may conflict with reported high-altitude radio emissions \cite{Mitra2017, Sun_{etal2025}}.

Clarifying this issue holds fundamental significance for pulsar physics due to the treatment of the binding energies of electrons and ions in pulsar surfaces. If both the charge-fluctuation scenario and the reported high-altitude radio radiation are reliable, the inner gap model should only be viable for bare strange stars [**?**]. Even so, we still require an additional channel to produce potential high-altitude radio radiation emanating from bare strange stars, such as the mechanism proposed in this manuscript.

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References

Arons, J. & Scharlemann, E. T. 1979, ApJ, 231, 854. doi:10.1086/157250
Beloborodov, A. M. 2008, ApJ, 683, L41
Bucciantini, N., Arons, J., & Amato, E. 2011, MNRAS, 410, 1, 381. doi:10.1111/j.1365-2966.2010.17449.x
de Jager, O. C., Harding, A. K., Michelson, P. F., et al. 1996, ApJ, 457, 253. doi:10.1086/176726
de Jager, O. C. 2007, ApJ, 658, 2, 1177. doi:10.1086/511950
Erber, T. 1966, Reviews of Modern Physics, 38, 4, 626. doi:10.1103/RevModPhys.38.626
Galtier, 2016, doi:10.1017/CBO9781316665961 Cambridge, Cambridge University Press,
Goldreich, P. & Julian, W. H. 1969, ApJ, 157, 869
Melrose, D. B. 2017, Reviews of Modern Plasma Physics, 1, 1, 5. doi:10.1007/s41614-017-
Mestel, L., Robertson, J. A., Wang, Y.-M., et al. 1985, MNRAS, 217, doi:10.1093/mnras/217.3.443
Mitra, D. 2017, Journal of Astrophysics and Astronomy, 38, 3, 52. doi:10.1007/s12036-017-
Philippov, A., Timokhin, A., & Spitkovsky, A. 2020, Phys. Rev. Lett., 124, 24, 245101. doi:10.1103/PhysRevLett.124.245101
Rees, M. J. & Gunn, J. E. 1974, MNRAS, 167, 1. doi:10.1093/mnras/167.1.1
Ruderman, M. A. & Sutherland, P. G. 1975, ApJ, 196, 51
Shukre, C. S. & Radhakrishnan, V. 1982, ApJ, 258, 121. doi:10.1086/160058
Sturner, S. J. 1995, ApJ, 446, 292. doi:10.1086/175788
Sturrock, P. A. 1971, ApJ, 164, 529. doi:10.1086/150865
Sun, S. N., Wang, N., Yan, W. M., et al. 2025, ApJ, 983, 2, 179. doi:10.3847/1538-4357/adc254
Timokhin, A. N. & Arons, J. 2013, MNRAS, 429, 1, 20. doi:10.1093/mnras/sts298
Timokhin, A. N. & Harding, A. K. 2015, ApJ, 810, 2, 144. doi:10.1108/0004-637X/810/2/144
Usov, V. V. 1987, ApJ, 320, 333. doi:10.1086/165546
Xu, R. X., Qiao, G. J., & Zhang, B. 1999, ApJ, 522, 2, L109. doi:10.1086/312226
Zhang, B. & Harding, A. K. 2000, ApJ, 532, 2, 1150. doi:10.1086/308598

Note: Figure translations are in progress. See original paper for figures.