Interaction of a Supernova Shock with the other Star in a Binary System

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Abstract

Interaction of a fast shock wave generated during a supernova explosion with a magnetized star-companion of the supernova precursor produces a current sheet. We consider the evolution of this current sheet and show that a singularity (shock) is formed in finite time within the ideal MHD framework. Charged particles (electrons) are accelerated in the vicinity of the singularity, and their distribution function has a plateau up to the energies of the order of $10^4 m c^2$. These fast particles radiate in the $\gamma$-range in the strong magnetic field of the current sheet ($B \sim 10^6 G$). Radiation is concentrated within a narrow angle around the current sheet, $\Delta \theta \sim 3 \cdot 10^{-4}$, and its spectrum has the maximum at several hundreds of keV. Presented calculations confirm the model of cosmological GRBs proposed by Istomin & Komberg (2002).

1 Introduction

Observations show that some gamma-ray bursts (GRBs) occur at cosmological distances, and some of the bursts are related to the supernova explosions. If the radiation of a cosmological GRB is isotropic, its energy release is estimated to be $10^{43} - 10^{44}$ erg at the hundreds of keV range. Massive supernovae emit comparable amounts of energy, mainly in the form of neutrinos. Thus a question about the origin of GRBs’ energy supply arises. Connection between the GRB phenomenon and that of a supernova seems to be supported by several facts. First, the paper by [3] reports the temporal and spatial concurrence of the “compact” supernova SN 1998 bw ($10^{45}$ erg, type SNIc) with GRB 980425, identified with a nearby galaxy which has $Z = 0.0085$. Secondly, “hills” are observed in the light curves of some optical transients several days after the GRB; they can be interpreted as a contribution from the supernova burst simultaneous with the GRB. Thirdly, narrow X-ray lines of ions characteristic

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of supernova shells are detected in the spectrum of the afterglow [9], with the velocity of the line-emitting material reaching $0.1c$. Obviously, the “energy crisis” arises if we attempt to relate an isotropic GRB to a supernova. One way to avoid this crisis is to suggest that GRBs are associated with strongly collimated relativistic ejections radiating within $\gamma$-range. This makes it possible to decrease by many orders the luminosity from cosmological GRBs corresponding to the observed fluxes. [1] proposed this connection between GRBs and strongly anisotropic supernova bursts. Summarizing, we see that current understanding of the nature of GRBs requires a combination of two conditions: a supernova explosion (possibly of special type) and the formation of a narrow (opening angle less than a degree) beam of relativistic particles radiating within X- to $\gamma$-range. [3] proposed a model where a beam of relativistic particles is formed by the interaction of a fast shock wave generated during a supernova explosion with a magnetized star-companion (a neutron star or a white dwarf) of the supernova precursor. The flow of the shock matter around the magnetosphere of the star-companion forms an almost parallel magnetospheric tail. Magnetic field in the tail becomes super strong (about $10^6$ G) due to the cumulative effect. Virtually any star-companion interacting with a supernova shock will produce such strong magnetic field, due to amplification during the explosive compression of conducting medium - the mechanism proposed by [3] to produce very strong magnetic fields. For the typical shock parameters, the density $\rho_{sh}$ is of order $10^{-8} g/cm^3$, and the velocity $V_{sh}$ is about $4 \cdot 10^9 cm/sec$, the value of the amplified magnetic field is $B = (4\pi \rho_{sh} V_{sh}^2)^{1/2} \approx 10^6$ Gauss. For a usual star like the sun (which has magnetic field of the order of 10G), the compression will be about $3 \cdot 10^2$ times and the transversal size of the tail will be equal to $R_\odot/3 \cdot 10^2 \approx 10^8 cm$. Hence, virtually any explosion of a supernova which had a star-companion, will result in a highly magnetized elongated plasma formation in the shadow of the star-companion. The magnetic field changes its direction from one edge of the tail to the other, which leads to the magnetic reconnection, magnetic energy release, and acceleration of charged particles. Similar processes take place in the magnetospheric tail of the Earth formed by the flow of the solar wind. In a strong magnetic field these relativistically accelerated particles produce synchrotron radiation concentrated in a narrow cone $\Delta \theta \approx 3 \cdot 10^{-4}$. The frequency of supernova explosions has the order of $3 \cdot 10^{-2} supernovae/yr$ per a galaxy, which gives $(3 \cdot 10^{-2})^2 \cdot 3 \cdot 10^{-2} \cdot 10^{11} \approx 300$ radiation bursts per year, directed towards the Earth (the Universe has $\approx 10^{11}$ galaxies). All this may provide an insight into the nature of GRBs. In this paper we provide an accurate quantitative theory of the outlined phenomena, namely the evolution of the current sheet, the acceleration of charged particles, and the formation of the radiation spectrum.

2 Stationary current sheet

Magnetic field of a star is almost dipolar at the large distances from the center of the star. Its direction changes from one side of the sheet to the other (with respect to the axis of the tail). Interaction of a dense plasma flow with this dipole magnetic field produces a long current sheet. A similar magnetospheric tail is formed by solar wind in the dipolar magnetic field of the Earth. Such situation may also arise in a close binary system of stars, consisting of a magnetized star.
and a massive supernova precursor. A shock wave from the exploding massive star interacts with the magnetized star and forms a current sheet, resulting in a GRB. Such model of GRBs was proposed in the papers of [5]. Yet another example of the current sheet formation comes from a binary system of neutron stars, one of which is a radio pulsar. Relativistic wind flow of the radio pulsar around the other star forms a magnetic tail. Discovery of a close binary system of neutron stars containing two radio pulsars PSR J0737-3039 A,B was reported in the paper of [11]. Observations of this system show that the wind flow of the more powerful pulsar 'A' around the magnetosphere of the radio pulsar 'B' takes place inside the light cylinder of the pulsar 'B'. Thus, a long and almost parallel tail of magnetic field is formed in the magnetosphere of 'B'.

To study the phenomena taking place in the magnetosphere of a magnetized star, we assume the sheet to be flat and infinitely long. Therefore, all characteristics depend only on one coordinate \( z \), which is perpendicular to the current sheet. Our approximation is one-dimensional, which makes possible the solution of non-stationary problems. The geometric structure of this problem is presented in Figure 1.

Non-uniform magnetic field is parallel to the \( x \) axis and the tail, the \( y \) axis is perpendicular to the \( x-z \) plane. Magnetic field depends linearly on the \( z \) coordinate

\[
\mathbf{B} = (Kz, 0, 0),
\]

where \( K \) is a constant determined by electric currents flowing in the \( y \) direction. A particle with charge \( q \) and mass \( m \) in such non-uniform magnetic field is subject to the gradient drift [4]

\[
u_y = -\frac{cp_|v_|}{2mqKz^2},
\]

(1)

Here \( v_| \) and \( p_| \) are the velocity and momentum components of the particle perpendicular to the magnetic field. Equation (1) is not applicable near the plane \( z \simeq 0 \). The Larmor radius of the particle is greater than the inverse gradient of the field there, so the drift approximation is not valid when \( Kz^2 < cp_|/q \). Electric current \( j_d \) produced by the drift motion of charged particles with density \( n \) is given by

\[
j_d = -n\frac{<p_|v_|>}{2m}c\frac{1}{Kz^2},
\]

(2)
and is independent of their charge. Brackets \(< \ldots >\) mean averaging over the distribution function of the particles. In addition to the drift current resulting from the drift of the particles in the non-uniform magnetic field, there also exists a magnetization current \(j_m\) due to non-uniform distribution of the magnetic moments of the particles, \(j_m = c \cdot \text{curl} \mathbf{M}, \mathbf{M} = n \langle \mu \rangle, \mu = -\mathbf{B}(p_{\perp}v_{\perp}/2B^2)\). The sum of the drift current and the magnetization current depends only on the gradient of the transversal energy density of particles
\[
j_y = -\frac{c}{B} \frac{d}{dz} \left( n \langle p_{\perp}v_{\perp} \rangle \right).
\] (3)

The total current \(j\) is formed by the electron and ion flows. Assuming plasma to be quasineutral and consisting of electrons and positively charged ions with mean charge Ze, we see that the electric current is
\[
j_y = -\frac{c}{B} \frac{d}{dz} \left[ n_i(1 + Z) \langle p_{\perp}v_{\perp} \rangle \right].
\] (4)

On the other hand, we have
\[
j = \frac{c}{4\pi} \text{curl} \mathbf{B} = \frac{cK}{4\pi} e_y.
\]

Finally,
\[
\frac{B^2}{8\pi} + n_i(1 + Z) \langle p_{\perp}v_{\perp} \rangle = \text{const} = P_0.
\] (5)

Equation (5) is virtually identical to the equilibrium condition of an inhomogeneous plasma in a non-uniform magnetic field, when the equilibrium is supported by external pressure \(P_0\).

3 Evolution of the current sheet

Now we add an extra assumption that the current sheet is compressed by some external force acting along the \(z\) axis, and consider the evolution of such sheet. We may assume that the evolution is fast enough and it can be described by ideal magnetohydrodynamics (MHD). MHD equations for the one-dimensional case (all parameters depend only on the \(z\) coordinate) are
\[
\frac{\partial B}{\partial t} + v_z \frac{\partial B}{\partial z} + B \frac{\partial v_z}{\partial z} = 0,
\]
\[
\frac{\partial v_y}{\partial t} + v_z \frac{\partial v_y}{\partial z} = 0,
\]
\[
\rho \left( \frac{\partial v_z}{\partial t} + v_z \frac{\partial v_z}{\partial z} \right) + \frac{\partial P}{\partial z} + \frac{1}{4\pi} B \frac{\partial B}{\partial z} = 0,
\]
\[
\frac{\partial \rho}{\partial t} + v_z \frac{\partial \rho}{\partial z} + \rho \frac{\partial v_z}{\partial z} = 0,
\]
\[
\frac{\partial P}{\partial t} + v_z \frac{\partial P}{\partial z} + \Gamma P \frac{\partial v_z}{\partial z} = 0.
\] (6)

During the one-dimensional compression (or expansion) of the current sheet magnetic field remains one-dimensional,
\[
\mathbf{B} = (B(z,t),0,0).
\]
In addition to the $y$ component of the velocity $v$, the $z$ component becomes time-dependent as well

$$v = (0, v_y(z, t), v_z(z, t)).$$

The plasma density is $\rho = \rho(z, t)$, its pressure is $P = P(z, t)$. We use the equation of state $P = C_2 \rho^\Gamma$, where $\Gamma$ is the adiabatic exponent. From the first and the fourth equations of system (6), we obtain

$$\rho \frac{dB}{dt} + \rho v_z \frac{dB}{dz} = -\rho B \frac{\partial v_z}{\partial z} = B \frac{\partial \rho}{\partial t} + B v_z \frac{\partial \rho}{\partial z},$$

which leads to the following condition

$$\rho \frac{dB}{dt} = B \frac{d\rho}{dt},$$

which means that the magnetic field is “frozen” into the plasma. Hence, we have $B \propto \rho$, so let $B = C_1 \rho$. Now, the third and the fourth equations of system (6) describe the evolution of the velocity $v_z(z, t)$ and plasma density $\rho(z, t)$

$$\rho \left( \frac{\partial v_z}{\partial t} + v_z \frac{\partial v_z}{\partial z} \right) + C_2 \Gamma \rho^{\Gamma - 1} \frac{\partial \rho}{\partial z} + \frac{C_2^2 \rho \partial \rho}{4\pi \partial z} = 0,$$

$$\frac{\partial \rho}{\partial t} + v_z \frac{\partial \rho}{\partial z} + \rho \frac{\partial v_z}{\partial z} = 0. \tag{7}$$

The change of variables $(\rho, v_z) \rightarrow (t, z)$ reduces the system of non-linear equations (7) to the following linear system

$$\frac{\partial z}{\partial \rho} - v \frac{\partial t}{\partial \rho} + C_2 \Gamma \rho^{\Gamma - 2} \frac{\partial t}{\partial v} + \frac{C_2^2 \rho \partial t}{4\pi \partial v} = 0,$$

$$\frac{\partial z}{\partial v} - v \frac{\partial t}{\partial v} + \rho \frac{\partial t}{\partial \rho} = 0. \tag{8}$$

$$\frac{\partial z}{\partial \rho} - v \frac{\partial t}{\partial \rho} + C_2 \Gamma \rho^{\Gamma - 2} \frac{\partial t}{\partial v} + \frac{C_2^2 \rho \partial t}{4\pi \partial v} = 0,$$

$$\frac{\partial z}{\partial v} - v \frac{\partial t}{\partial v} + \rho \frac{\partial t}{\partial \rho} = 0. \tag{9}$$

To simplify the notation we drop the index of $v_z$. Eliminating the derivatives of $z$ we obtain one equation for $t(\rho, v)$

$$\rho \frac{\partial^2 t}{\partial \rho^2} + 2 \frac{\partial t}{\partial \rho} - C_2 \Gamma \rho^{\Gamma - 2} \frac{\partial^2 t}{\partial v^2} - \frac{C_2^2 \partial^2 t}{4\pi \partial v^2} = 0. \tag{10}$$

Solution of this equation cannot be expressed in terms of known special functions for the arbitrary values of $\Gamma \neq 2$. However, we can obtain a quasiclassical approximation to the general solution after applying Laplace transform to the variable $v$. When $v > 0$, the quasiclassical solution to (10) is

$$t(\rho, v) = -\frac{1}{2\rho \sqrt{\xi(\rho)}} \int_{\rho^-}^{\rho^+} f(\rho_1) d\rho_1. \tag{11}$$

Here $\xi(\rho) = (C_2 \Gamma \rho^{\Gamma - 3} + C_2^2/(4\pi \rho))^{1/2} = (\Gamma P + B^2/(4\pi))^{1/2} \rho^{-3/2} = c_s/\rho$, and $c_s$ is the fast magnetic sound speed. The function $f(\rho)$ appearing in the solution depends on the initial conditions. We may assume that the velocity $v$ is zero when $t = 0$. Accordingly, $\partial v/\partial z|_{t=0} = 0$ as well. It is convenient to write the initial condition for the density $\rho$ at $t = 0$ in the following way $f_{\rho_0} f(\rho) d\rho = z$, where
where \( \rho_0 \) is the density at \( z = 0 \). From the continuity equation of system (7) it follows that \( \partial \rho / \partial t \big|_{t=0} = 0 \). From the first equation of the same system (the equation of motion) we see that \( \partial v / \partial t \big|_{t=0} = -\rho \xi^2(\rho) / f(\rho) \). After we change variables to \( (t(v, \rho), z(v, \rho)) \), our initial conditions become

\[
\begin{align*}
\frac{\partial t}{\partial v} \bigg|_{v=0} &= -\frac{f(\rho)}{\rho \xi^2(\rho)}; \quad \frac{\partial t}{\partial \rho} \bigg|_{v=0} = 0; \\
\frac{\partial z}{\partial v} \bigg|_{v=0} &= 0; \quad \frac{\partial z}{\partial \rho} \bigg|_{v=0} = f(\rho); \quad t \bigg|_{v=0} = 0.
\end{align*}
\]

In the obtained solution, \( \rho_+ \) satisfies the equation \( v = \int_{\rho_0}^{\rho_+} \xi(\rho_1) d\rho_1 \), and \( \rho_- \) satisfies the equation \( v = \int_{\rho_0}^{\rho_-} \xi(\rho_1) d\rho_1 \). For \( v > \int_{\rho_0}^{\rho_-} \xi(\rho_1) d\rho_1 \) the value \( \rho_- \) cannot be determined, because at the limiting value \( v = \int_{\rho_0}^{\rho_-} \xi(\rho_1) d\rho_1 \) the density \( \rho_- \) vanishes and cannot be negative later. Since \( \xi = c_s / \rho \), the condition \( v > \int_{\rho_0}^{\rho_-} \xi(\rho_1) d\rho_1 \) is virtually identical to \( v > c_s \), and quasiclassical solution (11) is not valid.

Substituting (11) into (9) gives the following approximation for \( z(v, \rho) \)

\[
z(v, \rho) = vt(v, \rho) - \frac{\partial}{\partial \rho} \left[ \rho \int_0^v t(v_1, \rho) dv_1 \right] + \int_{\rho_0}^{\rho} f(\rho_1) d\rho_1. \tag{12}
\]

The integrals appearing in expressions (11, 12) have their simplest form when \( \Gamma = 2 \), which is close to the adiabatic exponent for an ideal gas \( \Gamma = 5/3 \). In this case the solution is

\[
\begin{align*}
t(v, \rho) &= -\frac{4f_0}{\alpha^2} v p^k, \\
z(v, \rho) &= \frac{f_0(\rho^{k+1} - \rho_0^{k+1})}{k + 1} + \frac{(k - 1)\alpha^2}{8f_0} \rho^{-k} t^2. \tag{13}
\end{align*}
\]

Here we have chosen the function \( f(\rho) \) to be the power function \( f_0 \rho^k \) where \( k \) is an arbitrary value and \( \alpha = 2(2C_2 + C_1^2 / 4\pi)^{1/2} \). Solutions \( t(v, \rho) \) and \( z(v, \rho) \) (13) are always single-valued functions of \( (v, \rho) \), but the functions \( v(t, z) \) and \( \rho(t, z) \) may not be. The quantity \( \rho(t, z) \) can be found by reversing the second equation of system (13). This can be uniquely done until the moment when the shock occurs and the Jacobian of the transformation of the unknown variables vanishes

\[
J = \frac{\partial (t, z)}{\partial (\rho, v)} = \frac{4f_0^2}{\alpha^2} p^{2k} \left( 1 + 2k(1 - k) \frac{v^2}{\alpha^2 p} \right) = 0.
\]

For \( k < 0 \) and \( k > 1 \) the Jacobian \( J \) always vanishes at a value \( q = v^2 / (\alpha^2 p) \). When the shock occurs, the derivatives \( \partial \rho / \partial z \) and \( \partial \rho / \partial z \) become infinite which leads to the singularity in the electromagnetic field. The question whether there is a shock for almost arbitrary initial perturbations is very important, so we would like to obtain an exact solution of the MHD equations, and it is only possible for \( \Gamma = 2 \). This solution can be expressed in terms of hypergeometric
The hypergeometric function $2F_1(a, b, c, q)$ appearing in these equations has a branch point at $q = 1$, which means that solution (14) cannot be continued into the region $q > 1$.

Note that quasiclassical solution (13) for $\Gamma = 2$ corresponds to the first two terms in the expansion of exact solution (14) into powers of $q$.

Both in the exact solution and in the quasiclassical one, the singularity occurs for $0 < q < 1$ when the Jacobian $J$ vanishes (the shock point). However, in the exact solution the shock occurs only for negative values of $k$, namely for $k < -1/2$. These values of $k$ correspond to the more natural initial conditions when the initial value of the acceleration of the sheet $\partial v/\partial t|_{t=0} \propto \rho^{-k}$ does not have a singularity at $\rho = 0$. The behavior of the Jacobian $J$ of $q = v^2/(\alpha^2 \rho)$ for the exact and quasiclassical solutions is shown in Fig. 2.

The figure shows that the approximate solution describes the behavior of the sheet well right up to the shock point, so later we will use only quasiclassical solution (13).

An expression for the $y$ component of the plasma velocity should be added to formulas (13). It can be easily found from the equation $\partial v_y/\partial t + v \partial v_y/\partial z = 0$, whose solution is an arbitrary function of the variable $\rho^{k+2} \eta$, $v_y = \Phi(\rho^{k+2} \eta)$, where

$$\eta = 1 + \frac{k(k+2)\alpha^2}{8 f_0^2} \rho^{-2k-1} t^2 = 1 + 2k(k+2)q.$$  

The function $\Phi$ is determined by an initial condition for the motion of the sheet in the $y$ direction at $t = 0$. The initial velocity $v_y$ can be arbitrary for an
infinitely wide sheet. But since the real geometry of our problem is not flat, we
may assume that \( v_y \sim v = v_z \).

The electric and magnetic fields are given by
\[
B_x = B = C_1 \rho, \quad \frac{E_y}{B} = -\frac{v}{c} = \frac{x^2}{4f_0 c^2} p^{-k_t}, \quad \frac{E_z}{B} = \frac{v_y}{c} = \frac{\Phi}{c}.
\]

and the other components vanish.

4 Acceleration of charged particles

Evolution of the current sheet causes a drastic change of the magnetic field, its energy, and the appearance of the electric field. Due to the high speed of this process (it lasts for only about \( L/c \simeq 0.1 \text{sec} \)), the energy of the magnetic field is passed to the energy of the accelerated particles. Acceleration is most effective in a region with the fastest evolution, i.e. near the singularity. Let us consider acceleration of the charged particles in the vicinity of singularity. We assume that these particles are relativistic, so that \( \varepsilon = c p \), and that their fraction is small compared to the plasma density in the sheet. The singularity appears in the above solution of the MHD equations during the compression of the current sheet. The distribution function \( F_*(t, z, p_x, p_y, p_z) \) of relativistic particles in Cartesian coordinates satisfies collisionless Liouville’s equation
\[
\frac{\partial F_*}{\partial t} + \frac{e}{\varepsilon} p_z \frac{\partial F_*}{\partial z} + F_0 \frac{\partial F_*}{\partial p_\perp} = 0.
\]

Let us change the Cartesian coordinates \((p_x, p_y, p_z)\) to the cylindrical ones \((p_\parallel, p_\perp, \varphi)\) in the momentum space. Here \( p_\parallel \) is a component of the particle momentum along the magnetic field \( B = B_x \), \( p_\perp \) is that perpendicular to the magnetic field, and \( \varphi \) is a rotation angle of the particle around the magnetic field,
\[
p_x = p_\parallel, \quad p_y = p_\perp \sin \varphi, \quad p_z = p_\perp \cos \varphi.
\]
The force acting on the particles is the Lorentz force
\[
\mathcal{F} = e \left( E + \frac{\varepsilon}{c} [pB] \right).
\]

Thus we have the kinetic equation
\[
\frac{\partial F_*}{\partial t} + \frac{e}{\varepsilon} p_\perp \cos \varphi \frac{\partial F_*}{\partial z} + e \left( E_y \sin \varphi + E_z \cos \varphi \right) \frac{\partial F_*}{\partial p_\perp} + \\
\left[ \frac{e}{p_\perp} (E_y \cos \varphi - E_z \sin \varphi) + \omega_c \right] \frac{\partial F_*}{\partial \varphi} = 0,
\]

where \( \omega_c \) is the cyclotron frequency of the particle rotation, \( \omega_c = eB/p \). Let us expand the function \( F_* \) of \( \varphi \) into a Fourier series and leave only the first 3 terms
\[
F_* = F_0 + F_c \cos \varphi + F_s \sin \varphi.
\]
Hence we get the system of equations for the functions $F_0$, $F_c$, $F_s$

$$
\begin{align*}
\frac{\partial F_0}{\partial t} &+ \frac{c p_\perp}{2 p} \frac{\partial F_0}{\partial z} + \frac{1}{2} \frac{e E_y}{p} F_s + \frac{e E_y}{2} \frac{\partial F_s}{\partial p_\perp} + \\
\frac{1}{2} \frac{e E_y}{p} F_c + \frac{e E_y}{2} \frac{\partial F_c}{\partial p_\perp} &= 0, \\
\frac{\partial F_c}{\partial t} &+ \frac{c p_\perp}{2 p} \frac{\partial F_0}{\partial z} + \frac{e E_y}{2} \frac{\partial F_0}{\partial p_\perp} + \omega_c F_s = 0, \\
\frac{\partial F_s}{\partial t} &+ \frac{e E_y}{2} \frac{\partial F_0}{\partial p_\perp} - \omega_c F_c = 0.
\end{align*}
$$

Equation (17) has a simple physical meaning. All particles move in the $z$ direction with a velocity of the electric drift $-c E_y/B$, moreover the particles have an acceleration in this direction due to the variation of this velocity along the $z$ axis, and hence gain the transverse momentum $p_\perp$. It is important to note that the acceleration $\partial (c E_y/B)/\partial z$ tends to infinity at the shock point,

$$
\frac{\partial}{\partial z} \left( \frac{E_y}{B} \right) = -\frac{4kq}{t} (1 + 2k(1-k)q)^{-1}.
$$

The dependence of $E_y/B$ on the coordinate $z$ and the time $t$ is given by the formula (16). Now, it will be convenient to change variables from $(t, z, p_\perp)$ to $(\eta, \rho, p_\perp)$. In the new variables (17) reads

$$
(k + 2) \eta \frac{\partial F_0}{\partial \eta} - \rho \frac{\partial F_0}{\partial \rho} - \frac{1}{4} p_\perp \frac{\partial F_0}{\partial p_\perp} = 0. \tag{18}
$$

We may assume that the distribution function of fast plasma particles is Maxwellian at $t = 0 (\eta = 1)$, and their density is proportional to the plasma density, i.e. the fraction of the accelerated injected particles is constant along the entire sheet

$$
F_0(\eta = 1, \rho, p_\perp) = c_0 \rho e^{-c p_\perp/T_0}.
$$
The solution of kinetic equation (18) with such initial conditions is

\[ F_0(\eta, \rho, p_\perp) = c_0 \rho_0 \eta^{1/(k+2)} \exp \left( -cp_\perp \eta^{1/(2k+4)} / T_0 \right). \]

It shows that the distribution function of the particles remains Maxwellian, but the temperature is time-dependent \( T = T_0 \eta^{-1/(2k+4)} \). For \( -2 < k < 0 \) the temperature is rising as \( \eta \) tends to zero. At the moment \( \eta = 0 \) the temperature becomes infinite, which means that the distribution function is constant in the region \( p_\perp < \infty \). Note that the moment \( \eta = 0 \) occurs before the shock for \(-1/2 < k < 0\) and these moments coincide if \( k = -1/2 \). The case \( k = -1/2 \) corresponds to the initial distribution of density \( \rho(z) \) with constant initial plasma temperature along the entire sheet. Though formally the particles are accelerated up to infinite energies, it is clear that the cyclotron radius of the fastest particles should be less than the typical thickness of the sheet \( L \), and also, that their cyclotron frequency should be greater than the inverse typical evolution time \( \Delta t \) of the sheet. Therefore, we have two estimates for \( p_\perp m \): \( p_\perp m = p_{1m} = eL < B / c \) and \( cp_{2m} = T_0 \Delta \eta^{-1/(2k+4)} \), where \( \Delta \eta \) can be found from the definition of \( \eta \) and is estimated to be \( \Delta \eta \approx (c_\eta / (L \omega_0))^2 (p_{2m} / m c)^2 \). \( \omega_0 = eB/mc \). Finally we have \( p_{2m} / m c = (T_0 / (mc^2))^{(k+2)/(k+3)} (L \omega_0 / c_\eta)^{1/(k+3)} \). Later we will see that \( p_{1m} > p_{2m} \) for the actual values of the parameters, so we can assume that \( p_{m} = p_{2m} \).

### 5 Synchrotron radiation

The most interesting case for us is when \( k = -1/2 \), and in this case all plasma masses up at the moment of shock near the two planes \( z_{+,-} = \pm f_0 \sqrt{\rho_0} \), symmetric with respect to the center of the sheet. The estimate for \( z_+ \) is \( 2c_\eta^2 / \partial v / \partial t \big|_{t=0} \), it is reciprocal to the initial acceleration of the sheet \( \partial v / \partial t \big|_{t=0} \), and, of course, \( z_+ \) should be less than the thickness of the sheet \( L \). The plasma velocity \( v_z \) is positive for \( z > 0 \) and \( v_z < 0 \) for \( z < 0 \), i.e. plasma particles move away from the center of the sheet into the regions around \( z_{+,-} \) with a strong magnetic field. Due to the strong magnetic field, the synchrotron radiation becomes significant for the plasma particles (electrons) there. For \( z > z_+ \) and \( z < z_- \) the quasiclassical approximation used in the solution of the MHD equations is not applicable. We assume that the magnetic field is almost constant at the large distances \( z > z_+ \) and \( z < z_- \) from the center of the sheet, and its value is equal to the magnetic field at the boundary of the sheet. The electric field with the components \( E_y \) and \( E_z \) exists there as well, and plasma velocity in the electromagnetic field should be \( v \sim c_s \), where \( c_s \) is the fast magnetic sound speed.

Thus, fast particles with the constant distribution function obtained above move into the regions with constant fields due to the electric drift and strongly radiate there. Next, we find \( F \), the distribution function of these particles, and their radiation intensity \( I \).

We may assume that \( E_y / B \ll 1 \) and \( E_z / B \ll 1 \), i.e. the plasma motion is non-relativistic. In this case the drift velocity is \( \mathbf{v}_d = (0, cE_z / B, -cE_y / B) \).

The function \( F \) satisfies the following continuity equation in the phase space

\[
\frac{\partial F}{\partial t} + c \frac{p_z}{\rho} \frac{\partial F}{\partial z} + \frac{\partial F_\perp}{\partial p_\perp} F + F_\perp \frac{\partial F}{\partial p_\perp} = 0.
\]
The force $F_i$ acting on the particles has the simplest expression in the reference frame moving with the velocity of the electric drift, where there is no electric field. Since the drift is non-relativistic, this force is the same in the stationary reference frame. We use tildes for the values observed in the moving frame. The friction force due to radiation is

$$F_{fr} = -\frac{2e^4\tilde{B}^2}{3m^4c^6} \tilde{p} \left( \tilde{p}_y^2 + \tilde{p}_z^2 \right) \left( \tilde{p}_x, \tilde{p}_y, \tilde{p}_z \right),$$

note that the magnetic field in the moving reference frame is $\tilde{B} = B \left( 1 - \frac{(E_y^2 + E_z^2)}{B^2} \right)$, and components of the electric field are $\tilde{E}_y = 0$, $\tilde{E}_z = 0$. Let us denote

$$\kappa = \frac{2e^4}{3m^4c^6}.$$ 

The net force acting on the particles is

$$\mathcal{F}_x = \mathcal{F}_{fr,x},$$

$$\mathcal{F}_y = e\tilde{B} \tilde{p}_y + \mathcal{F}_{fr,y},$$

$$\mathcal{F}_z = -e\tilde{B} \tilde{p}_y + \mathcal{F}_{fr,z}.$$ 

Again, change the Cartesian coordinates $(p_x, p_y, p_z)$ to the cylindrical ones $(p_\parallel, p_\perp, \varphi)$ in the momentum space

$$p_x = \tilde{p}_x = p_\parallel,$$

$$p_y = \frac{E_z}{B} \tilde{p} + \tilde{p}_y = \frac{E_z}{B} \tilde{p} + p_\perp \sin \varphi,$$

$$p_z = -\frac{E_y}{B} \tilde{p} + \tilde{p}_z = -\frac{E_y}{B} \tilde{p} + p_\perp \cos \varphi.$$ 

Notice that a non-relativistic momentum addition rule was used here as the drift is non-relativistic.

Further, we neglect terms of the order higher than $E_y/B$ and $E_z/B$ in all calculations. Under this approximation $\tilde{B} = B$,

$$\mathcal{F}_i \frac{\partial F}{\partial p_i} = \mathcal{F}_i \frac{\partial F}{\partial \tilde{p}_i} - \mathcal{F}_y \tilde{p}_y \mathcal{E}_z + \mathcal{F}_z \tilde{p}_z \mathcal{E}_y \frac{\partial F}{\partial \tilde{p}_y} + \mathcal{F}_y \tilde{p}_y \mathcal{E}_z + \mathcal{F}_z \tilde{p}_z \mathcal{E}_y \frac{\partial F}{\partial \tilde{p}_z}. \quad (19)$$

Since divergence is a scalar value and the Lorentz force is solenoidal, we have

$$\frac{\partial \mathcal{F}_i}{\partial \tilde{p}_i} = \frac{\partial \mathcal{F}_i}{\partial \tilde{p}_i} = \frac{\partial \mathcal{F}_{fr,i}}{\partial \tilde{p}_i} = -4\kappa B^2 \frac{p_i^2}{\tilde{p}}.$$

Substituting the expressions obtained above in the original kinetic equation, and averaging it over the rotation angle $\varphi$, we obtain

$$\frac{\partial F}{\partial t} - \frac{eE_y}{B} \frac{\partial F}{\partial z} - \kappa B^2 \tilde{p}_y^2 \frac{\partial F}{\partial \tilde{p}} \left( \tilde{p}_x \frac{\partial F}{\partial \tilde{p}_x} + \tilde{p}_y \frac{\partial F}{\partial \tilde{p}_y} + \tilde{p}_z \frac{\partial F}{\partial \tilde{p}_z} \right) - 4\kappa B^2 \tilde{p}_y^2 F = 0. \quad (20)$$

The identity $\tilde{p} = p$ is used here, it can be proved by expanding $\tilde{p}$ into a Taylor series in powers of $E_y/B$ and $E_z/B$ and averaging over the angle $\varphi$. 

11
Let us consider the case \( z > 0 \) (\( z < 0 \) is analogous due to the symmetry of the problem). The particles with constant initial distribution function move into the region \( z > z_+ \) after accumulating at \( z_+ = 2|f_0|\sqrt{\rho_0} \), according to the kinetic equation obtained above. We can suppose without loss of generality that the origin of the \( z \) coordinate is at the point \( z_+ \), i.e. we can solve the kinetic equation with the initial condition

\[
F(t = t_*; z = 0; p_\parallel; p_\perp) = c_1 \Theta(p_m - p).
\]

Here \( t_* \) is the time of the constant distribution formation. Let us notice that the constant distribution function of accelerated particles obtained in the previous section was the function of the transverse momentum \( p_\perp \). Here, in contrast, we use the isotropic function with the same parameters. The reason is that the plasma with strongly anisotropic particle distribution is very unstable with respect to the excitation of the electrostatic waves leading to fast isotropization. The growth rate of such instabilities is of order of the cyclotron frequency \( \omega_c \) or the plasma frequency \( \omega_p \) [10]. We are looking for a stationary distribution function (the solution of (20) which does not depend on time) to find the spectrum of the radiation of the particles. This solution is

\[
F(z; p_\parallel; p_\perp) = c_1 \Theta\left( p_m - \frac{p}{1 + \frac{k_B^2 z^2 p_\perp^2}{c^2 E_p^2}} \right).
\]

Having found the distribution function, we can compute the spectrum and the directivity of the radiation, i.e. dependence of the radiation intensity on the direction and on the frequency.

First, let us find the synchrotron radiation from one particle moving with the velocity \( v \). The prime denotes that the variable is observed in the coordinate frame \( K' \) moving with the velocity \( v_\parallel \). To find the radiation, we will need to change the coordinate frame from the stationary one to \( K' \) and back. We denote the angle between the direction of the magnetic field and the velocity of the particle by \( \chi \) and the angle between the plane of the particle rotation and the direction of the radiation by \( \theta \). The radiation in the \( K' \) coordinate frame is

\[
dI' = d\Omega' \frac{e^4 B'^2 v'^2(1 - v'^2/c^2)}{8\pi^2 m^2 e^5} \times \frac{2 - \cos^2 \theta' - \frac{v'^2}{c^2} \left( 1 + \frac{3v'^2}{c^2} \right) \cos^4 \theta'}{\left( 1 - \frac{v'^2}{c^2} \cos^2 \theta' \right)^{7/2}}.
\]

Here

\[
d\Omega' = 2\pi \cos \theta' d\theta', \quad B' = \gamma B, \quad v'_\perp = \gamma v_\perp;
\]

\[
v_\perp = v \sin \chi, \quad v_\parallel = v \cos \chi, \quad \gamma = 1/\sqrt{1 - \beta^2}, \quad \beta = v_\parallel / c.
\]

Therefore, the radiation in the stationary coordinate frame is

\[
dI = (1 + \beta \sin \theta') dI', \quad \sin \theta = \frac{\beta + \sin \theta'}{1 + \beta \sin \theta'}, \quad \frac{d\theta}{\cos^2 \theta} = \gamma (1 + \beta \sin \theta') \frac{d\theta'}{\cos^2 \theta'}.
\]
\[
\frac{dI}{d\theta} = \gamma^3 \frac{\cos^3 \theta'}{\cos^2 \theta} \frac{e^4 B^2 v_\perp^2 (1 - \gamma^2 v_\perp^2 / c^2)}{4\pi m^2 c^5} \times \left[ 2 - \cos^2 \theta' - \frac{\gamma^2 v_\perp^2}{4c^2} \left( 1 + \frac{3\gamma^2 v_\perp^2}{c^2} \right) \cos^4 \theta' \right] / \left( 1 - \frac{\gamma^2 v_\perp^2}{c^2} \cos^2 \theta' \right)^{7/2}.
\]

The sought dependence of the radiation intensity on the direction \(\theta\) is

\[
\left( \frac{dI}{d\theta} \right) = \int_{p_{\min}}^{p_{\max}} dp \int_0^\pi p^2 dp \frac{\theta}{\sin \theta} \frac{\cos^3 \theta'}{\cos^2 \theta} \left( \frac{e^4 B^2 v_\perp^2 (1 - \gamma^2 v_\perp^2 / c^2)}{4\pi m^2 c^5} \times \left[ 2 - \cos^2 \theta' - \frac{\gamma^2 v_\perp^2}{4c^2} \left( 1 + \frac{3\gamma^2 v_\perp^2}{c^2} \right) \cos^4 \theta' \right] / \left( 1 - \frac{\gamma^2 v_\perp^2}{c^2} \cos^2 \theta' \right)^{7/2} \right) \times \left( \frac{c_1}{1 + \frac{\gamma B^2}{cE_y} p \sin^2 \chi} \right)^2 \Theta \left( p_m - \frac{p}{1 + \frac{\gamma B^2}{cE_y} p \sin^2 \chi} \right) \times 2\pi \sin \chi d\chi. \quad (22)
\]

Note, that only the particles with \(\gamma \gg 1\) make significant contribution to the radiation. In the \(K'\) coordinate frame the synchrotron radiation is concentrated near the plane of the particle rotation, which means that \(\theta' = 0\). The most part of the radiation is concentrated within the angle \(\Delta \theta' \simeq 1/\gamma \ll 1\). Note that \(\Gamma = (1 - \gamma^2 v^2 / c^2)^{-1/2} \gg 1\) since \(\gamma \gg 1\).

From Equation (21) it follows that \(\sin \theta \sim \beta = \frac{e}{c} \cos \chi \simeq \cos \chi\). Hence the particles with the pitch angle \(\chi\) radiate in the stationary coordinate frame mainly into the angle

\[
\theta = \frac{\pi}{2} - \chi. \quad (23)
\]

Let us perform the Taylor expansion of all expressions appearing in the integral over \(\chi\) in (22) in powers of \(1 - v^2 / c^2\) and leave only the first, most essential terms. Using the smallness of the angle \(\theta'\), we get

\[
\gamma \simeq 1/\sin \chi, \quad \cos \theta' \simeq 1, \quad v_\perp \simeq c \sin \chi,
\]

\[
(1 - \gamma^2 v_\perp^2 / c^2) \simeq \left( 1 - \frac{\gamma^2 v_\perp^2}{c^2} \cos^2 \theta' \right) \simeq \frac{1}{\Gamma^2 \sin^2 \chi},
\]

\[
\left[ 2 - \cos^2 \theta' - \frac{\gamma^2 v_\perp^2}{4c^2} \left( 1 + \frac{3\gamma^2 v_\perp^2}{c^2} \right) \cos^4 \theta' \right] \simeq \frac{7}{4\Gamma^2 \sin^2 \chi}.
\]

Substituting these equations into the integral (22) and using (23), we obtain that the integrand is approximately equal to

\[
\frac{7}{8} \Theta \left( p_m - \frac{p}{1 + \frac{\gamma B^2}{cE_y} p \cos^2 \theta} \right) \frac{e^4 B^2 v_\perp^2}{m^2 c^5} \times \left( \frac{\cos \theta}{1 + \frac{\gamma B^2}{cE_y} p \cos^2 \theta} \right)^4 \frac{d\chi}{\sin \chi}.
\]

Numerical computation shows that the radiation at the angle \(\theta\) is produced mainly by the particles with \(\Delta \chi \simeq 1.52/\Gamma\), i.e. the integral over \(d\chi\) is close to

\[
1.33 c_1 \frac{e^4 B^2 v_\perp^2}{m^2 c^5} \frac{\cos \theta}{\Gamma^2} \times \left( \frac{\cos \theta}{1 + \frac{\gamma B^2}{cE_y} p \cos^2 \theta} \right)^4 \Theta \left( p_m - \frac{p}{1 + \frac{\gamma B^2}{cE_y} p \cos^2 \theta} \right).
\]
The upper limit of the integration over $p$ in Equation (22) is

$$p = p_{cr} = \frac{p_m}{1 - \frac{\kappa B z}{eE_y} p_m \cos^2 \theta}.$$ 

Notice that $p_{cr} < p_m$ since $E_y/B < 0$. The lower limit of the integration, $p_{min}$, can be replaced by zero. Hence

$$\langle \frac{dI}{d\theta} \rangle = 1.33 c \int_0^z_{z_{max}} dz \int_0^{p_{cr}} p^4 \left( 1 + \frac{\kappa B^3 z}{eE_y} p \cos^2 \theta \right)^{-4} dp.$$ 

These two integrals can be evaluated analytically:

$$\langle \frac{dI}{d\theta} \rangle = 1.33 c_1 \frac{e^4 B^2}{m^4 c^5} \frac{p_{m}}{K}^5 \frac{z_{max}}{\cos \theta} \left( 3K^4 - \frac{4K^3}{\cos^2 \theta} + \frac{6K^2}{\cos^4 \theta} - \frac{12K}{\cos^6 \theta} + \frac{12}{\cos^8 \theta} \log(1 + K \cos^2 \theta) \right),$$

where $K = -(2e^2 B^3/3m^4 c^7 E_y) p_m z_{max} \gg 1$. Angular distribution of the intensity $\langle dI/d\theta \rangle$ is maximal at $\cos \theta \approx K^{-1/2}$, it vanishes at $\cos \theta = 0$, and has the width maximum $\Delta(\cos \theta)$ close to $K^{-1/2}$. That is, the radiation is strongly collimated along the magnetic field and the collimation angle $\Delta \theta$ can be estimated as $K^{-1/2} \ll 1$. Let us find the radiation spectrum. For a single particle the spectral intensity of the synchrotron radiation is [7]

$$\frac{dI}{d\omega} = \sqrt{\frac{3}{2\pi}} \frac{e^3 B}{mc^2} F\left(\frac{\omega}{\omega_c}\right).$$

Here

$$F(\xi) = \xi \int_\xi^\infty K_{5/3}(\xi) d\xi,$$

$$\omega_c = \frac{3eB \xi}{2mc} \left( \frac{\varepsilon}{mc^2} \right)^2 \approx \frac{3eB}{2m^3 c^3} p^2 \sin \chi,$$

$$B_{\perp} = B \sin \chi.$$ 

The total spectrum is given by

$$\langle \frac{dI}{d\omega} \rangle = c_1 \int_0^{z_{max}} dz \int_0^{p_{cr}} \left( 1 + \frac{\kappa B^3 z}{eE_y} p \sin^2 \chi \right)^{-4} \frac{dI}{d\omega} p^2 dp,$$

where

$$p_{cr} = p_m \left( 1 - \frac{\kappa B^3 z}{eE_y} p_m \sin^2 \chi \right)^{-1}.$$

Asymptotic behavior of the function $F(\xi)$ is known to be

$$F(\xi) \sim \sqrt{\frac{\pi \xi}{2}} e^{-\xi} as \xi \to \infty.$$ 

Using Laplace’s method, we have the asymptotics when $\omega \to \infty$

$$\int_{p_{min}}^{p_{cr}} \left( 1 + \frac{\kappa B^3 z}{eE_y} p \sin^2 \chi \right)^{-4} \frac{dI}{d\omega} p^2 dp \simeq \frac{3}{8} \times$$

$$\left( \frac{e^3 B}{\pi m^3 c^5 \omega} \right)^{1/2} p_m^4 \sin^{3/2} \chi \exp \left( -\frac{2m^3 c^3}{3eB} \frac{1}{p_{cr}^2} \sin \chi \right).$$
After substituting the expression for $p_{cr}$, we obtain the asymptotics for the two other integrals when $\omega \to \infty$

\[
\int_{0}^{z_{max}} dz \int_{p_{min}}^{p_{cr}} (...) dp \sim \frac{9c_{1}}{32\pi^{3/2}} \omega^{-3/2} p_{m}^{3} \times \left( \frac{m^{3}c^{7}E_{y}^{2}}{eB^{3}\sin\chi} \right)^{1/2} \exp \left( -\omega \frac{2m^{3}c^{3}}{3eB p_{m}^{2}} \sin\chi \right).
\]

Finally, we obtain the asymptotic of the spectral intensity

\[
\langle \frac{dI}{d\omega} \rangle \approx \frac{27\pi\sqrt{3}}{32} c_{1} p_{m}^{6} \frac{e|E_{y}|}{m^{3}c} \frac{1}{\omega^{2}} \exp \left( -\omega \frac{2m^{3}c^{3}}{3eB p_{m}^{2}} \right) \text{ as } \omega \to \infty.
\]

The exponential damping of the spectrum begins at the frequency

\[
\omega_{*} = \frac{3eB}{2mc} \left( \frac{p_{m}}{mc} \right)^{2}.
\]

Now let us find the asymptotics of the spectrum as $\omega \to 0$. Recall that the radiation is mainly concentrated in the angle $K^{-1/2}$ radian. Only the particles with $\chi \sim K^{-1/2}$ make a substantial contribution (equation (23)). The asymptotics of $F(\xi)$ is

\[
F(\xi) \sim 2^{2/3}\Gamma \left( \frac{2}{3} \right) \xi^{1/3} \text{ as } \xi \to 0.
\]

Therefore, when $\omega \to 0$, we have

\[
\int_{p_{min}}^{p_{cr}} \frac{dI}{d\omega} p^{2} \left( 1 + \frac{\kappa B^{3}z}{E_{y}} p \sin^{2}\chi \right)^{-4} dp \approx \frac{3^{7/6}\Gamma(2/3)}{7\pi} \frac{e^{8/3}B^{2/3}/c}{\omega^{1/3}p_{cr}^{7/3}\sin^{2/3}\chi}.
\]

Finally, after integrating over $\chi$ and $z$, we obtain

\[
\langle \frac{dI}{d\omega} \rangle \approx c_{1} \frac{27}{14} \frac{2^{3/2}}{\sqrt{\pi}} \cdot \Gamma \left( \frac{7}{6} \right) \frac{m^{4}e^{6}|E_{y}|}{e^{4/3}B^{2/3}p_{m}^{4/3}\omega^{1/3}} \text{ as } \omega \to 0.
\]

The asymptotic behavior of $\langle dI/d\omega \rangle$ changes from $\omega^{1/3}$ to $\omega^{-2}$ somewhere in the region $\omega < \omega_{*}$. An approximate value of the frequency $\omega = \omega_{0}$ when this change occurs is given by

\[
\omega_{0} \approx 0.8 \frac{eB}{mc} \left( \frac{p_{m}}{mc} \right)^{2}.
\]

Thus we see that the radiation spectrum of the accelerated particles of the whole sheet is maximal at $\omega \approx \omega_{0}$ with power gain for $\omega < \omega_{0}$ and power damping for $\omega > \omega_{0}$. The exponential damping occurs for the large frequencies $\omega > \omega_{*}$.

### 6 Summary

According to the papers of [3], the typical moderate values for the sheet are:

- the typical velocity of the fast magnetic sound is $c_{s} \approx 10^{9}$ cm/s,
- the thickness of the sheet is $L \approx 10^{8}$ cm,
- the typical value of the magnetic field is $B \approx 10^{6}$ G,
the plasma temperature in the sheet is $T_0 \simeq 10 \text{ keV}$.

The time of the sheet compression can be estimated from Equation (15) and the condition $z_+ \sim L$

$$\Delta t \sim \frac{f_0}{\alpha} \sim c_s \left| \frac{\partial v}{\partial t} \right|_{t=0} \sim \frac{L}{c_s} \sim 0.1 \text{ sec.}$$

Now we can estimate the maximum value of the Lorentz factor for the particles ($k = -1/2$):

$$\Gamma_m = \frac{p_m}{mc} = \left( \frac{T_0}{mc^2} \right)^{3/5} \left( \frac{L\omega_0}{c_s} \right)^{2/5} \simeq 10^4.$$  

The collimation $\Delta \theta$ of the radiation is determined by the parameter $K$, $\Delta \theta \simeq K^{-1/2}$,

$$K = \frac{2 \Gamma_m \omega_0^2 Lr_e}{c^2 \left| E_y \right|}.$$  

$r_e$ is the classical electron radius, $r_e = 3 \cdot 10^{-13} \text{ cm}$. Assuming $\left| E_y \right|/B \simeq c_s/c$, we obtain $K \simeq 10^7$, which yields the collimation angle, $\Delta \theta \simeq 3 \cdot 10^{-4}$. The intensity of the accelerated particle radiation in the sheet is maximal at the frequency $\omega_0 \simeq 8 \cdot 10^{20} \text{s}^{-1}$, which corresponds to the energies of the order of 500 keV. Asymptotic behavior of the spectrum as well as the frequency $\omega_0$ agree with the observed values of GRBs [2].

Thus, we see that the more precise calculations for the model of cosmological GRBs proposed by [5] confirm their estimations. Evolution of the current sheet, particle acceleration due to the magnetic energy release, and the synchrotron radiation turn out to agree with what was estimated. The total radiated energy corresponds to the release of the magnetic energy stored in the sheet ($10^{37} - 10^{40} \text{ erg}$, see the papers of [4]). The equivalent radiated energy in the isotropic models is $(3 \cdot 10^3)^2$ times greater, i.e. about $10^{44} - 10^{47} \text{ erg}$. The latter energy can be even greater for closer binary systems (where the distance between a supernova precursor and the star-companion is less than $10^{13} \text{ cm}$), and it may reach values of $10^{50} - 10^{53} \text{ erg}$. The proposed model explains the main characteristics of observed GRBs.
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