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Radiation of fast radio bursts by hot neutron stars

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ABSTRACT

A scenario for formation of fast radio bursts (FRBs) is proposed. Just like radio pulsars, sources of FRBs are magnetized neutron stars. The appearance of a strong electric field in the magnetosphere of a neutron star is associated with the close passage of a dense body near the hot neutron star. Thermal radiation from the surface of the star (the temperature is of the order of 10^8K) causes evaporation and ionization of matter in the dense body. Ionized gas (plasma) flows around the magnetosphere of the neutron star with velocity $u \simeq 10^7$ cm s^{-1} and creates an electric potential $\psi_0 \simeq 10^{11}$V in the polar region of the magnetosphere. A small fraction of electrons from the plasma flow are accelerated toward the star and gain a Lorentz factor of $\simeq 10^5$. Thermal photons moving toward precipitating electrons are scattered by them and produce gamma photons with energies $\simeq 10^5\, m_e c^2$. These gamma quanta produce electron–positron pairs in collisions with thermal photons. The electron–positron plasma produced in the polar region of the magnetosphere accumulates in a narrow layer at the bottom of a potential well formed on one side by a blocking potential ψ_0 and on the other by the pressure of thermal radiation. The density of electron–positron plasma in the layer increases with time and after a short time the layer becomes a mirror for thermal radiation of the star. The thermal radiation in the polar region under the layer is accumulated for a time $\simeq 500$ s, then the plasma layer is ejected outside. The ejection is observed as a burst of radio emission formed by the flow of relativistic electron–positron plasma.

Key words: radiation mechanisms: general – stars: neutron.

1 INTRODUCTION

Fast radio bursts (FRBs), first discovered in 2007 by Lorimer et al. (2007), are short radio signals of several milliseconds duration ($\tau \simeq 10^{-3}$s) with an energy flux from tens of milliJansky to several Jansky in the radio band of $\simeq$ 1–10 GHz. The main feature of FRBs is a high observed dispersion measure (DM). It is several times greater than the maximum value of the dispersion measure of a radio signal passing through our Galaxy. High DM suggests that the radio signal propagates through the intergalactic medium. Also, the rotation measure (RM), connected with the presence of a magnetic field in a plasma through which a radio signal propagates, is high. The ratio of the rotation measure to the dispersion measure, RM/DM, equals the average value of the longitudinal magnetic field along the line of sight. In contrast, this ratio is one order of magnitude smaller than the average magnetic field in the Galaxy, which also indicates that the radio signal propagates in the intergalactic medium. Finally, the repeated source FRB 121102 was identified with a galaxy having redshift $z = 0.192$ (Chatterjee et al. 2017). Assuming that large values of DM and RM are achieved in the intergalactic medium implies cosmological distances to sources of FRB of $z \simeq 1$. If we assume isotropic radiation into a solid angle $\simeq 4\pi$ and a cosmological distance to sources, the radiated power in the radio band should reach a value of the order of 10^{43} erg s^{-1} and a total radiated energy of 10^{40}erg. The beaming of FRBs can decrease these estimates by many orders of magnitude (Katz 2017a). If we do not take the relativistic time delay in the source into account, then the source size l is estimated to be of the order of 3×10^7cm. Further, the energy density in the source is equal to 3×10^{17} erg cm^{-3} in the radio range only. This corresponds to an electric field $E \simeq 10^{12}$V cm^{-1}, which is at least two orders of magnitude higher than the value of the atomic field, $E_a \simeq 0.5 \times 10^{10}$V cm^{-1}. At such large energy densities, we have to expect radiation in optical, X-ray and gamma ranges, which is not observed: it is difficult to understand what the extragalactic source could be. The nature of FRBs therefore remains unknown today.

There are several different consequences regarding the origin of FRBs. First of all, because of smallness of l, the source of a FRB is probably a neutron star. It is known that neutron stars are observed as radio pulsars, X-ray pulsars, magnetars and rotating radio transients (RRATs). High brightness temperature of a FRB, up to 10^{37}K, suggests a coherent mechanism of radio emission similar

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to that of radio pulsars, especially the mechanism of radiation of giant radio pulses observed from several radio pulsars (Soglasnov et al. 2004; Hankins & Eilek 2007). When radio sources of this kind with the same radiation mechanism are placed at cosmological distances, it is necessary to assume extreme values of neutron star parameters: large magnetic fields at the surface, $B_0 \simeq 10^{15}$ G, and fast rotation, $\Omega_{\rm NS} \simeq 10^4\ {\rm s}^{-1}$. However, neutron stars with such extreme parameters lose their rotation energy quite rapidly, during a time of the order of 10^3s, and cannot give repeated bursts. As for magnetars, in particular soft gamma repeaters (SGRs), the energy is stored in a magnetic field inside the star and they produce bursts in the gamma and X-ray energy ranges, which are not observed in FRBs. The energy released in a burst can be the energy of the electric field in the rotating magnetosphere of a neutron star (Katz 2017b). Katz calls such a phenomenon 'pulsar lightning'. However, it is not clear how and how fast transition from one configuration of the magnetosphere to another is possible. Earlier, the term 'lightning' was used for explanation of the radiation of RRATs (Istomin & Sobyanin 2011a,b). It was specifically shown that cascade formation of electron–positron pairs and their acceleration by a longitudinal electric field in the polar vacuum magnetosphere, caused by absorption of an energetic gamma quantum from Galactic and extragalactic backgrounds, should look like a flash of lightning on Earth during a thunderstorm.

The closest situation to that discussed in this article is provided by models of FRB origin resulting from the interaction of neutron stars with other bodies (planets, comets, asteroids). These are primarily direct collisions of bodies with a neutron star (Di & Dai 2017; Dai et al. 2016) and the interaction of a relativistic pulsar wind with the companion of a neutron star (Mottez & Zarka 2014).

However, apparently FRBs constitute a different class of source of radio emission, neutron stars. This can be seen from analysis of observations of the source FRB 121102.

2 FRB 121102

A large energy density in a source most likely suggests a cataclysmic event, such as an explosion and hence destruction of the source. Indeed, no recurrent events were recorded until recently. One radio source, FRB 121102, discovered in 2012 November, flared up ten times within 16 days (1.4×10^6s) 926 days later (2015 May). With the exception of one long time interval between consecutive flashes, also about 16 days, the time intervals between bursts were random, from $\simeq$ 20 s to $\simeq$ 1 h. The average duty cycle was $\simeq$ 670 s $\simeq$ 11 min (Spitler et al. 2016). Then, after 164 days (2015 November), six more bursts were recorded during 25 days, 2×10^6s, with two breaks of 6 and 18 days and with an average duty cycle of 430 s (Scholz et al. 2016). In 2016 September (after 287 days), the source FRB 121102 flashed four more times with an average time between bursts of $\simeq 10^3$ s (Marcote et al. 2017). Finally, after 340 days, 15 new bursts were recorded with an average duty cycle of 240 s in 2017 August (Gajjar et al. 2017).

We see that in all four series of bursts the time between successive flashes varies occasionally in a wide range from 10 s to $\simeq$ 1 h. However, the mean value of the duty cycle for all four series is $\tau_1 \simeq 500$ s. In addition, the duration of the series does not exceed a time of the order of $\tau_2 \simeq 20\,{\rm d} \simeq 2 \times 10^6$s, 16 days in the first series (2015 May) and 18 days in the second one (2015 November). The same time is a time of long breaks in the series. Finally, the time between series of bursts has an average value of $\tau_3 \simeq 200\,{\rm d} \simeq 6 \times 10^7$ s. These three times are very different, $\tau_1 << \tau_2 << \tau_3$.

Thus, the source FRB 121102 exhibits activity during a time $\tau_2 = 2 \times 10^6$s in the form of short bursts a few milliseconds in duration, $\tau \simeq 10^{-3}$s. The values of the duration of radio emission, intensity and duty cycle are close to the same values observed from so-called rotating radio transients (RRATs). However, although the duty cycle of RRATs is a random value, it is a multiplier of some constant time unit. This time unit is the period of rotation of a neutron star, which increases with time as for radio pulsars. The value of the time derivative of the period, $\dot{P}$, allows us to estimate the magnetic field strength at the stellar surface. It is $B \simeq 10^{12} - 10^{13}$G, which is similar to the value for radio pulsars. RRATs rotate more slowly than radio pulsars (the period of rotation is of the order of $\simeq$ 10 s). Because of this, there is no permanent generation of electron– positron plasma in the neutron star magnetosphere. Although the rotation of the magnetic field frozen into the star generates an electric field in the magnetosphere, $E \simeq v_{\rm rot}B/c$, the electric field is not sufficient for the continuous production of electron–positron plasma. However, a strong magnetic field and the electric field induced by the rotation of the magnetosphere of the neutron star can lead to a burst of production of electron–positron plasma under some external action. Such an external action in the case of RRATs is provided by Galactic and extragalactic gamma rays with energies above 1 MeV according to the mechanism for RRATs developed by Istomin & Sobyanin (2011a,b). Now a natural question arises: what will happen if a magnetized neutron star rotates even more slowly ($P > \tau_1 \simeq 500$ s) than the neutron star that is the source of RRAT? There is a strong magnetic field in its magnetosphere, but the electric field there is almost absent, which is necessary for acceleration of electrons and positrons in the magnetosphere and to begin a cascade plasma production process.

An electric field, $E \simeq uB/c$, can also arise in the magnetosphere when a sufficiently dense flow of charged particles moves through the magnetosphere with velocity u. Thus, we come to the conclusion that if the source of FRBs is a magnetized neutron star that rotates slowly enough and the birth of a relativistic plasma occurs in its magnetosphere, an external action is necessary to distort the magnetosphere and to create an electric field. The presence of a neutron star as a source of FRBs is indicated by the short duration of the burst of radio emission. Such an effect could be a flow of sufficiently dense plasma. Note that, from one flash to another, the dispersion measure is not constant but varies within 3 per cent during a time $\simeq \tau_1 = 500$ s. This would correspond to the motion of electron density inhomogeneities of size $<10^{13}$cm. This size is too small for either for the Galaxy or the previously mentioned intergalactic medium, but reflects the presence of plasma in the immediate vicinity of the neutron star. The characteristic values of the plasma flow scale L and its velocity u can be estimated from the following relations. First, $L/u = \tau_2 = 2 \times 10^6$ s. We assume that the time τ_2 is the time of close passage of a dense body (planet, comet, asteroid) near a neutron star. Secondly, we also assume that the repetition time of the burst series τ_3 is the orbiting period of a companion body around a neutron star, $P_{\rm orb} = \tau_3$. Since $\tau_2 << \tau_3$, the orbit of the body is strongly elongated. Therefore $u^2L = 2GM_{\rm NS} = 3.7 \times 10^{26}\ {\rm cm}^3\ {\rm s}^{-2}$. Here G is the gravitational constant and $M_{\rm NS}$ is the mass of the neutron star, which we put equal to $1.4\,{\rm M}_\odot$. Thus, we have

$$L = 10^{13}\ {\rm cm},\ u = 6 \times 10^6\ {\rm cm\,s}^{-1}. \qquad (1)$$

3 ELECTRIC FIELD

As a result, the scenario of interaction of a dense body orbiting around a magnetized hot neutron star (for FRB 121102) is as fol-

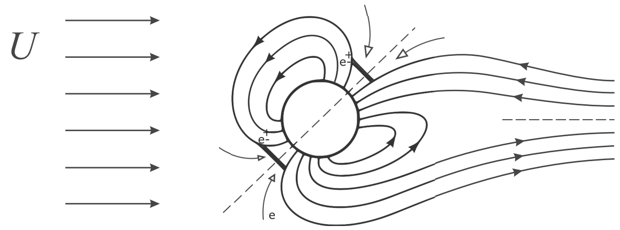

Figure 1. Scheme illustrating upflow of a magnetosphere of a magnetized neutron star by a plasma. Thin layers of electron–positron plasma are formed in the polar regions. The thermal radiation of the star is locked by these layers.

lows: a close pass of the body at a distance less than 1 au causes evaporation from the body and dense plasma flow around the magnetosphere of the neutron star, perturbing the magnetic field in the polar region and generating a longitudinal electric field (see Fig. 1).

For a temperature of the surface of the neutron star $T_{\rm NS}$ of the order of $10^8{\rm K} \simeq 10\,{\rm keV}$, the energy flux of X-ray photons is $S_{\rm NS} = 7 \times 10^{40}\,{\rm erg\,s^{-1}}$ (at a stellar radius of $R \simeq 10\,{\rm km}$), which exceeds the solar luminosity by 2×10^7 times. The evaporated plasma has a temperature of the order of the temperature of evaporation of a solid body, $T_{\rm b} \simeq 0.1\,{\rm eV}$, and its thermal velocity $v_{\rm T}$ is equal to $(2T_{\rm b}/Am_{\rm p})^{1/2} \simeq 10^5\,{\rm cm\,s^{-1}}$. Here we chose the mean atomic number of the evaporated ions to be of the order of $A = 20$ and $m_{\rm p}$ is the proton mass. Thus, $v_{\rm T} << u$ and the velocity of the flow around the neutron star magnetosphere is approximately equal to u. It should be noted that the temperature of the neutron star should not be too large, such that the radiation force acting on the ionized gas should not exceed the gravitation force, $T_{\rm NS} < (Am_{\rm p}cGM_{\rm NS}/ZR^2\sigma\sigma_{\rm T})^{1/4} = 0.3 \times 10^8 (A/Z)^{1/4} R_6^{-1/2}$ K. Here σ is the Stefan–Boltzmann constant and $\sigma_{\rm T}$ is the Thomson scattering cross-section. The pressure of the plasma flow destroys the magnetic field of the neutron star at distances from the centre of the star larger than a certain distance r determined by equality of pressures, $Am_{\rm p}n_{\rm e}u^2/2Z = B(r)^2/8\pi$. The value of Ze is the average ion charge, $A/Z \simeq 10$–20. We obtain $B(r) = (8\pi Am_{\rm p}n_{\rm e}u^2/Z)^{1/2}$. The region of perturbed magnetosphere at a distance r from the star, called cusp, has the same r. However, this decreases to a polar oval of small size at the stellar surface. Electron–positron plasma fills this polar region by a cascade process described below. The magnetic field at level r can be found by assuming that the magnetic field of the neutron star at large distances is dipole, $B = B_0(r/R)^{-3}$. Here B_0 is the value of the magnetic field on the surface of the NS, $B_0 = 10^{12}B_{12}$ G, $B_{12} = B_0/10^{12}$ G. Thus, $B(r) = 10^{12}B_{12}(r/R)^{-3}\,{\rm G} \simeq 10^3$ G for $r = 10^3R$. The electric field induced in the polar region, $E = uB/c$, is equal to $6 \times 10^{10}B_{12}(r/R)^{-3}\,{\rm V\,cm^{-1}} \simeq 60\,{\rm V\,cm^{-1}}$ for $r = 10^3R$. Accordingly, the rising voltage in the cusp is equal to $|\psi_0| = Er = 6 \times 10^{16}B_{12}R_6(r/R)^{-2}\,V \simeq 6 \times 10^{10}\,V$ for $r = 10^3R$ ($R_6 = R/10^6$cm). The electric field originates in the open region of the magnetosphere of the neutron star, in its polar region around the axis of stellar magnetic moment $\boldsymbol{M}$. On the surface of the star, this region is almost a circle of radius $R_0 \simeq R(R/r)^{1/2} \simeq 3 \times 10^4$ cm. This polar circle is the same as in neutron stars, which are radio pulsars. Knowing the magnitude of the magnetic field $B(r)$, we can determine the electron density $n_{\rm e}$ in the incoming plasma flow, $n_{\rm e} = n_0 = B^2(r)Z/8\pi Am_{\rm p}u^2 \simeq 10^{13}\,{\rm cm^{-3}}$ for $r = 10^3R$. A small fraction of electrons from the incoming stream penetrate into the polar region of the magnetosphere. Their flux S is equal to $S = \pi\eta n_0 ur^2 \simeq 2 \times 10^{39}\eta\,{\rm s^{-1}}$. Here η is the efficiency of penetration of electrons into the polar magnetosphere, $\eta << 1$.

The electric field arising in the magnetosphere at a distance $\simeq r$ penetrates deep into the magnetosphere in the polar region. Its dependence on height h above the surface of the star can be determined by solving the Laplace equation in the region bounded by the surface of radius $\rho(r) = h(h/r)^{1/2}$ with the following boundary conditions: $\psi(h = r) = \psi_0$, $\psi(h = R) = 0$, $\psi(h, \rho = h(h/r)^{1/2}) = 0$,

$$\frac{\partial^2\psi}{\partial h^2} + \frac{1}{\rho}\frac{\partial}{\partial\rho}\left(\rho\frac{\partial\psi}{\partial\rho}\right) = 0.$$

Replacing the transverse derivative $\Delta_\rho\psi$ by $-(4r/h^3)\psi$, we arrive at the equation

$$\frac{\partial^2\psi}{\partial h^2} = \frac{4r}{h^3}\psi.$$

The solution of this equation is $\psi = const \times (h/r)^{1/2}K_1[4(h/r)^{-1/2}]$, where $K_1(x)$ is the McDonald function of the first order. Since $h/r < 1$, $x > 4$, one can use an asymptotic presentation of the McDonald function for large arguments. As a result, we have

$$\psi(h) = \psi_0\left(\frac{h}{r}\right)^{3/4}\exp\left[-4\left(\frac{r}{h}\right)^{1/2} + 4\right]. \quad (2)$$

For practical purposes, the asymptotic presentation (2) does not differ from the exact solution expressed by the McDonald function and we will use expression (2) below. We see that the longitudinal electric field exists only in the upper part of the polar tube, $r > h > r^* \simeq 0.6r$. This field is exponentially suppressed when approaching the star. Thus, electrons falling into the polar region at $h \simeq r$ are accelerated toward the star, $\psi_0 < 0$ and obtain relativistic energy equal to $e|\psi_0|$, i.e. their Lorentz factor becomes equal to $\gamma_0 = 1.2 \times 10^{11}B_{12}R_6(r/R)^{-2} \simeq 10^5$.

4 PRODUCTION OF ELECTRON–POSITRON PLASMA

The process of production of electron–positron plasma in the polar magnetosphere of a magnetized hot neutron star is divided into several stages. This is primarily the production of gamma quanta as a result of inverse Compton scattering of thermal photons by relativistic primary electrons moving toward the surface of the star. Then the generated gamma quanta collide with thermal photons and give rise to $e^\pm$ pairs. Under the action of the radiation pressure of the star, pairs move outward and are captured into the bottom of a potential well formed on one side by radiation pressure and on the other by the blocking electrical potential ψ (2). We illustrate the birth process in Fig. 2.

Thermal photons with energies $E_{\rm ph} = m_{\rm e}c^2\epsilon$ propagating from the stellar surface are scattered by relativistic electrons and produce high-energy photons with energies $E'_{\rm ph} = m_{\rm e}c^2\epsilon'$:

$$\epsilon' = \epsilon\frac{4\gamma^2}{1 + \gamma^2\theta^2 + 4\gamma\epsilon}. \quad (3)$$

Here the angle θ is the angle of propagation of a scattered photon with respect to the velocity of a relativistic electron. Since, as can be seen from (3), ϵ' is large for photons scattered in the direction of propagation of the electron, $\theta \simeq \gamma^{-1} << 1$, we use the approximation $1 - \cos\theta \simeq \theta^2/2$. If we neglect the electron recoil, $4\gamma\epsilon < 1$, then the maximum energy of the scattered photon ($\theta = 0$) is $\epsilon'_{\rm m} = 4\gamma^2\epsilon$. For a Planck spectrum with temperature $T_{\rm NS}$, the maximum energy density is at the photon energy $E_{\rm ph} = 2.82T_{\rm NS} = 24T_8$keV. Here T_8 is the stellar surface temperature in units of 10^8 K, $T_8 = T_{\rm NS}/10^8$ K. Thus, the characteristic value of the thermal photon energy is $\epsilon \simeq 4.7 \times 10^{-2}T_8$. The value of $4\epsilon\gamma$ is equal to 2×10^4T_8. This means that,

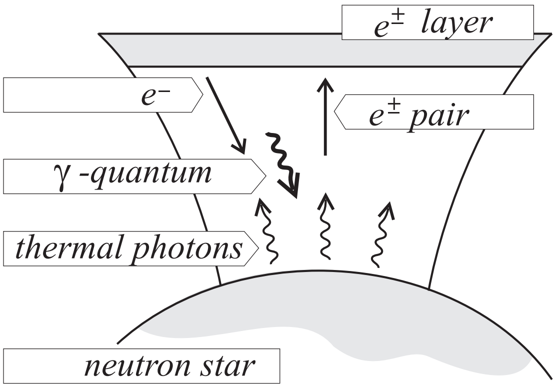

Figure 2. The birth of electron–positron plasma in the polar magnetosphere of a hot neutron star. A thermal photon is scattered by the primary electron e^- and generates a gamma quantum, which collides with thermal radiation and produces an $e^{\pm}$ pair. Electrons and positrons rise up under the action of the radiation pressure and accumulate in the narrow layer above the stellar surface.

for small scattering angles, $\theta < \theta^* = 2(\epsilon/\gamma)^{1/2} \simeq 1.4\times10^{-3}T_8^{1/2}$, the electron completely loses its energy, $\epsilon' = \gamma$. At larger angles, $\theta > \theta^*$, the energy of the scattered photon is $\epsilon' = 4\epsilon/\theta^2 < \gamma$. The scattered photon ϵ' propagates toward the star,where the density of thermal photons increases $\propto h^{-2}$ and the magnetic field also grows $\propto h^{-3}$. Therefore it is possible to create electron–positron pairs in two ways: (1) in collisions of scattered photons with thermal photons moving towards, $\epsilon' + \epsilon = \gamma^- + \gamma^+$, and (2) birth in a magnetic field, when the scattered photon intersects the line of strong magnetic field at the angle β, $\epsilon' \sin\beta > 2$. It should be noted that, in contrast to radio pulsars, energetic electrons γ and photons ϵ' propagate toward the star, where the birth of pairs is more efficient, but not outwards. In addition, curvature photons, which play a major role in initiation of cascade production of pairs in magnetospheres of radio pulsars, have small energies here and are incapable of producing pairs. Indeed, the energy of a curvature photon ϵ_c is $\epsilon_c = \bar{\lambda}\gamma^3/\rho_c \simeq 5\times10^{-4} << 1$. Here $\bar{\lambda}$ is the Compton wavelength of an electron, $\bar{\lambda} = \hbar/m_e c = 3.86\times10^{-11}$ cm, and ρ_c is the radius of curvature of magnetic field lines in the polar region, $\rho_c = 4(rh)^{1/2}/3 > 4(rR)^{1/2}/3 \simeq 10^8$ cm.

Let us consider the process of production of electron–positron pairs by scattered photons.

(1) For the production of a pair in collision consisting of a scattered photon with a thermal photon, $\epsilon'\epsilon > 1$, taking into account (3) we obtain the condition

$$\theta < [(2\epsilon - 1/\gamma)^2 - 2/\gamma^2]^{1/2} \simeq 2\epsilon,\ \epsilon > \epsilon_1 = (2^{1/2}+1)/2\gamma.$$

The inverse Compton scattering cross-section in the laboratory coordinate system associated with the neutron star has the form (Berestetskii, Lifshitz & Pitaevskii 1982)

$$d\sigma = \frac{8\pi r_e^2\gamma^2\theta d\theta}{(1+\gamma^2\theta^2+4\epsilon\gamma)^2}\left[\frac{1}{(1+\gamma^2\theta^2)^2} - \frac{1}{1+\gamma^2\theta^2} + \frac{1}{4}\left(\frac{1+\gamma^2\theta^2+4\gamma\epsilon}{1+\gamma^2\theta^2} + \frac{1+\gamma^2\theta^2}{1+\gamma^2\theta^2+4\gamma\epsilon}\right)\right]. \quad (4)$$

Here r_e is the classical radius of an electron, $r_e = 2.82\times10^{-13}$ cm. Integrating the cross-section (4) over an angle θ from 0 to $[(2\gamma\epsilon - 1)^2 - 2]^{1/2}/\gamma$ and then averaging over the Planck spectrum, we obtain the value of the thickness τ gained by an electron moving toward the neutron star when it is scattered by thermal photons:

$$\tau = \int_h^{r^*} dh' \int_{\epsilon_1}^{\infty} \frac{dn_{ph}}{d\epsilon} d\epsilon \int_0^{2\epsilon} \frac{d\sigma}{d\theta} d\theta = \frac{4a}{\pi}\frac{r_e^2 T^2 R}{\bar{\lambda}^3\gamma}\left(\frac{R}{h} - \frac{R}{r^*}\right). \quad (5)$$

The coefficient a is proportional to the integral of the scattering cross-section (4) over the Planck spectrum, starting from photon energies $\epsilon > (2 + {}^{1/2}1)/2\gamma$ to infinity. It depends on the electron energy γ and the temperature of the star, $T = T_{NS}/m_e c^2$. For the parameters of interest, $T_8 \simeq 1$, $\gamma \simeq 10^5$, the value of a is $\simeq 1$. The coefficient $4ar_e^2T^2R/\pi\bar{\lambda}^3\gamma$ on the right-hand side of (5) is large, $\simeq 5\times10^3 T_8^2\gamma_5^{-1}R_6$ ($\gamma_5 = \gamma_0/10^5$, $R_6 = R/10^6$ cm). This means that the thickness τ becomes of the order of unity fairly fast, $(r^* - h)/r^* \simeq 0.12$ and the electron completely loses energy, emitting a gamma quantum with energy $\epsilon' \simeq \gamma$.

Let us now find the efficiency of production of electron–positron pairs produced by collisions of energetic photons ϵ' with thermal photons ϵ. The cross-section for pair production is (Berestetskii et al. 1982)

$$\sigma^{\pm} = \frac{\pi r_e^2}{2}(1-v^2)\left[(3-v^4)\ln\frac{1+v}{1-v} - 2v(2-v^2)\right],$$

where v is equal to $v = (1 - 1/\epsilon\epsilon')^{1/2}$. Again averaging over the Planck spectrum, starting from the energy $\epsilon = 1/\epsilon'$ and assuming $\epsilon' \simeq \gamma$, we obtain the value of the thickness τ_e gained by a gamma quantum of energy ϵ' with respect to production of a pair,

$$\tau_e = \frac{b}{2\pi}\frac{r_e^2T^2R}{\bar{\lambda}^3\gamma}\left(\frac{R}{h} - \frac{R}{h_i}\right). \quad (6)$$

The numerical coefficient b for $T_8 = 1$, $\gamma = 10^5$ is $b \simeq 25$. Here h_i is the initial height from which the gamma photon begins to move toward the star. We see that the birth of an electron–positron pair is even more effective, $b/8a \simeq 3$ times, than gamma-ray radiation by an electron. Thus, the primary electron, passing a distance $\simeq qr^*$, $q \simeq 0.12(1 + 1/3) = 0.16$ through thermal radiation, produces a pair with electron and positron energies $\simeq \gamma/2$. In turn, this newborn pair produces a new pair in the photon field of the star, etc., so cascade production of electrons and positrons occurs. The minimum number of possible cascades K, if we do not take into account the decrease in energy of secondary electrons (γ in the expression 5), is determined by the condition

$$q\sum_{n=1}^{K}(1-q)^n = 1 - \frac{R}{r^*}.$$

Summing up, we obtain $K + 1 = [\ln(R/r^*)/\ln(1-q)]$. Substituting characteristic values $r^* \simeq 6\times10^2 R$, $q \simeq 0.16$, we find $K \simeq 35$. Since $2^{35} \simeq 3\times10^{10} >> \gamma_0$, the number of pairs λ generated by one primary electron is determined by the relation $\lambda \simeq \gamma/_0 2$.

(2) We now consider the single-photon production of electron–positron pairs by scattered photons ϵ' in the magnetic field of a neutron star. We need to know the number of photons generated with energy $\epsilon' > 2$. This is larger than the previously calculated density of energetic photons $\epsilon' > \epsilon^{-1} \simeq 10^2$, for photons that collide with thermal photons to produce pairs. Expression (3) for the energy of scattered photons determines the regions of angle θ and energy ϵ of primary photons:

$$\theta < (2\epsilon - 4\epsilon\gamma^{-1} - \gamma^{-2})^{1/2} \simeq (2\epsilon)^{1/2},\ \epsilon > 1/2\gamma(\gamma-2) \simeq 1/2\gamma^2.$$

As before, using the cross-section (4), integrating over the Planck spectrum, we obtain an expression for the thickness τ_1:

$$\tau_1 = \frac{4a_1}{\pi}\frac{r_e^2 T^2 R}{\lambdabar^3 \gamma}\left(\frac{R}{h} - \frac{R}{r^*}\right), \tag{7}$$

where the constant $a_1 = 3.8$ in the region of parameters $T_8 \simeq 1,\ \gamma \simeq 10^5$. Thus, the mean free path l_e of a fast electron with respect to the production of secondary photons with energy $\epsilon' > 2$ is 3.8 times smaller than the mean free path for production of energetic photons with energies $\epsilon' > 10^2,\ l_e \simeq 3 \times 10^{-2} r^*$.

Scattered secondary photons emitted by electrons along magnetic field lines and propagating toward the surface of the star begin to cross magnetic field lines at an angle $\beta \neq 0$ due to the curvature of magnetic field lines in the polar region of the magnetosphere. Thus, the angle β is equal to

$$\beta = \int_h^{h_i} \rho_c^{-1}(h')\,dh' = \frac{3}{2}\left[\left(\frac{h_i}{r}\right)^{1/2} - \left(\frac{h}{r}\right)^{1/2}\right].$$

Here the value of h_i is equal to the initial altitude from which the photon begins to propagate. As we will see, pair production occurs below the height $h \simeq 10R,\ h << r$. Because of this and because of the strong dependence of the magnetic field strength on the height, $B \propto h^{-3}$, angle β can be considered to be a constant, $\beta \simeq 3h_i^{1/2}/2r^{1/2}$. The probability of pair production in the magnetic field per unit time is (Berestetskii et al. 1982)

$$w = \frac{3^{3/2}\alpha c}{2^{9/2}\lambdabar} b|\sin\beta| \exp\left\{-\frac{8}{3\epsilon' b|\sin\beta|}\right\}\Theta(\epsilon'|\sin\beta| - 2). \tag{8}$$

The quantity b is the magnetic field intensity in units of the critical field, $b = B/B_\hbar$, $B_\hbar = 4.4 \times 10^{13}$G, α is the fine-structure constant, $\alpha = 1/137$, and $\Theta(x)$ is the stepwise theta function. The height, h, at which a pair is born with probability $\simeq 1$ is determined by the condition

$$\frac{1}{c}\int_h^{h_i} w(\epsilon', h')\,dh' = 1.$$

Setting $\sin\beta \simeq \beta$, we obtain

$$\left(\frac{h}{R}\right)^3 = \frac{9}{16}\epsilon' b_0 \left(\frac{h_i}{r}\right)^{1/2}\Lambda, \tag{9}$$

where the value of b_0 is $b_0 = B_0/B_\hbar = 2.3 \times 10^{-2} B_{12}$ and

$$\Lambda = \ln\left\{\frac{3^{1/6}}{2^{17/6}}\frac{R\alpha}{\lambdabar}(\epsilon')^{-2/3} b_0^{1/3}\left(\frac{h_i}{r}\right)^{1/6}\right\} - \frac{5}{3}\ln\left\{\frac{3^{1/6}}{2^{17/6}}\frac{R\alpha}{\lambdabar}(\epsilon')^{-2/3} b_0^{1/3}\left(\frac{h_i}{r}\right)^{1/6}\right\}.$$

The characteristic values of Λ are $\simeq$ 15–20. We see that photons with energy $\epsilon' \simeq 5$ produce electron–positron pairs only near the stellar surface, $h \simeq R$, while energetic photons, $\epsilon' \simeq 10^5$, can produce pairs at a distance $h/R \simeq 30$. After pairs are born, they lose their transverse momenta, emitting synchrotron photons with energy $\epsilon_s \simeq 3b\gamma^2|\sin\beta|/2 = 3b_0/2|\sin\beta| = 3b_0(r/h_i)^{1/2}/2 \simeq 1$. They cannot produce new pairs. Thus, the cascade single-photon production of electron–positron pairs in a strong magnetic field is actually absent in our case, in contrast to the cascade production of pairs in the magnetospheres of radio pulsars.

5 PLASMA TRAP

Thus, we see that primary electrons with a Lorentz factor $\gamma_0 \simeq 10^5$ effectively produce electron–positron pairs by interacting directly with thermal photons. In view of the cascade character of the process, most of them have multiplicity $\lambda \simeq 10^5/2$ and energies of $\gamma \simeq 1$ with an energy spread $\Delta\gamma \simeq 1$. However, they do not reach the surface, because the flux of thermal photons from the surface pushes them out. As a result, electrons and positrons are accelerated outward from the star:

$$\frac{d\gamma}{dt} = \sigma_T \frac{\pi^2 c T^4}{60\lambdabar^3}\left(\frac{R}{h}\right)^2.$$

Here σ_T is the Thomson cross-section, $\sigma_T = 8\pi r_e^2/3$, and we substitute $\sigma = \pi^2/60\hbar^3 c^2$ for the Stefan–Boltzmann constant in energy units. Integrating and setting $d\gamma/dt = c\,d\gamma/dh$, we obtain

$$\gamma = \frac{2\pi^3}{45}\frac{r_e^2 T^4 R}{\lambdabar^3}\left(1 - \frac{R}{h}\right).$$

In the region $R < h < r^*$, where, factually, there is no electric field (2), secondary electrons and positrons receive energy γ_f:

$$\gamma_f = \frac{2\pi^3}{45}\frac{r_e^2 T^4 R}{\lambdabar^3} = 1.6 \times 10^5 T_8^4 R_6.$$

Positrons escape outside freely, while electrons obeying the condition

$$\kappa = 0.6\gamma_5/T_8^4 R_6 > 1$$

are reflected from the electric potential (2) and become trapped. This trap is formed on one side by the electric potential created by the plasma flow in the magnetosphere and on the other by the thermal radiation of the star itself. Electrons stop at a distance $h = h_r$ from the star, defined by the condition $\gamma_f = e|\psi(h_r)|/m_e c^2$:

$$h_r = r\left[1 + \frac{1}{4}\ln\kappa + \frac{3}{8}\ln\left(1 + \frac{1}{4}\ln\kappa\right)\right]^{-2}.$$

After reflection, electrons begin to move toward the surface of the star under the action of the trapping potential $\psi(h)$ and, if there were no interaction with thermal photons moving toward them, they would arrive at the stellar surface with an energy corresponding to a Lorentz factor of γ_0/κ. However, they emit gamma quanta, which produce new electron–positron pairs, are decelerated to non-relativistic energies and are picked up again by the radiation of the star. The force acting on electrons moving outward from the star is potential, since the cross-section for scattering of thermal photons is equal to the Thomson cross-section, which does not depend on the electron energy. The corresponding potential ψ_T is equal to

$$\psi_T = m_e c^2 \gamma_f \frac{R}{h}.$$

The plot of the potential $\psi_T/e|\psi_0|$ ($\kappa = 3$) is shown in Fig. 3 The figure also shows the potential of the electric field $\psi(h)/\psi_0$ (equation 2) and the total potential $(e|\psi(h)| + \psi_T)/e|\psi_0|$.

As a result of pair production, electrons are accumulated at the bottom of the potential well formed by the total potential $e|\psi(h)| + \psi_T(h)$. The location of the bottom, h_m, is equal to (with logarithmic accuracy)

$$h_m = 7r\left[1 + \frac{1}{4}\ln\kappa + \frac{1}{4}\ln\left(\frac{2r}{R}\right) - \frac{5}{8}\ln\left(1 + \frac{1}{4}\ln\kappa + \frac{1}{4}\ln\left(\frac{2r}{R}\right)\right)\right]^{-2} < h_r. \tag{10}$$

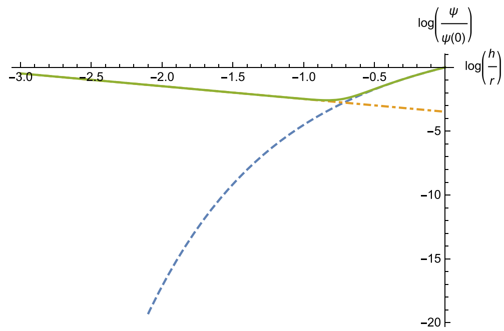

Figure 3. The logarithms of the potential $\psi_{\rm T}$ acting on electrons from the thermal stellar radiation (dot–dashed line), the potential $e|\psi(h)|$ (equation 2) (dotted line) and the total potential $e|\psi(h)| + \psi_{\rm T}$ (solid line) as functions of the height h.

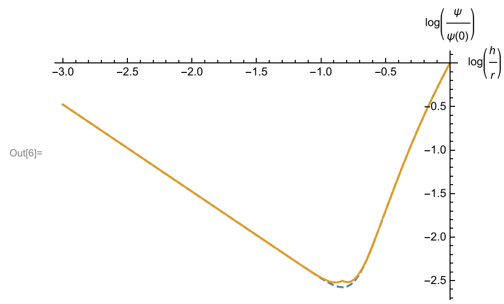

Figure 4. Logarithm of the total potential $e|\psi + \psi_{\rm sh}| + \psi_{\rm T}$ acting on electrons (solid line) and of the potential $e|\psi| + \psi_{\rm T}$, at which electrons are trapped initially (dotted line).

Thus, a primary electron with energy γ_0, after almost reaching the surface of the star and losing its energy, returns, receiving from thermal photons an energy less than the original one, γ_0/κ, then again moves toward the star and again produces quanta and pairs until it settles near the height $h_{\rm m}$. The multiplicity, i.e. the number of generated electrons per one primary electron, increases as $\lambda \simeq \gamma_0\kappa/2(\kappa - 1)$. It should be noted that the decelerating force acting on an electron when it moves toward the star is not conservative, since the scattering cross-section is inversely proportional to the electron energy, $\propto \gamma^{-1}$, as can be seen from expression (5). Therefore, the motion of energetic electrons toward the star is not the motion in the total potential $e|\psi(h)| + \psi_{\rm T}(h)$. Their deceleration occurs faster, ${\rm d}\gamma^2/{\rm d}h \propto h^{-2}$, but not ${\rm d}\gamma/{\rm d}h \propto h^{-2}$ as for motion in the potential $\psi_{\rm T}$. For subrelativistic electrons, $\Delta\gamma \simeq 1$, however, located near the bottom of the potential well, the cross-section for scattering of thermal photons is the Thomson one and their motion is potential in the potential $e|\psi| + \psi_{\rm T}$.

The thickness of the layer Δh, where electrons are accumulated, is small. It is determined by the spread of electron energies near the surface of the star after production of gamma quanta and pairs, $\Delta\gamma \simeq 1$:

$$\Delta h = h_{\rm m}\frac{\Delta\gamma}{2\gamma_{\rm f}}\left(\frac{h_{\rm m}}{r}\right)^{1/2} << h_{\rm m}.$$

This layer creates an additional electric potential $\psi_{\rm sh}$. The equation for the potential $\psi_{\rm sh}$ is as follows :

$$\frac{{\rm d}^2\psi_{\rm sh}}{{\rm d}h^2} - \frac{4r}{h^3}\psi_{\rm sh} = en_l(h)\frac{4r}{h^3}, \qquad (11)$$

where $n_l(h)$ is the electron column density associated with the density n by the relation $n_l = \pi nh^3/r$ due to the dependence of the cross-section of the magnetic tube s on height, $s = \pi h^3/r$. Since the thickness of the electron layer is initially small, we can assume $n_l = N_{\rm e}\delta(h - h_{\rm m})$, where $\delta(x)$ is the Dirac delta function. The solution of equation (11) is

$$\psi_{\rm sh} = -\frac{eN_{\rm e}r^2}{h_{\rm m}^3}\left(\frac{hh_{\rm m}}{r^2}\right)^{3/4}\begin{cases}\exp\left(-\frac{4r^{1/2}}{h^{1/2}} + \frac{4r^{1/2}}{h_{\rm m}^{1/2}}\right), & h < h_{\rm m},\\ \exp\left(-\frac{4r^{1/2}}{h_{\rm m}^{1/2}} + \frac{4r^{1/2}}{h^{1/2}}\right), & h > h_{\rm m}.\end{cases}$$

We see that the electric field of the electron layer decreases exponentially on both sides. The characteristic width of the additional potential, δh, does not depend on the value of the potential and is equal to $\delta h \simeq h_{\rm m}(h_{\rm m}/r)^{1/2}$, i.e. it is equal to the width of the magnetic tube at the altitude $h = h_{\rm m}$. The total potential $(e|\psi + \psi_{\rm sh}| + \psi_{\rm T})/e|\psi_0|$ is shown in Fig. 4 taking account of the additional electric field created by electrons trapped in the minimum of the potential.

At height $h \simeq h_{\rm m}$, electrons form a local potential well, to which positrons are attracted. Fig. 5 shows the potential acting on positrons, $[e(\psi + \psi_{\rm sh}) + \psi_{\rm T}]/e|\psi_0|$. In the absence of the potential of the electron layer $\psi_{\rm sh}$, positrons are accelerated by thermal radiation up to an energy γ_0/κ, then receive even more energy from the potential ψ. However, if we do not take into account the interaction of positrons passing through the layer with electrons, then positrons are not trapped inside the layer. Deceleration of positrons due to bremsstrahlung radiation in the potential $\psi_{\rm sh}$ is not enough for them to lose energy γ_0/κ. Here we should take into account the collective interaction of passing positrons with Lorentz factor $\gamma^+ = \gamma_0/\kappa$ with electrons and positrons trapped earlier. The plasma density in the layer is equal to $n_{\rm p} = n^- + n^+$, so their plasma frequency is $\omega_{\rm p}^2 = 4\pi n_{\rm p}e^2/m_{\rm e}$. A beam of positrons with density $n_{\rm b} = S\lambda r/\pi ch_{\rm m}^3$ excites plasma oscillations in the layer, which decelerate positrons due to beam instability. The increment of the beam instability, $\nu_{\rm b}$, is

$$\nu_{\rm b} = \frac{3^{1/2}}{2^{4/3}}\left(\frac{n_{\rm b}}{n_{\rm p}\gamma^{+3}}\right)^{1/3}\omega_{\rm p},\ n_{\rm p} > n_{\rm b}/\gamma^{+3}.$$

Positrons are trapped into the well $\psi_{\rm sh}$ if the condition $\nu_{\rm b}\tau_{\rm p} > 1$ is satisfied, where $\tau_{\rm p}$ is the time of passage of positrons though the plasma layer, $\tau_{\rm p} = (h_{\rm m}^3/r)^{1/2}/c$. Let us show for characteristic values of quantities that the condition $\nu_{\rm b}\tau_{\rm p} > 1$ is already satisfied at low densities of plasma in the layer. We take characteristic values: $h_{\rm m} \simeq 10^8$ cm, $r \simeq 10^9$ cm, $s_{\rm m} = \pi h_{\rm m}^3/r \simeq 3\times10^{15}\,{\rm cm}^2$, $\tau_{\rm p} = (h_{\rm m}^3/r)^{1/2}/c \simeq 10^{-3}$ s, $\lambda \simeq 10^5$, $\gamma^+ \simeq 3\times10^4$, $\omega_{\rm p} = 5.6\times10^4 n_{\rm p}^{1/2}$, $n_{\rm b} = S\lambda/s_{\rm m}c \simeq 2\times10^{18}\eta\,{\rm cm}^{-3}$. As a result we obtain the condition

$$n_{\rm p} > 2\times10^{16}\left(\frac{n_{\rm b}}{1\,{\rm cm}^{-3}}\right)^{-2}\,{\rm cm}^{-3},$$

which is obviously valid for $n_{\rm p} > n_{\rm b}/\gamma^{+3} \simeq 10^5\eta\,{\rm cm}^{-3}$. Thus, we see that a thin layer of width $(h_{\rm m}^3/r)^{1/2}$ arises in the polar region, where all electrons and positrons produced in the polar magnetosphere are accumulated. The plasma is an almost neutral, subrelativistic electron–positron plasma in this layer. Electrons are confined there by the locking potential ψ (equation 2), while positrons are held by confined electrons.

There is a slight displacement of the position of the plasma layer from the height $h = h_{\rm m}$. Indeed, the drag force of the thermal radiation of the star becomes twice as large as the drag force acting separately on electrons and positrons. In the expression for the value of $h_{\rm m}$ (10), it is necessary to replace κ by $\kappa/2$. As a result, the shift

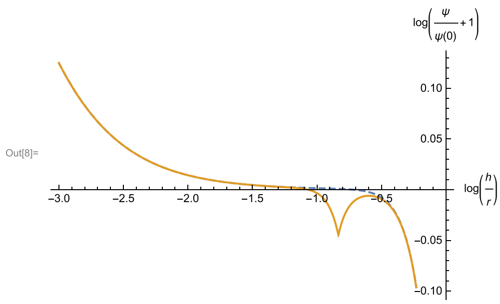

Figure 5. Logarithm of the potential $e(\psi + \psi_{\rm sh}) + \psi_{\rm T}$ acting on positrons (solid line) and of the potential $e\psi + \psi_{\rm T}$ without influence of trapped electrons (dotted line).

of the position of the plasma layer $\Delta h_{\rm m}$ is equal to

$$\frac{\Delta h_{\rm m}}{h_{\rm m}} = \frac{\ln 2}{2}\left(\frac{h_{\rm m}}{r}\right)^{1/2}\left[1 - \frac{5}{8}\left(\frac{h_{\rm m}}{r}\right)^{1/2}\right] \simeq 0.1.$$

6 BURST

Electron–positron plasma constantly accumulates in the layer. The number of pairs N grows linearly with time, $N = S\lambda t \simeq 2 \times 10^{44}\eta t$. Accordingly, the plasma column density $\int n_{\rm p}\,{\rm d}h = N/s_{\rm m} = Nr/\pi h_{\rm m}^3 \simeq 6 \times 10^{28}\eta t$ cm^{-2} grows. The thickness relative to scattering of thermal radiation of the star by the electron–positron plasma of the layer, $2\sigma_{\rm T}\int n_{\rm p}\,{\rm d}h$, also grows with time and becomes equal to unity at time $t^* = s_{\rm m}/2S\lambda\sigma_{\rm T} \simeq 10^{-5}\eta^{-1}$ s. This means that the light of the star will start to reflect effectively from the plasma layer. If the reflection coefficient from the layer and the surface of the star is η_r, then the energy density of thermal radiation in the polar region under the plasma layer will begin to grow exponentially, with characteristic time $2h_{\rm m}/c\eta_r \simeq 10^{-2}\eta_r^{-l}$ s. The effective temperature of trapped radiation $T_{\rm eff}$ will grow in time until radiation pushes the plasma layer outside the magnetosphere. This happens when the radiative force acting on an electron and a positron of the plasma layer, $2\sigma\sigma_{\rm T}T_{\rm eff}^4(R/r)^2/c$, at height $h = r$ becomes equal to the force acting on electrons from the locking potential at the same height, $e{\rm d}|\psi|/{\rm d}h|_{h=r} = 11e|\psi_0|/4r$. As a result, we have

$$T_{\rm eff} = 10^8\,{\rm K}\left(\frac{11r}{8R}\kappa\right)^{1/4} \simeq 10^9\,{\rm K}.$$

The released energy can be estimated as the energy of thermal radiation of the neutron star accumulated during time τ_1, $\mathcal{E} > \pi\sigma T_{\rm NS}^4\tau_1 R^3/r \simeq 10^{40}T_8^4$ erg. Here we chose the minimum area of trapped radiation $s_R = \pi R^3/r \simeq 3 \times 10^9$ cm^2, which corresponds to the area of the polar region on the stellar surface. At height $h = h_{\rm m}$, the area of trapped radiation is much larger, $s_{\rm m} \simeq 3 \times 10^{15}$ cm^2. In view of the rapid growth of the energy of trapped stellar radiation, ejection of the plasma layer will occur during a short time period $(r - h_{\rm m})/c \simeq 3 \times 10^{-2}$ s. Therefore, the power W of energy release is equal to $W = c\mathcal{E}/r > 3 \times 10^{41}T_8^4$ erg s^{-1}.

All the energy released is energy from the stream of relativistic electrons and positrons emitted from the polar magnetosphere. The Lorentz factor of the e$^{\pm}$ plasma is the same as that of the secondary pulsar plasma, $\gamma \simeq 10^2$. However, the densities here are much larger: 10^{17} cm^{-3} at an altitude $r \simeq 10^9$cm, while radio pulsars have the same density, 10^{17} cm^{-3}, only near the stellar surface. Radio emission arises as the collective radiation of a stream of electrons and positrons accelerating along magnetic field lines. Acceleration can be across the magnetic field associated with the curvature of the magnetic field in the polar magnetosphere. Here, the radius of curvature $\rho_{\rm c} \simeq 10^8$–10^9cm is the same as in the polar magnetosphere of radio pulsars. Also, during a burst, electrons and positrons of a layer are accelerated in the longitudinal direction with the same characteristic time $(r - r_{\rm m})/c \simeq \rho_{\rm c}/c$. Thus, the radio emission mechanism in this case is of the same nature as that of radio pulsars. For the coefficient of transformation of the energy flux of the relativistic plasma into radio flux for pulsars, $\alpha_r \simeq 10^{-4}$ (Beskin, Gurevich & Istomin 1993), the power emitted in the radio range is $W_r \simeq 3 \times 10^{37}T_8^4$ erg s^{-1}. This corresponds to an energy density of the electromagnetic field $E^2/8\pi = W_r/r^2c \simeq 10^9$ erg cm^{-3}. Such an energy density, which is nine orders of magnitude higher than the energy density of the plasma incident from the companion ($n_{\rm e}T_{\rm p} \simeq 10^{13}$ cm $^3\times$ 0.1 eV = 1 erg cm^{-3}), forms a vacuum 'hole' in the streaming plasma through radio-wave ponderomotive action. Thus, the powerful radio emission together with the flow of relativistic electron–positron plasma will leave the interaction area of the magnetized neutron star with the companion body unhindered and exit into the surrounding interstellar matter without causing any additional radiation.

The energy density of a powerful radio wave becomes comparable with the energy density of the incoming plasma only at large distances from the neutron star, $l \simeq r[(E^2/8\pi)/n_{\rm e}T_{\rm p}]^{1/2} \simeq 3 \times 10^4 r = 3 \times 10^7 R$. From these distances, the radio wave begins to propagate in a plasma and to undergo dispersion. The characteristic value of the electron density at such distances can be estimated from the observed variations of the dispersion measure $\delta(DM)$ of 3 per cent: $\delta(DM) \simeq 17$ cm^{-3}pc $= 5 \times 10^{19}$ cm^{-2}. Thus, $n_{\rm e} \simeq 10^6$ cm^{-3}, which is much less than the density of the incoming plasma near the neutron star at a distance $r \simeq 10^9$ cm, $n_{\rm e} \simeq 10^{13}$ cm^{-3}. We see that the near region of interaction of a dense wind from a companion with a neutron star does not contribute to the observed dispersion measure, because of the high power of the radiating radio emission. The contribution is given by the far regions. It is also possible to estimate the value of the characteristic magnetic field at distances $\simeq l$ from a NS. Since $l \simeq 3 \times 10^7 R$, the dipole field of a neutron star is practically zero at such distances. However, the plasma flow destroys and captures the magnetic field of the star at distances larger than r, where the magnetic field is $B < 10^3$G. Because of the magnetic field being frozen into the conducting plasma, the magnetic field at large distances $\simeq l$, where $n_{\rm e} \simeq 10^6$ cm^{-3}, is $B < 10^3{\rm G}(10^6/10^{13}) = 10^{-4}$G. Such a field has a random orientation. Therefore, the possible observable rotation measure RM due to Faraday rotation on the periphery of the interaction does not exceed a value of $RM < 10^3$ rad m^{-2} in this case. Of course, this estimate corresponds to the characteristic values chosen here. In a particular burst with other parameters, other values of the rotation measure can be observed. For example, the observed RM values lie in the range 10–1000 rad m^{-2} for one-off FRBs.

For positrons with total number of the order of $10^{44}T_8^4$, scattering in interstellar matter, cooling down and annihilating with the electrons of this matter will yield gamma quanta with an energy 511 keV. If we take characteristic values of the temperature and density of interstellar matter of $T \simeq 0.1$ eV and $n \simeq 0.1$ cm^{-3}, respectively, then the luminosity in the 511-keV gamma line will be $\simeq 10^{26}$ s^{-1} during the time after the burst. Such a flux cannot be seen from a distance of $\simeq 1$ Gpc.

The energy radiated in the radio range, $\mathcal{E}_r$, is equal to $\mathcal{E}_r \simeq 10^{36}T_8^4$ erg. It is considerably smaller than the estimate given at the beginning of the article, $\mathcal{E}_r \simeq 10^{40}$ erg. The latter value was obtained under the assumption that the radiation is isotropic and has no directivity. It is clear, however, that electrons and positrons that receive energy from the flux of photons at an altitude r from the surface of the star have directional motion spread over angles $\Delta\theta = R/r \simeq 10^{-3}$. Besides this, the radio emission will have directivity $\Delta\theta \simeq \gamma^{-1}$, where γ is the Lorentz factor of accelerated particles. The force $\sigma\sigma_{\rm T}T^4(R/r)^2/c \simeq 1.3 \times 10^{-13}T_8^4$ din will lead to acceleration inside the region $r \simeq 10^9$cm up to an energy $\gamma \simeq 2 \times 10^2 T_8^4$. This value will determine the directivity of the radio emission, $\Delta\theta \simeq 5 \times 10^{-3}$. It should be noted that this value of directivity follows from observations of the source FRB 121102. The observed breaks in the series of bursts have approximately the same duration as the series themselves. If one interprets breaks as departures of the polar region from the line of sight of the observer due to rotation of the star with period P, then $P \simeq \tau_2 \simeq 20$ d. The radiation directivity of $\simeq 10\tau_1/\tau_2 \simeq 2.5 \times 10^{-3}$ agrees well with the value of γ, the energy of accelerated electrons and positrons produced in the polar cap of the magnetosphere of the star. Taking into account such directivity, the effective energy radiated in the radio range will increase $4/\Delta\theta^2 \simeq 6 \times 10^5$ times, when recalculated to the total solid angle 4π, which gives $\mathcal{E}_r^{\rm eff} \simeq 6 \times 10^{41}$ erg. This is sufficient to place a source of FRB at cosmological distances. As for the duration of the burst, τ, it is defined by the size of the region of acceleration of the plasma layer, its thickness $\simeq r \simeq 10^9$cm and the Lorentz factor of accelerated electrons and positrons, $\tau \simeq (r/c)/\bar{\gamma}$. For $\gamma = 2 \times 10^2$, the value of $1/\bar{\gamma} = \ln(\gamma)/\gamma$ is $\simeq 3 \times 10^{-2}$. Thus, the time $\tau \simeq 10^{-3}$ s is in agreement with the time observed in FRBs.

7 DISCUSSION

The short burst time $\tau \simeq 10^{-3}$s tells us about the compactness of the source of radiation: apparently a neutron star. Neutron stars observed as radio pulsars have strong magnetic fields $B_0 \simeq 10^{12}$ G and small rotation periods $P = 10^{-3}$–1 s. An electric field, which has non-zero projection on to the magnetic field, arises in the magnetosphere due to the rotation of a magnetized body. This occurs in the open magnetosphere, in which magnetic field lines reach the light cylinder surface $c/\Omega_{\rm NS}$. Near the stellar surface the polar cap is formed, in which continuous generation of electron–positron plasma and an electric current takes place. For a fixed value of the magnetic field B_0, continuous plasma generation is possible for sufficiently fast rotation of the star, $P < (B_0/10^{12}\,{\rm G})^{8/15}$ s. The radio emission from a pulsar is due to the flow of relativistic plasma in the open magnetosphere (Beskin et al. 1993). The energy of radio emission is a small fraction of the energy of the plasma flow, which in turn is fed from the energy of rotation of the star. Continuous plasma generation becomes impossible for slower rotation. However, plasma generation can occur in flares, as in the case of RRATs. The electric field in the polar magnetosphere begins to accelerate electron–positron pairs produced by an accidental energetic gamma quantum. This leads to cascade plasma production and to a flash of radio emission. For even slower stellar rotation, the electric field, $E \simeq \Omega_{\rm NS}RB/c$, is not sufficient for cascade plasma production in the magnetosphere. However, the external action of plasma flow with velocity u relative to the neutron star produces a sufficiently strong electric field. It generates cascade plasma production in the polar magnetosphere. The process of plasma birth in this case has special characteristics in comparison with that taking place in the magnetosphere of a radio pulsar. Acceleration of electrons, injected from the plasma stream into the magnetosphere, occurs toward the star (aurora). In this case, the production of gamma rays by inverse Compton scattering of thermal photons is most efficient. The multiplication factor λ (the number of pairs per one primary electron) reaches a value of 10^5. The electron–positron plasma created is accumulated in the polar magnetosphere in the form of a thin layer in the region of a minimum of the total potential created by the electric field, which is generated by flowing plasma and the radiation pressure of the thermal radiation of the star (see Fig.4). The electron–positron plasma density in the layer increases linearly with time and grows so high that the layer becomes a mirror for thermal radiation. After that, the effective temperature of radiation of the star under the mirror begins to grow exponentially up to a value of $T_{\rm eff} \simeq 10^9$K. Finally, the stellar radiation pushes the electron–positron layer outward, overcoming the external blocking electric potential for a short time $\simeq 3 \times 10^{-2}$s. A 'hole' in the magnetosphere of a neutron star and in the stream of plasma flowing around the star forms. Accelerated up to energies $\gamma \simeq 2 \times 10^2$, the flow of electrons and positrons generates radio emission by a mechanism analogous to the mechanism of radio emission from pulsars. It is worth noting that the Lorentz factor of accelerated plasma is the same as in the secondary electron–positron plasma in the magnetosphere of a pulsar. Because the flow is relativistic, the radio emission is directed into a small solid angle, 6×10^{-5} times smaller than 4π. This allows for the total energy of the burst radiated in the radio range, $\mathcal{E}_r \simeq 10^{36}$ erg, to appear as effective energy of an isotropized burst more than five orders of magnitude higher, $\mathcal{E}_r^{\rm eff} \simeq 10^{41}$ erg. The duration of radio emission is small, $\tau \simeq 10^{-3}$s, also because of the relativistic nature of the plasma flow.

From the model discussed, it follows that the region of the polar magnetosphere $h \simeq h_{\rm m}$ is a source of gamma radiation in the range of the annihilation line at 511 keV just before the burst of radio emission during time τ_1, when there is an accumulation of electron–positron plasma in the layer. The plasma density in the layer can be estimated from the condition $2n_{\rm p}h_{\rm m}\sigma_{\rm T} \simeq 1$, when the layer becomes the mirror, $n_{\rm p} \simeq 10^{16}\,{\rm cm}^{-3}$. This value will also be the density of gamma quanta $n_\gamma = n_{\rm p}$. Correspondingly, the almost isotropic flux of gamma radiation from the layer in the annihilation line is $S_\gamma = n_\gamma c s_{\rm m} \simeq 10^{42}\,{\rm s}^{-1}$ and the luminosity in the annihilation line is $\simeq 10^{36}\,{\rm erg\,s}^{-1}$. The observation of the 511-keV annihilation line from a FRB would confirm the presence of an electron–positron plasma in the source of the FRB and would allow us to determine the distance to the source independently from the line shift.

Although our results were based on observations of the repeating source FRB 121102, we think that the proposed mechanism for the origin of FRBs is valid for both the repeating source 121102 and single FRBs. Apparently, the repetition is due to the fact that the neutron star in this case is an almost coaxial rotator ($P \simeq 20$ d). In this case, the direction of the axis of the magnetic dipole coincides with the axis of rotation. Each close passage of the companion body leads to the observation of a burst. If the neutron star has a finite angle of inclination of its axes, then for $P > \tau_2$ only one flare event is observed. When the radio emission is highly directed, $\Delta\theta \simeq 5 \times 10^{-3}$, the ratio of the number of repeated sources of FRBs to the number of single sources will be $\simeq \Delta\theta$, i.e. for 200 single bursts one repeating source has to be observed. In addition, with strong evaporation of a body, large loss of its mass occurs and if a body is not massive enough it can receive sufficient additional moment that will lead it away from orbit around the neutron star.

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