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Some astrophysical effects of nonlinear vacuum electrodynamics in the magnetosphere of a pulsar

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ABSTRACT

In this study, we consider the propagation of hard electromagnetic emissions in the magnetosphere of a pulsar based on General Relativity and nonlinear vacuum electrodynamics. We show that the radiation will propagate at different velocities in the magnetosphere of a pulsar and form two normal modes, which are polarized in mutually orthogonal planes. We calculate the delay between the two orthogonal modes as they propagate from the pulsar to the detection device.

1. Introduction

Theory [1] and experiments [2] show that the electrodynamics are nonlinear in a vacuum. Therefore, studies [3–10] of various manifestations of nonlinearity in electrodynamics are of undoubted interest.

A number of recent studies [11–19] have considered various effects of nonlinear vacuum electrodynamics as well as the possibility of their measurement in the laboratory. However, the magnetic fields that can be created in the laboratory $B \sim 10^5$ G are significantly smaller than the quantum value $B_q = m^2c^3/(e\hbar) = 4.41 \times 10^{13}$ G, so their observation will only be possible in the future after further developments in measurement technology. Therefore, at present, the main focus is on the astrophysical effects of nonlinear vacuum electrodynamics that occur in the magnetic fields $B \sim B_q$ of pulsars and collapsars.

The effects of nonlinear corrections of the vacuum electrodynamics on the polarization and directivity of the radiation from X-ray pulsars were first studied in [20–22]. These calculations showed that an electromagnetic wave passing through the magnetic field of a pulsar should experience nonlinear electrodynamic birefringence, where it should be split into two normal modes with mutually orthogonal polarization to be distributed via non-coincident rays at different speeds. However, in the magnetosphere of a pulsar, X-rays should also experience birefringence due to the presence of plasma.

As shown in [22], at a certain concentration of plasma, the birefringence caused by nonlinear vacuum electrodynamics exceeds that induced by plasma. For gamma radiation, due to its higher frequency, the distorting effect of the plasma is not noticeable. Therefore, information about the distribution and density of the field in the emitting region can be obtained by studying the polarization states of the hard radiation of pulsars [22].

These studies were continued in [23–30]. For some special cases (the propagation of an electromagnetic wave in the plane of the magnetic equator and the magnetic meridian of pulsars) [27–29], it was shown that the main nonlinear electrodynamic effect, which can be registered on Earth, differs between the velocities of the normal modes in the magnetic field of a pulsar [29,30]. However, in previous studies [24,25,27], these calculations were performed for only a few simple cases.

Therefore, we aim to calculate this effect in the most general case where an electromagnetic pulse passes through the magnetic field of a pulsar in a random direction. We assume that at some point in the magnetosphere or on the surface of a pulsar, there is a relatively short burst of hard radiation. We also assume that the electromagnetic pulse generated during this event is either unpolarized or it has elliptical polarization [21].

In the strong magnetic field of a pulsar, the radiation pulse splits into two pulses due to birefringence, which are polarized in mutually perpendicular planes with different velocities. Therefore, the two electromagnetic pulses emitted at the same time from the same source will arrive at the recording device installed on a near-Earth

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satellite over different beams at different times $t_2 \neq t_1$. This device will first register the arrival of the front part of a more rapid pulse and the polarization of the detected radiation will be linear over time $t_2 - t_1$. After this period of time, the front part of the second pulse with orthogonal polarization will arrive at the recorder. Therefore, the further polarization of the total momentum will be arbitrary. To observe this effect, electromagnetic radiation detectors for pulsars must be equipped with devices that can measure the polarization state of the radiation.

The manifestation of nonlinear electrodynamic birefringence in the general case where a relatively short electromagnetic pulse passes the magnetic field of a pulsar in an arbitrary direction, as well as the calculation of the period $t_2 - t_1$, are obtained for the first time in the present study.

2. Equations for nonlinear vacuum electrodynamics and gravitation

We consider the nonlinear post-Maxwell electrodynamics, which are a direct consequence of quantum electrodynamics [1]. The Lagrangian [31] has the form

$$L = \frac{\sqrt{-g}}{32\pi}\{2I_2 + \xi[(\eta_1 - 2\eta_2)I_2^2 + 4\eta_2 I_4]\} - \frac{\sqrt{-g}}{c} j^\beta A_\beta,$$

where j^β is the current density four-vector, g is the determinant of the metric tensor, $\xi = 1/B_q^2$, $I_2 = F_{\beta\sigma}F^{\sigma\beta}$, and $I_4 = F_{\beta\sigma}F^{\sigma\nu}F_{\nu\mu}F^{\mu\beta}$ are invariants of the electromagnetic tensor $F_{\beta\sigma}$, and according to quantum electrodynamics, $\eta_1 = e^2/(45\pi\hbar c) = 5.1\times 10^{-5}$, $\eta_2 = 7e^2/(180\pi\hbar c) = 9.0\times 10^{-5}$.

The field equations derived from this Lagrangian have the form

$$\frac{1}{\sqrt{-g}}\frac{\partial}{\partial x^\beta}\{\sqrt{-g}Q^{\sigma\beta}\} = -\frac{4\pi}{c} j^\sigma, \tag{2.1}$$

$$Q^{\sigma\beta} = 8\pi\frac{\partial L}{\partial F_{\beta\sigma}} = \{1 + \xi(\eta_1 - 2\eta_2)I_2\}F^{\sigma\beta} + 4\xi\eta_2 F^{\sigma\nu}F_{\nu\mu}F^{\mu\beta}.$$

The second pair of electrodynamics equations agrees with the corresponding equations from Maxwell's theory

$$\frac{\partial F_{\mu\beta}}{\partial x^\nu} + \frac{\partial F_{\beta\nu}}{\partial x^\mu} + \frac{\partial F_{\nu\mu}}{\partial x^\beta} = 0. \tag{2.2}$$

The metric tensor in Eq. (2.2) satisfies Einstein equations [32]

$$R_{\beta\sigma} - \frac{1}{2}g_{\beta\sigma}R = -\frac{8\pi G}{c^4}T_{\beta\sigma}, \tag{2.3}$$

where $R_{\beta\sigma} = R^\nu_{\beta\sigma\nu}$ is the Ricci tensor and $T_{\beta\sigma}$ is the energy–momentum tensor of the matter and all fields, including electromagnetic. The system of equations (2.1)–(2.3) in our problem is addressed by the method of successive approximations with a precision linear in the small dimensionless parameters: the gravitational potential and post-Maxwell amendments. The gravitational field of the pulsar is assumed to be spherically symmetric and in the harmonic Fock coordinates [32], the metric will be expanded in the small parameter α/r with the required accuracy

$$g_{00} = 1 - \frac{2\alpha}{r}, \quad g_{rr} = -1 - \frac{2\alpha}{r},$$
$$g_{\theta\theta} = r^2 g_{rr}, \quad g_{\phi\phi} = g_{\theta\theta}\sin^2\theta, \tag{2.4}$$

where $\alpha = \gamma M/c^2$, γ is a gravitational constant, and M is the mass of the pulsar.

Suppose that at time $t = 0$, a hard radiation impulse is emitted from the point $\mathbf{r} = \mathbf{r}_0$ in the pulsar magnetosphere. Then, because of birefringence in the magnetic field of the pulsar, the impulse will split [33] into two impulses with orthogonal polarizations, which move at different speeds.

For convenience, we introduce the following spherical coordinate system. Consider a beam for the first normal mode and draw a tangent to it at the point $\mathbf{r} = \mathbf{r}_0$. The axis of the spherical coordinate system will be directed such that the tangent to the chosen beam and the center of the pulsar will lie in the same plane, and $\theta = \pi/2$, and the azimuthal coordinate ϕ of the source of hard radiation will be equal to $\phi = 0$.

Without loss of generality, we assume that in this coordinate system, the vector of the magnetic dipole moment of the pulsar $\mathbf{m}$ is directed to a point with spherical coordinates θ_0 and ϕ_0. Then, the Cartesian components of the magnetic dipole moment $\mathbf{m}$ take the form

$$m_x = |\mathbf{m}|\sin\theta_0\cos\phi_0, \quad m_y = |\mathbf{m}|\sin\theta_0\sin\phi_0, \quad m_z = |\mathbf{m}|\cos\theta_0.$$

As is accepted [34] in problems of celestial mechanics, instead of the radial coordinate r, we introduce the coordinate $u = 1/r$. Then, the non-zero components of the dipole electromagnetic field tensor of the pulsar in the coordinate system u, θ, ϕ, with the required accuracy for our purposes will be

$$\begin{aligned}
F_{u\theta} &= -F_{\theta u} = |\mathbf{m}|\sin\theta_0\sin(\phi - \phi_0),\\
F_{u\phi} &= -F_{\phi u} = |\mathbf{m}|\sin\theta[\sin\theta_0\cos\theta\cos(\phi - \phi_0) - \sin\theta\cos\theta_0],\\
F_{\phi\theta} &= -F_{\theta\phi} = 2|\mathbf{m}|u\sin\theta[\sin\theta_0\sin\theta\cos(\phi - \phi_0) + \cos\theta\cos\theta_0].
\end{aligned} \tag{2.5}$$

In electrodynamics, the eikonal method [32,33,35,36] is used when solving most problems, which allows us to study the motion of electromagnetic impulses based on their beams. Applications of this method to nonlinear electrodynamics have shown [33,37] that the propagation of electromagnetic waves in external electromagnetic and gravitational fields in nonlinear electrodynamics with field equations (2.1)–(2.3) is equivalent to the propagation of the normal modes through the isotropic geodesics in effective space-time, for which the metric tensor $G^{eff(1,2)}_{\nu\mu}$ has the form

$$G^{eff\ (1,2)}_{\nu\mu} = g_{\nu\mu} - 4\eta_{(1,2)}\xi F_{\nu\beta}g^{\beta\sigma}F_{\sigma\mu}. \tag{2.6}$$

Therefore, studies of the laws of propagation for electromagnetic impulses in the magnetic (2.5) and gravitational (2.4) fields of a pulsar can be performed conveniently not by using Eqs. (2.1)–(2.2), but based on the analysis of isotropic geodesics in space-time with the metric tensor (6).

Let us substitute Expressions (2.4) and (2.5) into (2.6), and write the components of the metric tensor of the effective space-time $G^{eff\ (1,2)}_{\nu\mu} \equiv G^{(1,2)}_{\nu\mu}$ explicitly as

$$\begin{aligned}
G^{(1,2)}_{00} &= 1 - 2\alpha u,\\
G^{(1,2)}_{uu} &= -\frac{1+2\alpha u}{u^4} - 4\mathbf{m}^2\xi\eta_{1,2}u^2\{\sin^2\theta_0\sin^2(\phi - \phi_0)\\
&\quad + [\sin\theta_0\cos\theta\cos(\phi - \phi_0) - \sin\theta\cos\theta_0]^2\},\\
G^{(1,2)}_{u\theta} &= 8\mathbf{m}^2\xi\eta_{1,2}u^3[\sin\theta_0\sin\theta\cos(\phi - \phi_0) + \cos\theta\cos\theta_0]\\
&\quad \times[\sin\theta_0\cos\theta\cos(\phi - \phi_0) - \sin\theta\cos\theta_0],\\
G^{(1,2)}_{u\phi} &= -8\mathbf{m}^2\xi\eta_{1,2}u^3[\sin\theta_0\sin\theta\cos(\phi - \phi_0) + \cos\theta\cos\theta_0]\\
&\quad \times \sin\theta\sin\theta_0\sin(\phi - \phi_0),\\
G^{(1,2)}_{\theta\theta} &= -\frac{(1+2\alpha u)}{u^2} - 4\mathbf{m}^2\xi\eta_{1,2}u^4\{\sin^2\theta_0\sin^2(\phi - \phi_0)\\
&\quad + 4[\sin\theta_0\sin\theta\cos(\phi - \phi_0) + \cos\theta\cos\theta_0]^2\},\\
G^{(1,2)}_{\theta\phi} &= -4\mathbf{m}^2\xi\eta_{1,2}u^4[\sin\theta_0\cos\theta\cos(\phi - \phi_0) - \sin\theta\cos\theta_0]\\
&\quad \times \sin\theta\sin\theta_0\sin(\phi - \phi_0),\\
G^{(1,2)}_{\phi\phi} &= -\Bigg\{\frac{(1+2\alpha u)}{u^2} + 4\mathbf{m}^2\xi\eta_{1,2}u^4\{[\sin\theta_0\cos\theta\cos(\phi - \phi_0)
\end{aligned}$$