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LETTER

Letter to the Editor

A jet-disk symbiosis model for Gamma Ray Bursts: SS 433 the next?

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Abstract. We consider a jet-disk symbiosis model to explain Gamma Ray Bursts and their afterglows. It is proposed that GRBs are created inside a pre-existing jet from a neutron star which collapses to a black hole due to massive accretion. In our model we assume that the initial energy due to this transition is all deposited in the jet by magnetic fields, using fully well explored concepts from jets and disks in active galactic nuclei and compact active stars in binary systems. The observed emission is then due to an ultrarelativistic shock wave propagating along the jet. We show that a good agreement between model predictions and observational data can be obtained for systems with accretion rates as high as in the Galactic jet source SS433. Specifically, we are able to reproduce the typical observed afterglow emission flux and its spectrum as a function of time.

Key words: gamma rays: bursts – shock waves – ISM: jets and outflows – radiation mechanisms: non-thermal

1. Introduction

Gamma Ray Bursts (GRBs) have been a mystery for almost 30 years. Recently, thanks to the Italian-Dutch satellite BeppoSax, it has been possible to detect for the first time their X-ray, optical, and radio afterglows. Many articles have been published to report the main characteristics of GRBs, (e.g., Guarnieri & al. 1997; Piro & al. 1998; Frail & al. 1997; Metzger & al. 1997; Gorosabel & al. 1998; Kulkarni & al. 1998). By now we know that GRBs are isotropically distributed over the sky, at least three of them are cosmological, they show a time variability in the γ-emission of the order of milliseconds, and long complex bursts. Many, but not all of them have an X-ray afterglow, and it seems that for most of them a host galaxy can be found.

Modern theoretical attempts to interpret the data are based on ideas by Mészáros and Rees (1993), Mészáros (1994), Panaitescu and Mészáros (1998), as well as Paczyński (1986), and Paczyński and Rhoads (1993). In these models a relativistic shock is caused by a relativistic fireball in a pre-existing gas, such as the interstellar medium or a stellar wind, producing and accelerating electrons/positrons to very high energies, which produce the gamma-emission and the various afterglows observed. The low level of associated radiation at other wavelengths limits the baryonic load of the emitting regions to very low amounts, and constrains the scale of the emitting region to lengths much larger than a neutron star.

One serious question is whether the overall energetics of the fireball—assumed to be isotropic—are reasonable (Sari and Piran 1997; Dar 1997) or actually exceed the level given by any conceivable model of neutron star mergers or other stellar collapses. Another question is whether a fireball model that uses external or internal shock waves can solve the baryonic mass load problem to explain the γ-emission, starting with an initial energy of 10^{51} erg in the spherical shell rest frame.

Many authors suppose the validity of the fireball model and provide evidence for its agreement with the observations (Waxman 1997a, Waxman 1997b, Vietri 1997), but the observational results for GRB971214, requiring an initial energy $E \simeq 10^{53}$ erg (if the emission is isotropic), pose a serious challenge to existing models.

Here we propose an anisotropic model where the γ-ray emission and the afterglow is produced inside a preexisting jet and calculate the temporal evolution of the corresponding flux.

2. GRB jet model

In our model GRBs develop in a pre-existing jet. We consider a binary system formed by a neutron star and an O/B/WR companion in which the energy of the GRB is due to the accretion-induced collapse of the neutron star to a black hole.

The high speed in the energy flow in AGN-jets is generally believed to be initiated by strong magnetic field coupling to the accretion disk of the black hole (e.g. Falcke & Biermann 1995, Romanova & Lovelace 1997, Romanova et al. 1997). We use the same mechanism here. We assume that in this transition a large amount of energy is anisotropically released as Poynting flux along the polar axis. One may think of this process as a violent and rapid twist, such as occasionally debated as a possible cause for supernova explosions (Kardashev 1970, Bisnovatyi-Kogan 1970, Le Blanc & Wilson 1970, and more recently Biermann 1993). This energy release naturally initiates an ultrarelativistic shock wave in the pre-existing jet. The emission microphysics are as in existing fireball models.

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To fix the jet parameters, we use the basic ideas of the jet-disk symbiosis model by Falcke & Biermann (1995, 1999). In this model the direction of the magnetic field is mostly perpendicular to the axis of the jet and the values of the total particle number $n_{\rm j}$ (relativistic electrons + thermal protons) and magnetic field $B_{\rm j}$ are calculated using the equipartition between magnetic field energy density and kinetic plasma energy density in the umperturbed jet. Both, $B_{\rm j}$ and $n_{\rm j}$, are a function of the jet ejection rate $\dot M_{\rm jet} \simeq 0.05\dot M_{\rm disk} \simeq (10^{-5}M_\odot/{\rm yr})\dot M_{-5{\rm j}}$, the equipartition parameter ϵ and the bulk velocity of the jet $\beta_{\rm j}\gamma_{\rm j} = 0.3v_{0.3}$, where $\gamma_{\rm j}$ is the bulk Lorentz factor of the jet prior to the burst. The values of the particle number density and the magnetic field in the unperturbed jet are then given as

$$n_{\rm j}(z_{\rm j}) \simeq 80.34\ (\dot M_{-5{\rm j}}v_{0.3}^{-1})\theta_{-1{\rm j}}^{-2}z_{17,{\rm j}}^{-2} \quad {\rm cm}^{-3} \tag{1}$$

$$B_{\rm j}(z_{\rm j}) \ = \ 1.05\ (\dot M_{-5{\rm j}}^{1/2}v_{0.3}^{-1/2})(\epsilon^{-1/2}\theta_{-1{\rm j}}^{-1}\gamma_{\rm m,2}^{1/2}) \times z_{17,{\rm j}}^{-1} \quad {\rm Gauss} \tag{2}$$

where $\gamma_{\rm m,2}$ is the minimum electron Lorentz factor $\gamma_{\rm em} = 100\gamma_{\rm m,2}$ of the relativistic electrons (assumed to be in a power law distribution), $\theta_{-1{\rm j}} = \theta_{\rm j}/(0.1{\rm rad})$ is the opening angle of the jet, and $z_{\rm j} = 10^{17}z_{17,{\rm j}}\,{\rm cm}$ is the distance along the jet. z is the redshift.

The advantages of using the propagation of an ultrarelativistic shock wave inside a jet can be summarized in the following two points

- The initiation of the shock by magnetic fields does not necessarily involve considerable matter. It implies a low amount of baryonic matter, of order $5 \times 10^{-8}\ M_\odot$, in the jet.
- The initial amount of energy deposited in the jet is of the order $E_{51} = E/(10^{51}{\rm erg})$, and we can obtain an apparent isotropic energy as high as $E_{\rm app} = 10^{53.3}E_{51}\theta_{-1{\rm j}}^{-2}$ erg. This fits well the requirements of GRB971214.

The evolution of the shock Lorentz factor $\gamma_{\rm sh}$ with distance $z_{\rm j}$ along the jet axis can be obtained using the conservation of the unperturbed jet gas energy in the shock rest frame, $E/2 = \pi{\theta_{\rm j}}^2 z_{\rm j}^3\gamma_{\rm sh}^2 m_{\rm p}c^2 n_{\rm j}$, that is

$$\gamma_{\rm sh}(z_{\rm j}) \approx 11.48\ (E_{51}^{1/2}\dot M_{-5{\rm j}}^{-1/2}v_{0.3}^{1/2})z_{17,{\rm j}}^{-1/2}. \tag{3}$$

The characteristic time scale to see the emission across the entire region when the shock has reached a distance $z_{\rm j}$ along the axis of the jet is given by $t^{\rm (ob)} = z_{\rm j}/(2\gamma_{\rm sh}^2 c)$. Substituting this into the formula for $\gamma_{\rm sh}$, we get the time evolution of the ultrarelativistic shock front

$$z_{\rm j}(t) \approx 2.81 \times 10^{17}(E_{51}^{1/2}\dot M_{-5{\rm j}}^{-1/2}v_{0.3}^{1/2})t_5^{1/2} \quad {\rm cm} \tag{4}$$

$$\gamma_{\rm sh}(t) \approx 6.84\ (E_{51}^{1/4}\dot M_{-5{\rm j}}^{-1/4}v_{0.3}^{1/4})t_5^{-1/4}, \tag{5}$$

where $t_5 = t/(10^5{\rm s})$.

A change in the emission properties occurs when the opening angle of our jet passes the relation $1/\gamma_{\rm sh} = \theta_{\rm j}$. Using the formula of the bulk Lorentz factor as a function of time, we can see at what time $t^\star$ this relation holds: $t^\star \approx 2.20 \times 10^4(E_{51}\dot M_{-5{\rm j}}v_{0.3})\theta_{-1{\rm j}}^4$ s. Hence, after about 6 hours this limit is reached. Prior to this, the observed emission is limited by the Lorentz boost to a conical section of the shock front of angle $1/\gamma_{\rm sh}$. For the following calculation of the afterglow emission several hours after the burst we can therefore consider $1/\gamma_{\rm sh} > \theta_{\rm j}$.

In the flow behind a steady shock front, relativistic particles are usually accelerated and magnetic fields can be amplified. After the shock, $B'_\parallel \simeq B_\parallel$ and $B'_\perp \simeq 4\gamma_{\rm sh}B_\perp$, relative to the jet axis. The resulting magnetic field in the shocked plasma will have a strength of roughly $B \approx (16\gamma_{\rm sh}^2+1)^{1/2}B_\perp$ and the shock wave compresses the magnetic field component perpendicular to the jet axis. The jump conditions for the density particle number give $n^{\rm (sf)} \simeq 4\gamma_{\rm sh}n_{\rm j}$, (De Hoffmann and Teller 1950, Marscher and Gear 1985). We use here the approximation for an ultrarelativistic shock so that $n_2/n_1 = 4\gamma_{12} \simeq 4\gamma_{\rm sh}$, where 1 and 2 are related to the zone before and after the shock. This allows a straight and simple limit to nonrelativistic shocks. In our model, this corresponds to

$$n_{\rm j}^{'\rm (sf)}(t) \ = \ 2.78 \times 10^2(E_{51}^{-3/4}\dot M_{-5{\rm j}}^{7/4}v_{0.3}^{-7/4}) \times \theta_{-1{\rm j}}^{-2}t_5^{-5/4} \quad {\rm cm}^{-3} \tag{6}$$

$$B_{\rm j}^{'\rm (sf)}(t) \ = \ 10.24\ (E_{51}^{-1/4}\dot M_{-5{\rm j}}^{3/4}v_{0.3}^{-3/4}) \times (\epsilon^{-1/2}\theta_{-1{\rm j}}^{-1}\gamma_{\rm m,2}^{1/2})t_5^{-3/4} \quad {\rm Gauss} \tag{7}$$

where (sf) shows the quantities in the shock frame and (ob) the corresponding values in the observer frame.

In our jet we consider an electron power law distribution $N(\gamma_{\rm e})d\gamma_{\rm e} = C_{\rm e}\gamma_{\rm e}^{-{\rm p}}d\gamma_{\rm e}$ where $\gamma_{\rm em} < \gamma_{\rm e} < \gamma_{\rm max}$, with a cutoff at a constant minimum electron Lorentz factor $\gamma_{\rm em} \simeq 100$. In fact, if p-p collisions inject a population of electrons in the unperturbed jet, then one would expect that in the unperturbed jet the electron Lorentz factor $\gamma_{\rm e}$ goes from $\gamma_{\rm em} \simeq 100 \simeq m_\pi/m_{\rm e}$ up to some large value (Falcke & Biermann 1995).

The equation for $C_{\rm e}$ has been obtained considering that in the shock frame the powerlaw electron energy density is taken as a fraction $\delta \leq 1$ of the pre-existing relativistic electron energy density and using the value p=2

$$C_{\rm e}^{\rm (sf)}(t) \ = \ 1.91 \times 10^4\delta(E_{51}^{-1/2}\dot M_{-5{\rm j}}^{3/2}v_{0.3}^{-3/2}) \times (\theta_{-1{\rm j}}^{-2}\gamma_{\rm m,2})t_5^{-3/2} \quad {\rm cm}^{-3} \tag{8}$$

If we were to relate to the proton energy density in the shock, $\delta \leq 20 \simeq m_{\rm p}/(m_{\rm e}\gamma_{\rm em})$.

To calculate the boosting factor of the transition from the shock frame to observer frame, we assume the angle between the jet-axis and the line of sight to the observer is $\theta_{\rm obs} < 1/\gamma_{\rm sh}$ some time during the γ-ray burst. The critical synchrotron frequency after the shock front in the observer frame is given by $\nu_{\rm m} \approx \frac{2}{1+z}\gamma_{\rm sh}^3\gamma_{\rm em}^2\frac{eB}{m_{\rm e}c}$, where we consider the minimum electron Lorentz factor evolves in phase space with the bulk Lorentz factor $\gamma_{\rm em} \times \gamma_{\rm sh}$:

$$\nu_{\rm m}(t) \approx \frac{2.61 \times 10^{14}}{1+z}E_{51}^{1/2}(\epsilon^{-1/2}\theta_{-1{\rm j}}^{-1}\gamma_{\rm m,2}^{5/2})t_5^{-3/2} \quad {\rm Hz} \tag{9}$$

We assume isotropic emission in the shock rest frame. Only when we transform the radiation emitted into the observer frame, the photons are concentrated in the forward direction, lying within a cone of half-angle $1/\gamma_{\rm sh}$. For the afterglow, however, the solid angle is determined by the actual opening angle of the jet.

At the critical frequency, $\nu_{\rm m}$, in our model the synchrotron cooling time is less than the dynamical time, $z_{\rm j}/(\gamma_{\rm sh}c)$. In fact, in the shock frame

$$t_{\rm s}^{\rm (sf)}/t_{\rm d}^{\rm (sf)} \approx 7.88\times10^{-3}(\dot{M}_{-5\rm j}^{-1}v_{0.3})(\epsilon\theta_{-1\rm j}^{2}\gamma_{\rm m,2}^{-2})t_5 \qquad (10)$$

so the synchrotron time scale is shorter than the dynamical time scale for our standard parameters.

Considering the cooling evolution of our electron power law spectrum, to calculate the flux we use the formula $F_\nu^{\rm (ob)}(t) = \left(\frac{dP}{d\nu}\right)_{\rm p=2}^{\rm (ob)} z^2 x_\nu \pi\theta_{\rm j}^2\frac{1}{4\pi D^2}$, where $\left(\frac{dP}{d\nu}\right)_{\rm p=2}^{\rm (ob)}$ is the total emission power per unit volume per unit frequency for an electron power law distribution with an index $p=2$, (e.g. Sect. 6.4 of Rybicki & Lightman, 1979). In our model this corresponds to a flux in the observer frame given by

$$F_\nu^{\rm (ob)}(t) = 4.00\times10^{-18}C_{\rm e}^{\rm (sf)}(t)B^{3/2\rm (sf)}(t)\nu^{-1/2}z^2(t)\times x_\nu\theta_{\rm j}^2\gamma_{\rm sh}^3\times\frac{1}{4D^2}. \qquad (11)$$

Here, the quantity x_ν represents the thickness of the radiating shock front, extending behind the shock front to a point where the local $\nu_{\rm m}$ drops below ν

$$x_\nu = 4.70\times10^{13}(E_{51}^{3/8}\dot{M}_{-5\rm j}^{-9/8}v_{0.3}^{9/8})\times(\epsilon^{3/4}\theta_{-1\rm j}^{3/2}\gamma_{\rm m,2}^{-3/4})t_5^{9/8}\nu_{14}^{-1/2}\quad {\rm cm} \qquad (12)$$

where $\nu_{14}=\nu/(10^{14}{\rm Hz})$. In both equations (11) and (12) ν is in the shock frame.

Substituting the values of the equations (4), (5), (7), (8) and (12) in (11), and with ν in the observer frame

$$F_\nu^{\rm (ob)}(t) = 7.45\times10^{-28}\delta(E_{51}^{5/4}\dot{M}_{-5\rm j}^{-1/4}v_{0.3}^{1/4})\times\gamma_{\rm m,2}D_{28.5}^{-2}t_5^{-5/4}\nu_{14}^{-1}\quad {\rm erg\ cm^{-2}s^{-1}Hz^{-1}} \qquad (13)$$

where D is the luminosity distance, $D_{28.5}=D/(10^{28.5}{\rm cm})$ corresponds to a redshift of about $z\simeq1.5$, using $q_{\rm o}=1/2$ and $H_{\rm o}=65\ {\rm km\ s^{-1}Mpc^{-1}}$.

The flux at frequencies ν below the critical frequency, $\nu_{\rm m}$, can be obtained using the flux $F_{\nu_{\rm m}}(t)$ at frequencies $\nu>\nu_{\rm m}$, calculated at the frequency $\nu_{\rm m}$, that is $F_{\nu<\nu_{\rm m}} = F_{\nu_{\rm m}}(t)\times\left(\frac{\nu}{\nu_{\rm m}}\right)^\beta$. The corresponding flux is given by

$$F_{\nu<\nu_{\rm m}}^{\rm (ob)}(t) = (2.84\times10^{-28})(2.61)^{-\beta}\delta(E_{51}^{3/4-\beta/2}\dot{M}_{-5\rm j}^{-1/4}\times v_{0.3}^{1/4})(\epsilon^{1/2+\beta/2}\theta_{-1\rm j}^{1+\beta}\gamma_{\rm m,2}^{-3/2-5\beta/2})D_{28.5}^{-2}\times t_5^{1/4+3\beta/2}\nu_{14}^{\beta}(1+z)^{1+\beta}\ {\rm erg\ cm^{-2}s^{-1}Hz^{-1}} \qquad (14)$$

For optically thin emission, β is equal 1/3 and the time dependence is $t^{3/4}$. This suggests a gentle optical rise.

Another important check is to calculate the synchrotron self-absorption coefficient α_ν in our model (Sect. 6.8 of Rybicki & Lightman, 1979). Integrating it over an isotropic distribution of particles and using the value obtained for the coefficient $C_{\rm e}$, we find $\alpha_\nu(t)$. The optical depth along the path of a travelling ray is equal to the product of the absorption coefficient $\alpha_\nu(t)$ times the width of the shell (equation 12), and it will be equal to unity at the time

$$t_{\tau=1} = 0.12\times\delta^{8/15}(E_{51}^{-1/3}\dot{M}_{-5\rm j}v_{0.3}^{-1})(\epsilon^{-2/15}\theta_{-1\rm j}^{-4/3}\times\gamma_{\rm m,2}^{2/3})\nu_{14}^{-28/15}\quad {\rm s} \qquad (15)$$

3. Discussion and conclusion

We now check our model against the observed data. Here we assume a fixed opening jet angle and a minimum electron Lorentz factor, and the equipartition parameters ϵ and δ are referenced to unity.

For GRB970208, using the equation (9), at a time of $t_{\rm opt} = t/(7.6\times10^4 s)$, we obtain a frequency $\nu_{\rm m}(t_{\rm opt}) \approx (3.95\times10^{14})(1+z)^{-1}E_{51}^{1/2}(\epsilon^{-1/2}\theta_{-1\rm j}^{-1}\gamma_{\rm m,2}^{5/2})t_{\rm opt}^{-3/2}$ Hz.

Substituting the same numbers into the equations for the flux (13), we have $F_\nu(t_{\rm opt}) \approx 3.09\times10^{-28}\ \delta(E_{51}^{5/4}\dot{M}_{-5\rm j}^{-1/4} v_{0.3}^{1/4})\gamma_{\rm m,2}\,D_{28.5}^{-2}\,\nu_{\rm opt}^{-1}\,t_{\rm opt}^{-5/4}\ {\rm erg\ cm^{-2}s^{-1}\,Hz^{-1}}$, where $\nu_{\rm opt} = \nu/(3.4\times10^{14}{\rm Hz})$.

Again using GRB970208 as reference, we obtain for the flux in the X-band at a frequency of 10^{18} Hz and a time 3.6×10^4 s: $F_\nu(t_{\rm X}) \approx 2.67\times10^{-31}\delta(E_{51}^{5/4}\ \dot{M}_{-5\rm j}^{-1/4}v_{0.3}^{1/4})\gamma_{\rm m,2}D_{28.5}^{-2}\nu_{\rm X}^{-1}t_{\rm X}^{-5/4}\ {\rm erg\ cm^{-2}s^{-1}Hz^{-1}}$.

For the optical depth in the optical band we obtain $\tau_\nu(t_{\rm opt}) \approx 2.00\times10^{-13}\delta(E_{51}^{-5/8}\dot{M}_{-5\rm j}^{15/8}\ v_{0.3}^{-15/8})(\epsilon^{-1/4}\theta_{-1\rm j}^{-5/2}\gamma_{\rm m,2}^{5/4})\ \nu_{\rm opt}^{-7/2}t_{\rm opt}^{-15/8}$, hence, within our model the optical emission region is optically thin, except very early.

In our model, the condition $1/\gamma_{\rm sh}<\theta_{\rm j}$ corresponds to the early phase, when $\nu_{\rm m}$ is actually in the gamma-ray regime. In this phase, the corresponding maximum flux in the observer frame is $F_{\rm max} \approx 8.66\times10^{-7}\delta E_{51}D_{28.5}^{-2}\theta_{-1\rm j}^{-2}\gamma_{\rm m,2}t^{-1}\ {\rm erg\ cm^{-2}s^{-1}}$. This is derived from using all available energy in energetic electrons, and redistributing it into emission; photons are generated by various emission processes, including pion decay, upscattered by inverse Compton emission, and redistributed in photon energy by pair opacity (Rachen & Mészáros 1998). The integration of this flux in time from 10^{-3} s to 10 s gives a value of about $10^{-5}\delta E_{51}D_{28.5}^{-2}\theta_{-1\rm j}^{-2}\gamma_{\rm m,2}\ {\rm erg\ cm^{-2}}$. This we identify with the initial gamma-ray emission.

Concerning the variability at times when $1/\gamma_{\rm sh}>\theta_{\rm j}$, we note that in observed jets, one often finds inhomogeneities on the scale of a few jet-diameters parallel to the axis, and down to some fraction of the diameter perpendicular to the axis. However, any variability derived from such inhomogeneities is smeared out by arrival time differences for the observer. Therefore, using the jet diameter as a reference scale, this smearing limits any temporal variability

$$\frac{\Delta t}{t}\gtrsim(\gamma_{\rm sh}\theta_{\rm j})^2 \simeq 0.47(E_{51}^{1/2}\dot{M}_{-5\rm j}^{-1/2}v_{0.3}^{1/2})t_5^{-1/2} \qquad (16)$$

Inhomogeneities on transverse scales smaller than the jet diameter can shorten this. This variability may explain the complex features of the optical rise in GRB970508 (Galama et al. 1998).

The way in which our model can reproduce the observed properties of GRBs depends strictly on two sets of parameters, one set which is verifiable, because it derives from known active binary stars, and a second set which characterizes the explosion. The explosion depends on the initial shock energy E_{51}, and the equipartition parameter δ. On the other hand, the binary set is composed of the mass loss rate $\dot{M}_{-5\rm j}$ of the jet, the jet opening angle $\theta_{\rm j}$, the bulk velocity in the jet $v_{0.3}$ and the equipartition parameter ϵ. To produce shock waves and the observed emission in the jet, it is necessary that the initial amount of matter in the shock is comparable with the mass in the jet. This implies that the binary system necessary in our model has to have a super Eddington accretion rate to get the required mass loss rate. In our Galaxy, one such binary system is known, SS433 (Murata and Shibazaki, 1996), for which $\dot{M}_{\rm disk} \simeq (2 \times 10^{-4} M_{\odot}/{\rm yr})$. In our model, we have used the parameters characteristic of this system, implying that there are system like SS433 where the central object indeed is a neutron star.

In the context of the Falcke $\&$ Biermann jet-disk model, SS433 is a radio-weak jet-disk system, while here we have used their radio-loud model; but it has been noted that systems such as GRS1915+105 and perhaps also SS433 can switch from radio-weak to radio-loud in the terminology of Falcke & Biermann (1995, 1999). Considering that the jet ejection rate is bound to the disk accretion rate by the relation $\dot{M}_{\rm jet} \simeq 0.05 \dot{M}_{\rm disk}$, we find that for an initial bulk Lorentz factor $\Gamma_{\rm o} = 10^4$, the mass in the shock $M_{\rm sh} \approx 5 \times 10^{-8} M_{\odot}$ is equal to the mass present in the jet. Therefore, SS433 might be a good candidate to explode as the next violent GRB in our Galaxy.

To summarize, our model can explain the initial gamma ray burst, the spectrum and temporal behaviour of the afterglows, the low baryon load, an optical rise, and do all this with a modest energy budget. Moreover, this GRB model is developed within an existing framework for galactic jet sources, using a set of well determined parameters.

Obviously, we have simplified in many places, using a naive version of shock acceleration, only strong shocks, a conical jet geometry, etc., but the good agreement with the data we obtain within our framework shows that more detailed calculations may be worthwhile.

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