

LETTER TO THE EDITOR

Flares in GRB afterglows from delayed magnetic dissipation

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ABSTRACT

One of the most intriguing discoveries made by the *Swift* satellite is the flaring activity in about half of the afterglow lightcurves. Flares have been observed on both long and short duration GRBs and on time scales that range from minutes to ~ 1 day after the prompt emission. The rapid evolution of some flares led to the suggestion that they are caused by late central engine activity. Here, I propose an alternative explanation that does not need reviving of the central engine. Flares can be powered by delayed magnetic dissipation in strongly magnetized (i.e. with initial Poynting to kinetic flux ratio $\gtrsim 1$) ejecta during its deceleration due to interaction with the external medium. A closer look at the length scales of the dissipation regions shows that magnetic dissipation can give rise to fast evolving and energetic flares. Multiple flares are also expected in the context of the model.

Key words. gamma rays: bursts – magnetohydrodynamics (MHD) – instabilities

1. Introduction

A rich early afterglow phenomenology has recently been revealed thanks, to a large extent, to the X-ray detector on board to the *Swift* satellite. The systematic study of X-ray afterglows (Zhang et al. 2006; Nousek et al. 2006) has shown that they are typically characterized by an initial steep decay, associated with the end of the prompt emission phase, followed by a flatter decay that lasts for about one hour and probably corresponds to the onset of the afterglow. At later times, the lightcurves are characterized by the “normal” decay observed before *Swift*.

Another exciting *Swift* discovery is the afterglow flares observed in about half of the bursts (Burrows et al. 2005). Flares appear on different timescales t_f after the prompt emission ranging from ~ 100 to $\sim 10^5$ s, they are characterized by duration δt_f such that $\delta t_f/t_f \lesssim 1$, and show spectral evolution with respect to the smooth decay part of the afterglow. The flares are involving substantial energy release that ranges from a small fraction up to an amount equal to that of the prompt GRB emission (as in the case of GRB 050502b; Falcone et al. 2006) and they have been observed in the afterglow lightcurves of both long and short duration GRBs (Campana et al. 2006).

Although bumps in the afterglow lightcurves can be a result of slower shells catching up with the decelerating ejecta (e.g. Rees & Mészáros 1998) or interaction of the ejecta with a clumpy external medium, the rapid evolution of some of the flares has been used as an argument against their external shock origin. This led to the suggestion that the central engine revives at later times (ranging from hundreds of seconds to days) giving rise to flares through late internal shocks or magnetic dissipation (Burrows et al. 2005; Zhang et al. 2006). Suggestions on how to revive the central engine have been made by King et al. (2005), Dai et al. (2006), Fan et al. (2006), Perna et al. (2006).

Here, I make an alternative suggestion which does not need late time activity of the central engine. I investigate the possibility that the deceleration of strongly magnetized ejecta leads to revival of MHD instabilities, that were delayed by time dilation

before the deceleration phase. The resulting dissipation produces flares in localized reconnection regions in the flow.

2. Magnetic dissipation in a decelerating shell

Magnetic fields may be the main agent extracting energy from the central engine, leading to a Poynting-flux-dominated flow (e.g., Thompson 1994; Mészáros & Rees 1997; Spruit et al. 2001). Magnetic dissipation through instabilities in an axisymmetric flow (Lyutikov & Blandford 2003; Giannios & Spruit 2006) or directly through reconnection in a highly non-axisymmetric flow (Drenkhahn & Spruit 2002) can power the prompt emission with high radiative efficiencies and accelerate the flow to ultrarelativistic speeds. The photospheric emission of a dissipative MHD flow is highly non-thermal and can be responsible for the prompt GRB emission (Thompson 1994; Giannios 2006).

Dissipation of Poynting flux is partial, however, if it depends on global magnetic instabilities. MHD instabilities (such as the kink instability) grow on the Alfvén crossing time in the fluid frame. Time dilation in the central engine frame by the bulk Lorentz factor of the flow Γ results in slowing down of the dissipation in ultrarelativistic flows. Giannios & Spruit (2006) have shown that for large initial Poynting to kinetic flux ratios, the flow remains magnetically dominated up to radii $\gtrsim 10^{16}$ cm, i.e. close to the distance where the deceleration of the flow against the external medium is expected.

The deceleration of the flow can, however, naturally lead to reviving of the instabilities in the regions which come into causal contact again and to substantial magnetic dissipation in the afterglow phase. Another (and probably complementary) possibility is that dissipation of magnetic energy is triggered by the crossing of the reverse shock through the magnetized shell (Thompson 2005).

Here, I investigate the possibility that late time dissipation takes place in this way and explore whether it can power the commonly observed flares in the GRB afterglows. I discuss in

particular the energetics and the variability expected from such events.

2.1. Deceleration of a magnetized shell

The issue of flow acceleration and of prompt emission in a Poynting flux dominated flow has been studied in a number of works in the context of GRB flows (e.g., Drenkhahn & Spruit 2002; Vlahakis & Königl 2003; Giannios & Spruit 2006). At a sufficient distance from the source, the flow has accelerated to its terminal Lorentz factor and the magnetic field is dominated by its toroidal component. After the ultrarelativistic flow has swept up enough mass from the external medium (which may be the interstellar medium or the wind material), it starts decelerating, driving a shock into the external medium. In this work, I focus on this deceleration phase.

I assume that a cold, strongly magnetized shell with a dominant toroidal magnetic field is moving with initial bulk Lorentz factor $\Gamma_0 \gg 1$ into the external medium. The magnetic content of the ejecta can be conveniently parameterized by the magnetization σ defined as the ratio of Poynting to kinetic energy flux $\sigma = B'^2/4\pi w$, where B' and $w = \rho c^2 + e + p$ are the magnetic field and the enthalpy density of the ejecta respectively as measured by an observer comoving with the flow. The enthalpy density consists of the rest mass energy density ρc^2 , the energy density e and the gas pressure p . Since the shell is assumed to be cold initially, its magnetization is $\sigma_0 = B'^2/4\pi\rho c^2$.

The initial phase of the interaction of the magnetized shell with the external medium has been studied by Zhang & Kobayashi (2005) who have shown that, for “typical” GRB parameters, a reverse shock forms in the ejecta as long as $\sigma_0 \lesssim 100$. In the shocked part of the ejecta, both gas and magnetic field contribute to the total pressure; their sum balances the pressure of the shocked external medium at the contact discontinuity. For high enough initial magnetization of the ejecta $\sigma_0 \gtrsim 0.1$, the magnetic pressure dominates the gas pressure in the shocked ejecta irrespective of whether the reverse shock is Newtonian, mildly relativistic or relativistic. This can be shown by solving for the MHD shock conditions (Kennel & Coroniti 1984; Zhang & Kobayashi 2005), where it is also shown that shock compression leads to an increase of the magnetization of the ejecta, i.e. $\sigma_{\text{sh}} \gtrsim \sigma_0$.

Thus, for $\sigma_0 \gtrsim 0.1$, one can neglect the gas pressure of the shocked ejecta with respect to the magnetic pressure. The pressure balance at the contact discontinuity yields

$$\frac{B_{\text{sh}}'^2}{8\pi} = \frac{4}{3}\Gamma_{\text{sh}}^2\rho_e c^2, \quad (1)$$

where B_{sh}' , Γ_{sh} are magnetic field strength and the bulk Lorentz factor of the shocked ejecta and ρ_e is the density of the external medium. From this point on, I focus on the shocked ejecta and omit the subscript “sh” for simplicity in the notation.

2.2. The magnetic dissipation region

I have argued that the deceleration of the flow (and/or the crossing of the reverse shock) revives MHD instabilities that lead to dissipation of magnetic energy through reconnection of magnetic field lines at different locations in the flow. The energy released in a dissipative event and its observed duration depend on the characteristic dimensions of the dissipation region. Let l'_1 and l'_2 be the characteristic length scales in the radial and perpendicular to the bulk motion directions (measured by a comoving

observer) of this region respectively. Furthermore, I consider l'_1 to be the scale over which the field changes polarity and that magnetic field lines are advected in the middle plane of this region where they reconnect. The reconnection plane is assumed to be perpendicular to the radial direction. More generally, one should consider an arbitrary angle between the radial direction and the normal to the reconnection plane. However, I have checked that such a generalization only complicates the analysis that follows, without introducing any essential changes on the results, and has been avoided.

The electromagnetic energy that is contained in the dissipation region is (in the central engine frame)

$$E_d = \Gamma l'_1 l'_2 \frac{B'^2}{4\pi} = \frac{8}{3} l'_1 l'_2 \Gamma^3 \rho_e c^2, \quad (2)$$

where Eq. (1) has been used in the last step. If a large fraction of the dissipated energy leads to fast moving electrons, they can efficiently radiate most of it *if* the electrons’ cooling timescale is shorter than the expansion time scale $t_{\text{exp}} = r/\Gamma c$ of the flow. Here, I show that this is the case in the dissipation region under consideration.

An efficient cooling mechanism in the strongly magnetized flow is synchrotron cooling with a time scale

$$t_{\text{syn}} = 3m_e c / 4\sigma_T \beta_e^2 \gamma_e U_B, \quad (3)$$

where γ_e is the characteristic Lorentz factor at which electrons are accelerated in the reconnection region and $U_B = B'^2/8\pi$ is the magnetic energy density. The available magnetic energy per proton (and electron) is $\gtrsim m_p c^2$ for a flow with $\sigma_0 \gtrsim 1$, which is the range of σ_0 that I focus on in this study. Dissipation of most of this energy to the electrons leads to $\gamma_e \gtrsim m_p/m_e$. By using, for an order of magnitude estimate, $\Gamma = 50$, $r = 3 \times 10^{17}$ cm, $\rho_e = 10^{-24}$ gr cm $^{-3}$, and $\gamma_e = 10^{3.5}$ in Eqs. (1) and (3), I find $t_{\text{syn}}/t_{\text{exp}} \sim 10^{-2}$ and the characteristic synchrotron photon energy $E_{\text{syn}} = 3\Gamma\gamma_e^2 q_e \hbar B' / 2m_e c \sim 100$ eV, i.e. in the soft X-rays. Smaller values for $t_{\text{syn}}/t_{\text{exp}}$ and higher for E_{syn} are found for “typical” parameters that correspond to the case where the external medium is the stellar wind.

Thus, it is possible that a large fraction of the reconnected energy is promptly radiated away in the X-rays through synchrotron emission. This energy is emitted in the $\theta \sim 1/\Gamma$ forward cone, i.e., with a beaming factor $4\pi/2\pi(1 - \cos\theta) \simeq 4\Gamma^2$. The beaming leads to isotropic equivalent energy $E_{\text{t}}^{\text{iso}}$ that is $4\Gamma^2$ times larger than the dissipated energy E_d

$$E_{\text{t}}^{\text{iso}} = 4\Gamma^2 E_d. \quad (4)$$

If the released magnetic energy is to be observed as a flare, the dissipation must take place sufficiently fast. This leads to conditions on the length scales l'_1 and l'_2 that have to hold. These conditions are investigated in the next section.

2.3. Variability constraints

In the dissipation region, magnetic fields are advected along a distance $\sim l'_1$ into the reconnection plane which has a characteristic length l'_2 . The observed duration of the flare δt_f is the sum of two contributions. The first is the “radial” duration δt_r coming from the fact that while dissipation takes place, the flow expands radially resulting at different arrival times for photons that are emitted at different stages of the dissipation event. The second contribution comes from the “angular” duration δt_{ang} that is related to the delay in arrival of photons that are emitted in different locations on the reconnection plane. The duration of the flare

is $\delta t_f = \delta t_r + \delta t_{\text{ang}}$. I turn to these two contributions and study them separately.

The “radial” duration of the flare depends on how fast reconnection takes place. The speed v_r at which magnetic reconnection proceeds in a strongly magnetized plasma is enhanced by the relativistic kinematics of the problem (e.g., Lyutikov & Uzdensky 2003; Lyubarsky 2005) and is a substantial fraction $\varepsilon \sim 0.1$ of the Alfvén speed which is close to the speed of light for $\sigma_0 \gtrsim 1$. Magnetic field lines are advected for a distance $l'_1/2$ above and below the reconnection plane with speed v_r , which leads to a co-moving dissipation timescale $t'_d \simeq l'_1/(2v_r) \simeq l'_1/(2\varepsilon c)$. Dividing this timescale with the expansion timescale $r/\Gamma c$, one has an estimate on the fractional duration $\delta t_r/t_f$ of the flare because of the radial motion of the reconnecting region

$$\delta t_r/t_f = \Gamma l'_1/2\varepsilon r. \quad (5)$$

The “angular” duration is related to the characteristic length l'_2 of the reconnecting region. A first constraint to l'_2 comes from causality arguments that yield $l'_2 \lesssim r/\Gamma$. The observer sees an emitting region with length l'_2 from a characteristic angle $\sim 1/\Gamma$. The difference in the arrival time of two photons coming from the two edges of the emitting region is $\delta t_{\text{ang}} \simeq l'_2/c\Gamma$. How the observer time is related to the radius and bulk Lorentz factor of the blast wave depends on the details of the deceleration profile which, in turn, depend on the density profile of the external medium and on the magnetic content of the ejecta. To keep the study general, I assume that $t_f = r/\alpha c\Gamma^2$, where α is a parameter with typical values \sim a few (e.g., for baryonic ejecta decelerated in constant density medium $\alpha = 4$, Waxman 1997, and in stellar wind $\alpha = 2$, Pe'er & Wijers 2005). Using the expressions for δt_{ang} and t_f , I have

$$\delta t_{\text{ang}}/t_f = \alpha\Gamma l'_2/r. \quad (6)$$

The observed duration of the flare δt_f is given by the sum of the radial and angular timescales. Using Eqs. (5) and (6), I arrive at

$$\delta t_f/t_f = \Gamma l'_1/2\varepsilon r + \alpha\Gamma l'_2/r. \quad (7)$$

The observed fractional duration of the flare thus constrains both characteristic length scales of the reconnection region

$$l'_1 \lesssim 2\varepsilon r(\delta t_f/t_f)/\Gamma, \quad (8)$$

$$l'_2 \lesssim r(\delta t_f/t_f)/\alpha\Gamma. \quad (9)$$

In the process of deriving the angular time scale, I have implicitly ignored fast fluid motions in the frame moving with the bulk of the flow. The material leaves the reconnection region with the Alfvén speed (e.g. Lyubarsky 2005) however, which is close to the speed of light. This can affect the estimate (9) by allowing for larger values of l'_2 (and ultimately for larger energy release in a flare of a specific $\delta t_f/t_f$; see the rest of the section). This effect is particularly pronounced in $\sigma_0 \gg 1$ flows which can contain emitting regions with high “internal” Lorentz factors and lead to more rapid variability than that estimated here (Lyutikov & Blandford 2003; Lyutikov 2006).

Expressions (8) and (9), combined with Eqs. (4) and (2), yield a limit on the isotropic equivalent energy for a flare that can be produced by a single reconnection event

$$E_f^{\text{iso}} \lesssim 20\varepsilon(\delta t_f/t_f)^3\Gamma^2 r^3 \rho_e c^2/\alpha^2, \quad (10)$$

where the numerical factor of the last expression $64/3$ was set to 20. Equation (10) can be rewritten in a more compact form

by using the fact that $4\pi r^3 \rho_e/3 \simeq M_{\text{swept}}$, where M_{swept} is the mass of the external medium that has been swept up by the blast wave. Furthermore, the forward shock conditions dictate that $\Gamma^2 M_{\text{swept}} c^2$ is to be identified with the energy that has been passed onto the forward shock $E_{\text{fs}}^{\text{iso}}$. Expression (10) can, thus, be rewritten

$$E_f^{\text{iso}} \lesssim 5\varepsilon(\delta t_f/t_f)^3 E_{\text{fs}}^{\text{iso}}/\alpha^2. \quad (11)$$

From the last expression, it is clear that the model predicts that the energy available to power the fastest evolving flares with $\delta t_f/t_f \sim 0.1$ is 3 orders of magnitude less than the energy that can be released during a smoother flare with $\delta t_f/t_f \sim 1$. The energy of the flare depends on $E_{\text{fs}}^{\text{iso}}$, the radial dependence of which has not been studied for strongly magnetized ejecta. The key open question is how fast is the Poynting flux passed onto the forward shock. This has an important effect on the bulk Lorentz factor of the flow as a function of radius $\Gamma(r)$ (see Zhang & Kobayashi 2005; Lyutikov 2006) and, therefore, on the (observer) timescales over which magnetic dissipation can result in powerful flares. A detailed study of the dynamics of the deceleration of ejecta with $\sigma_0 \gtrsim 1$ can settle these issues and yield valuable constraints on the model.

Nevertheless, one can have a rough estimate of the energetics of a flare by assuming that the energy in the forward shock at some radius r is a large fraction of the energy that was initially carried by the ejecta $E_{\text{ej}}^{\text{iso}}$. In view of Eq. (11), it means that a fraction

$$f \lesssim 5\varepsilon(\delta t_f/t_f)^3/\alpha^2 \quad (12)$$

of the initial blast wave energy can power a flare of duration δt_f at (observer) time t_f . As an arithmetic example, I set $\alpha^2 \simeq 10$ and $\varepsilon \simeq 0.2$, from which I find $f \lesssim 0.1(\delta t_f/t_f)^3$. One can go a step further and compare the energy of the flare with that of the prompt GRB emission by assuming that the latter is a fraction $\eta \sim 0.1$ of the energy of the blast wave. This means that a flare coming from late magnetic dissipation can reach a fraction $\sim(\delta t_f/t_f)^3$ of the prompt GRB emission. So the model may explain powerful and fast evolving flares at the same time. On the other hand, there is a clear prediction that fast evolving flares are less energetic than the smoother ones.

2.4. Multiple flares

Very often, the afterglow light curves are characterized by multiple flares. In the context of this model, this corresponds to multiple reconnection regions. The physical scales of such region are constrained by Eqs. (8) and (9) and one can check that they correspond to a moderate fraction of the volume of a shell with $1/\Gamma$ opening angle that emits toward the observer. So, a single flare does not use up all the Poynting flux in the observer’s light cone and repetition is possible. It is even likely that dissipation in one region can induce instabilities and magnetic dissipation in neighboring regions, which results in a sequence of flares.

3. Conclusions

The afterglow flares observed by *Swift* have been suggested to indicate late time activity of the central engine. As an alternative, I propose here that flares are produced in the deceleration phase of strongly magnetized ejecta. This comes out naturally from Poynting models where instabilities give partial magnetic dissipation (prompt emission and acceleration; see, e.g., Giannios & Spruit 2006). For high bulk Lorentz factor of the flow, the instabilities slow down because of the time dilation effect and

the flow remains strongly magnetized in the afterglow phase. The deceleration of the flow due to interaction with the external medium (and possibly the crossing of the reverse shock; Thompson 2005) revives the instabilities and leads to delayed magnetic dissipation.

I have looked at the magnetic energy available in the flow to power such flares and how it depends on the length scales of the dissipation region. For a flare to have the observed rapid evolution, the scales of this region cannot be arbitrarily large, and this constrains the available energy per flare. I estimate that flares caused by delayed reconnection can be powerful for both long and short GRBs, with isotropic equivalent energies up to a fraction $\sim 0.1 \cdot (\delta t_f / t_f)^3$ of the blast wave energy. Although this estimate contains uncertainties related to the speed of magnetic reconnection and the deceleration profile of the ejecta (see Eq. (12)), it indicates that powerful flares through delayed magnetic dissipation are possible and that smooth flares are expected to be more energetic than spiky ones. Multiple flares are expected in this model as a result of dissipation in multiple neighboring regions in the decelerating flow. Furthermore, inverse Compton scattering of the flare photons that takes place in the forward shock can result in GeV–TeV flares as shown by Wang et al. (2006).

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