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New Journal of Physics

The electromagnetic model of gamma-ray bursts

M Lyutikov

University of British Columbia, 6224 Agricultural Road, Vancouver, BC V6T 1Z1, Canada and Department of Physics and Astronomy, University of Rochester, Bausch and Lomb Hall, PO Box 270171, 600 Wilson Boulevard, Rochester, NY 14627-0171, USA E-mail: lyutikov@phas.ubc.ca

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Abstract. The electromagnetic model (EMM) of gamma-ray bursts (GRBs) and a contrast of its main properties and predictions with the hydrodynamic fireball model (FBM) and its magnetohydrodynamical extension are described. The EMM assumes that rotational energy of a relativistic, stellar-mass central source (black hole-accretion disk system or fast rotating neutron star) is converted into magnetic energy through a unipolar dynamo mechanism, propagated to large distances in the form of relativistic, subsonic, Poynting flux-dominated wind and is dissipated directly into emitting particles through current-driven instabilities. Thus, there is no conversion back and forth between internal and bulk energies as in the case of the fireball model. Collimating effects of magnetic hoop stresses lead to strongly non-spherical expansion and formation of jets. Long and short GRBs may develop in a qualitatively similar way, except that in the case of long burst ejecta expansion has a relatively short, non-relativistic, strongly dissipative stage inside the star. EMMs and FBMs (as well as strongly and weakly magnetized fireballs) lead to different early afterglow dynamics, before deceleration time. Finally, the models are discussed in view of latest observational data in the Swift era.

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1. Short introduction

Gamma-ray bursts (GRBs) are conventionally divided into two classes, short-hard and long-soft, distinguished by their duration (with a division near ~ 2 s) and spectrum hardness (Kouveliotou *et al* 1993). Detection of Type Ic supernovae nearly coincident with long GRBs unambiguously linked them with deaths of massive stars (Hjorth *et al* 2003, Stanek *et al* 2003). Studies of the host galaxies of long GRBs, which turned out to be actively star-forming, further strengthen this association (Djorgovski *et al* 1998). Recent progress in observations of short bursts showed that on one hand they show qualitatively similar afterglow behaviour (but without any supernovae signature), while on the other hand their energetics was two to four orders of magnitude smaller and they are preferentially (at the time of writing three out of four) associated with older stellar populations (Covino *et al* 2005, Fox *et al* 2005, Gehrels *et al* 2005, Prochaska *et al* 2005, Retter *et al* 2005, Villasenor *et al* 2005). This indirect evidence is consistent with formation of short GRBs in compact star mergers (double neutron stars or black hole–neutron star binaries) and formation of a black hole (e.g. Rosswog *et al* 2003, Aloy *et al* 2005).

2. Short and long GRBs

The association of two types of GRBs with different astronomical objects is somewhat surprising given their apparent similarity (perhaps less surprising in view of the fact that some short GRBs may be associated with nearby SGRs). One possible reason is that though short and long GRBs occur in different astrophysical settings, their appearance is governed by similar physical processes related to formation and early evolution of stellar mass relativistic compact objects. (Similarities of temporal and spectral properties of the first 2 s in long bursts and short bursts (Nakar and Piran 2002, Ghirlanda *et al* 2003) may be an indication of this.) But then one expects that during a merger of, e.g. two neutron stars, the resulting black hole has large angular momentum and thus can potentially release much more energy than observed (one can invoke different efficiencies, but a naive guess would be that it is harder to extract and propagate energy from a compact object inside a stellar core, contrary to observations).

So why are short GRBs so under-energetic when compared with long ones? One possibility is that the presence of a disk is a necessary condition for extracting energy from a black hole, so after the disk disappears energy extraction stops. In this case, the energy and angular momentum that will power a GRB outflow effectively come not from the central black hole but from the surrounding disk (in the sense that it is the energy and/or lifetime of the disk and not the energy in the black hole that determines the resulting energy of a GRB outflow) (for related discussion see also van Putten (2005)).

In case of neutron star mergers, a black hole forms fairly early, while the mass of the accretion disk is small, $\leq 0.1 M_{\odot}$, with short viscous timescales, $\sim 0.1-1$ s (e.g. Ruffert and Janka 2001). On the other hand, a black hole inside a collapsing core of a massive star may accrete several solar masses of material (e.g. MacFadyen *et al* 2001) (at any given moment the mass of the disk is small, but a large amount of mass, $\sim 1-10M_{\odot}$ passes through the disk during accretion). In addition, the amount of the rotational energy stored in the case of core collapse depends on core rotation before the explosion (which, in turn, depends on metallicity through wind angular momentum loss (Woosley and Heger 2005)), resulting in a broad spread of rotational energies.

3. Principal issues: electromagnetic and fireball models (FBMs)

In this paper, the electromagnetic model (EMM) of GRBs, which assumes that the energy that will power a GRB comes from rotational kinetic energy of a central source, is described. The energy is extracted through a magnetic field, which can be generated by local dynamo mechanisms (e.g. Pugliese *et al* 1999, Proga *et al* 2003, McKinney and Gammie 2004, De Villiers *et al* 2005, Hawley and Krolik 2005). As argued above, the GRB energy should then be related not to the total rotational energy of a central black hole but to the disk around it. Another possibility for long bursts is the formation of a 'millisecond magnetar', a fast rotating strongly magnetized protoneutron star.

Whether magnetic fields play an important dynamical role at any stage in the outflow remains, in our view, one of the principal issues in GRBs physics. Currently, the overwhelming point of view, advocated by the FBM, is that magnetic fields do not play any major dynamical role (except, perhaps, at a very early stage, after which fields are dissipated quickly). FBM advocates that in the emission region magnetic fields are re-created locally (e.g. through development of Weibel instability (Medvedev and Loeb 1999)), with energy density typically much smaller

than plasma energy density. Fields are small scale, with correlation length l_c much smaller than the 'horizon' length $l_c \ll R/\Gamma$ (*R* is the radius of the outflow in the laboratory frame and Γ is its bulk Lorentz factor). An alternative approach, advocated by MHD and EMMs (i.e. Usov 1992, Blandford 2002, Lyutikov and Blandford 2003, Vlahakis and Königl 2003) is that there are dynamically important large scale fields with 'super-horizon' correlation length $l_c \ge R/\Gamma$, which are created at the source and which may play a major role in driving the whole outflow in the first place.

To quantify the dynamical importance of large-scale magnetic fields, it is useful to introduce the magnetization parameter σ as a ratio of Poynting $F_{Poynting}$ to (cold) particle F_p fluxes (or as a ratio of rest frame energy densities)

$$\sigma = \frac{F_{\text{Poynting}}}{F_{\text{p}}} = \frac{B^2}{4\pi\Gamma\rho c^2} = \frac{b'^2}{8\pi\rho' c^2},\tag{1}$$

where *B* and ρ are magnetic field and plasma density in the laboratory frame, *b'* and ρ' are magnetic field and plasma density in the plasma frame (where the electric field is zero). For $\sigma \ll 1$ magnetic fields are dynamically unimportant (this is assumed within a framework of a conventional FBM), while for $\sigma \ge 1$ magnetic fields start to play an important dynamical role. For $\sigma \gg 1$, there is an important qualitative change in the dynamical behaviour of the flow at $\sigma_{\rm crit} = \Gamma^2/2$. For $\sigma < \sigma_{\rm crit}$ the flow is super-Alfvenic, while for $\sigma > \sigma_{\rm crit}$ the flow is sub-Alfvenic (and sub-fastmagnetosonic). The difference is somewhat analogous to the difference between subsonic and supersonic flows in hydrodynamics.

Thus, depending on the parameter σ three qualitatively different regimes for expansion of the ejecta may be identified, which can be called (i) FBM (below) $\sigma \ll 1$, (ii) MHD models $1 \leq \sigma < \sigma_{crit}$ and (iii) EMM (below) $\sigma > \sigma_{crit}$. These three possibilities lead to *a qualitatively different dynamic behaviour of flows*. Let us next describe qualitatively the main features of the models. (As the FBM and EMM are at the extreme range of σ we discuss those first.)¹

FBM (e.g. Piran 2004). The defining characteristic of the FBM is that at intermediate distances (far from the central source but before most energy is transferred to the forward shock) most energy produced by the central source is carried by the bulk motion of ions. In temporal order, the transformations of energy are as follows. Initially, the energy that will power the GRB and its afterglow is thermalized near the central source, so that most of it is converted into lepto-photonic plasma. This internal energy is then converted to the bulk motion of ions, and reconverted back into internal at internal shocks; at the same time, small-scale magnetic fields are generated. The energy of these generated magnetic fields is then used to accelerate leptons via the Fermi mechanism to highly relativistic energies.²

EMM (Lyutikov and Blandford 2003). The defining characteristic of the EMM is that the bulk energy of the flow is carried *subsonically* by the magnetic field. In temporal order, the evolution of the energy proceeds as follows. The energy that will power a GRB comes from kinetic rotational energy of the central source (millisecond pulsar or BH–disk system). It is then converted to magnetic energy using a unipolar inductor (as in pulsars), transported to large distances in a form of strongly magnetized wind and is used to accelerate particles in the emission region. Acceleration of particles is done via magnetic dissipation (not through shocks).

¹ The definitions and discussion below are based not on the nature of the central object but on ejecta content at large distances from the central region and before production of γ -rays occurs.

 $^{^2}$ The energy that goes to non-thermal particles is *electromagnetic* even in the FBM: Fermi-type acceleration is done by turbulent EMF associated with fluctuations of magnetic field.

MHD model (e.g. Drenkhahn and Spruit 2002, Vlahakis and Königl 2003). In this case, most energy is also carried by magnetic field (similar to Lyutikov and Blandford 2003), but the flow is supersonic, similar to FBM. In the version of Drenkhahn and Spruit (2002), magnetic field energy is first converted into bulk motion and then dissipated through internal shocks, similar to FBM. This is, in principle, not necessary, so that magnetic field energy may be dissipated directly into emitting particles, similar to EMM.

The principal differences between the EMM and MHD approaches are that MHD-type outflows usually cross the fast magnetosonic critical surface after which moment they become causally disconnected from their source (Goldreich and Julian 1970). Initially the flow is expanding freely, so that the flow dynamics is determined by the internal structure of the flow. Only after the flow reaches the terminal velocity does the interaction with the media become important. Unlike MHD, force-free flows are sub-fast magnetosonic, so no conditions at the fast critical surface appear. In this case, it is the *interaction with boundaries that determines the properties of the flow* (similar to subsonic hydrodynamic flows). Thus, the distinctive feature between MHD and force-free flows is whether the wind becomes fast supersonic (MHD regime, $\sigma < \sigma_{crit}$) or not (force-free regime, $\sigma > \sigma_{crit}$). This important difference leads to somewhat different dynamics of the flow and can be tested with observations, as discussed in subsection 10.2.

4. Source formation and energy release in EMM

4.1. Electromagnetic luminosity and currents

In this section, the main ingredients of the EMM, stressing its principal difference and predictions from the FBM, have been described. EMM assumes that the GRB 'prime mover' is a relativistic, fast rotating, near-stellar-mass source. As discussed above, in order to reconcile energies of short and long GRBs the 'prime mover' should not be the black hole but the disk around it. For numerical estimates, it is assumed that a central source generates luminosity $L = L_{50} \times 10^{50}$ erg s⁻¹ for a time t_s , where $t_s \sim 100$ s for long bursts and $t_s \sim 1$ s for short ones (there are indications that both long and short bursts are powered by the same luminosity, but for different time (Fox *et al* 2005)). The mass of the central source is $\sim 0.01 M_{\odot}$ for short bursts and $\sim M_{\odot}$ for long bursts (we stress again that this is the total mass passing through the disk, and not in the black hole). The source is assumed to rotate with a spin frequency \sim kHz. In addition, it is assumed that a source possesses a large magnetic field of $B_s \sim 10^{14} L_{50}^{1/2}$ G. If initially the core is fast rotating (e.g. Woosley and Heger 2005), the total rotational energy in the disk, $\sim M_{disk} R_{disk}^2 \Omega^2$, in the case of core collapse is much larger, $\sim 7.9 \times 10^{52} (M_{disk}/M_{\odot}) (R_{disk}/\times 10^6 \text{ cm})^2 (\omega/6.28 \times 10^3 \text{ rad s}^{-1})^2 \text{ erg}$, than in the case of mergers, $\sim 7.9 \times 10^{50}$ erg with $M_{disk} \sim 0.01 M_{\odot}$ in the above estimate. The source is expected to be active for $t_s \sim E/L \sim 100$ s for long bursts and ~ 1 s for short bursts.

Rotational energy of the central object is extracted by magnetic fields through a unipolar induction mechanism, similar to the prevailing model of active galactic nuclei (AGN) jets (e.g. Lovelace 1976, Ferrari 2004). Magnetic fields both launch the jet (e.g. through Blandford–Znajek mechanism) and collimate it by hoop stresses. The latest full relativistic MHD numerical simulations of accretion disk–black hole systems do show formation of the strongly magnetized axial funnel (e.g. McKinney and Gammie 2004, De Villiers *et al* 2005). Thus, *large-scale, energetically dominant magnetic fields may be expected in the launching region of GRB jets.*

Qualitatively, in the immediate vicinity of the source, the plasma is separated into two phases: an internal matter-dominated phase in which large currents are flowing and an external magnetically dominated phase. Strong magnetic fields and magnetic flux are generated in the dense medium (a disk or a differentially rotating neutron star-like object), while relativistic outflow is generated in the magnetically dominated phase. In this case, matter loading may be expected to be small (e.g. analogous to pulsar wind). As the source remains active for about a thousand to a million dynamical times, the flow will be able to settle down quickly to a quasi-steady state evolving slowly as the hole or neutron star slows down. The separation into matter-and magnetic field-dominated phases is somewhat similar to the Sun, where the dynamo operates in the tachocline, deep below the surface, while magnetic energy and, most importantly, magnetic flux are dissipated outside the star.

The key assumption of the model is that the dissipation rate at the source remains low enough such that the power continues to be dominated by the electromagnetic component (rather than the heat of a fireball) well out into the emission region. Thus, electrical currents flow all the way out to the expanding blast wave, rather than being dissipated close to the source. For electromagnetically dominated outflows, the value of the total current may be related to the total luminosity of the source by

$$I \sim \sqrt{\frac{Lc}{4\pi}} \sim 3 \times 10^{20} L_{50}^{1/2} \text{A},$$
 (2)

where the notation $X_n = (X/10^n)$ has been adopted. Under general electromagnetic and relativistic conditions, the total impedance of the source and the emission region is close to the impedance of free space $Z \sim 100 \Omega$.

In case of long bursts, associated with collapse of massive stars, the source will initially inflate a non-relativistically expanding electromagnetic bubble inside the star. This magnetized cavity is separated from the outside material by the (tangential) contact discontinuity (CD) containing a surface Chapman–Ferraro current. This current terminates the magnetic field and completes the circuit that is driven by the source. On a microphysical level, the current is created by the particle of the surrounding medium completing a half turn in the magnetic field of the bubble, so that the thickness of the current-currying layer is of the order of the ion gyro-radius. After the breakout, the density that controls the ejecta expansion falls down, so that the expansion becomes relativistic. In case of short bursts, associated with merger of two neutron stars, a somewhat similar process will happen, except that there is no non-relativistic stage, so that the bubble is directly inflated in the circumburst medium.

4.2. Distribution of current in the wind: structured jet

The form of expanding bubble depends on lateral distribution of source luminosity, which within a framework of EMM mode is given by lateral distribution of current. At a relativistic stage, expansion is nearly ballistic (Shapiro 1979) while at a non-relativistic stage inside a star, a flow may be collimated both through the action of magnetic hoop stresses *and* interaction with the surrounding gas (see subsection 5.2). One particular stationary outflow configuration, which captures the essential features of the outflow, is that the outgoing current is confined to the poles and the equatorial plane and closes along the surface of the bubble, figure 1. This current distribution minimizes the total energy given a total toroidal magnetic flux and has been advocated in relativistic stationary winds (Heyvaerts and Norman 2003). The magnetic field in the bubble is



Figure 1. Current flow in the electromagnetic bubble. Current flow is mostly along the axis, on the surface of the magnetic shell, along equator and close-up at the trailing part of the shell. The magnetic shell is preceded by the forward shock, typically r/Γ^2 ahead of it. Non-sphericity of the shell, which is of the order $\sim 1/\Gamma^2$, is enhanced.

inversely proportional to the cylindrical radius, $B_{\phi} \propto 1/(r \sin \theta)$. Accompanying this magnetic field is a poloidal electrical field so that there is a near-radial Poynting flux carrying energy away from the source. Thus, as the magnetic field strength is strongest close to the symmetry axis, the bubble will expand fastest along the polar direction. The internal structure of an outflow then corresponds to a 'structured jet' with $L_{\theta} \sim \theta^{-2}$ (Lipunov *et al* 2001, Rossi *et al* 2002), so that the central source releases an equal amount of energy per decade of θ .

5. Long GRBs: expansion inside a star

5.1. How important is dissipation?

Some fraction of the central source luminosity is likely to be dissipated close to the source. The FBM implicitly assumes that all of the energy released is quickly converted into heat, while in the EMM this does not happen and the energy flows away from the light cylinder mainly in the form of an electromagnetic Poynting flux and the load impedance is located in the emission region.

The issue of dissipation is somewhat complicated, as discussed next. On one hand, somewhat paradoxically, it becomes harder to convert electromagnetic energy directly to the pair plasma the stronger the magnetic field becomes. The reason is that the maximum potential drop that is available for dissipation will be limited by various mechanisms of pair production. Typically, after an electron has passed through a potential difference $\Delta V \sim 10^9-10^{12}$ V it will produce an electron–positron pair either through the emission of curvature photon or via inverse Compton scattering. This will be followed by an electromagnetic cascade and the newly born pairs will create a charge density that would shut-off the accelerating electric field. In other words, the pair density required to supply the electrical current and space charge scales linearly with the field strength, while the electromagnetic energy density scales as its square. The stronger

the field, the more likely it is to persist into the outflow. This is because GRBs are so powerful that the dissipation in the source is probably low.

There is an important caveat to the preceding discussion which applies to very early, nonrelativistic, stages of bubble expansion in case of long bursts. Somewhat similar to pulsar wind nebulae (cf Rees and Gunn 1974), non-relativistic, ideal, homologous expansion of strongly magnetized nebulae which is causally disconnected from the source and injects toroidal magnetic field at nearly the speed of light cannot occur. The reason is that magnetic flux and energy are supplied to the inflating bubble by a rate that cannot be accommodated in the bubble. The rate of supply is determined by the processes inside the light cylinder of a newly formed, compact object, while inflation of the bubble is controlled by the external gas density. (Even if the wind remains subsonic, it is unlikely that processes at the edge of the inflating bubble would influence wind generation region near the light cylinder.) This leads to the following 'contradiction'. Magnetic flux (integrated over the meridional plane) is supplied to the bubble at a rate $\dot{\Phi} \sim 2Ic$. Similarly, energy is supplied to the bubble at a rate $\dot{U}_{\rm EM} \sim \mathcal{E}I$. We can also compute the magnetic flux $\Phi = \mathcal{L}I$ and the energy stored within the bubble $U_{\rm EM} = \mathcal{L}I^2/2$, using the self-inductance \mathcal{L} . Let the bubble radius be $R(\theta, t)$ where θ is the polar angle measured from the symmetry axis defined by the spin of the compact object. If the magnetic field in the bubble is predominantly toroidal between cylindrical radii ϖ_{\min} and $\varpi_{\max} = R \sin \theta$, this is given by

$$\mathcal{L} \sim \frac{\mu_0}{2\pi} \int \mathrm{d}z \ln\left(\frac{\overline{\varpi}_{\mathrm{max}}}{\overline{\varpi}_{\mathrm{min}}}\right). \tag{3}$$

We therefore see that, if the bubble expands sub-relativistically, the rate of supply of both flux and energy exceeds the rate at which the flux and energy can be stored by a factor $\sim [\ln(\varpi_{\max}/\varpi_{\min})(dz/dt)/c]^{-1}$ (cf Rees and Gunn 1974). Therefore, for non-relativistic expansion of the bubble, too much flux and too much energy are generated by the source.

The way out of the 'paradox' is that dissipation must become important, which will destroy some magnetic energy and most importantly eliminate most of the toroidal flux. Most of the dissipation is likely to occur near the axis where the current density is highest and the susceptibility to current-driven instability is greatest. In this case, a lateral flow of energy will set in, carrying the poloidal field lines with it towards the axis. This, in turn, will lead to the pile-up of magnetic field near the axis and to faster radial expansion near the axis (the toothpaste tube effect) (Lyutikov and Blandford 2003).

5.2. Form of the expanding bubble

The dynamics of a non-spherically expanding bubble may be described using the method of Kompaneets (1960), and Zel'dovich and Raizer (1967). Consider a small section of non-spherical non-relativistically expanding CD with radius $R(t, \theta)$. The CD expands under the pressure of a magnetic field so that the normal magnetic stress at the bubble surface is balanced by the ram pressure of the surrounding medium. At the spherical polar angle θ , the CD propagates at an angle tan $\alpha = -\partial \ln R/\partial \theta$ to the radius vector. Balancing the pressure inside the bubble $B^2/(8\pi) = I^2/(2\pi c^2 R^2)$ with the pressure of the shocked plasma gives

$$\left(\frac{\partial R}{\partial t}\right)^2 = \kappa \frac{I^2(t)}{2\pi R^2 \sin^2 \theta \rho(R,\theta)} \left[1 + \left(\frac{\partial \ln R}{\partial \theta}\right)_{\rm T}^2\right],\tag{4}$$

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where κ is a coefficient of the order of unity which relates the pressure at the CD to the pressure at the forward shock.

Equation (4) shows that non-spherical expansion inside the star is due to both the anisotropic driving by magnetic fields *and* collimating effects of the stellar material (the term in parenthesis, which under certain conditions tends to amplify non-sphericity). The rate of expansion of the bubble inside the star depends upon the density profile of the stellar envelope and the time evolution of the luminosity (or, equivalently, of the current I(t)). For a given dependence $\rho(R, \theta)$ and I(t), equation (4) determines the velocity of the CD. Generally, solutions will be strongly elongated along the axis. A simple analytical solution for I, $\rho \sim$ const. is

$$R(t,\theta) = \left(\frac{2}{\pi}\frac{I^2}{\rho c^2}\right)^{1/4} \frac{\sqrt{t}}{\sin\theta}.$$
(5)

(current is related to the luminosity by equation (4)). Qualitatively, the bubble and the forward shock will cross the iron core ($r_c \sim 2.5 \times 10^8$ cm) in $t \sim r_c \sqrt{\rho}/B(r_c) \sim 0.3$ s, short enough to produce an ample supply of ⁵⁶Ni (Woosley *et al* 2003). If M(R) is the stellar mass external to radius *R*, then the breakout time is

$$t_{\text{breakout}}(\theta) \sim 1\theta_{-1}^2 \left(\frac{M}{M_{\odot}}\right)^{1/2} \left(\frac{R}{R_{\odot}}\right)^{1/2} L_{50}^{-1/2} \text{s.}$$
 (6)

The electromagnetic bubble can be confined equatorially by the star for the duration of the burst $t_{\text{breakout}}(\pi/2) \sim 100 \text{ s}$ and will expand non-relativistically as we have assumed. However the expansion along the axis proceeds on a short timescale and breakout should occur early in the burst.

Thus, along the jet axis non-relativistic expansion lasts for several seconds, which is much shorter than the burst duration. After breakout, the flow quickly becomes relativistic so that there is no longer any need for dissipation. Thus, relative fraction of dissipated energy is small, so that overall the flow magnetization remains large. Most of the dissipation described above will result in creation of a lepto-photonic plasma, which decouples after photosphere, so that the remaining flow remains strongly magnetized. In summary, when the bubble expands non-relativistically, it must be dissipative, while after breakout, the expansion becomes relativistic and the resistance falls so that the electromagnetic energy that is still being supplied by the source is mostly absorbed by the inflating bubble and by doing work against the surroundings.

6. Optically thick expansion: mini-fireball

Under the electromagnetic hypothesis, most of the energy released by the source comes out in the form of Poynting flux. However, there must be some dissipation that would lead to creation of a lepto-photonic component (as discussed in subsection 5.1); in the case of non-relativistic stage of bubble expansion inside a star, dissipation may be considerable. This will create an optically thick 'warm' fireball (in a sense that it is dominated by magnetic field energy, but also has considerable pressure). Expansion of this 'warm' fireball will create a thermal precursor, similar to the conventional FBM, but modified by presence of magnetic field (Lyutikov and Usov 2000, Lyutikov and Blandford 2003).

At early stages (before the breakout in case of long bursts), the plasma enthalpy is strongly dominated by a lepto-photonic plasma with a temperature

$$T \sim \left(\frac{L}{a\Delta\Omega r^2\beta c\Gamma^2(1+\sigma)}\right)^{1/4},\tag{7}$$

where $\Delta\Omega$ is a typical opening solid angle, and luminosity of the source can then be written as $L = \int d\Omega \Gamma^2 r^2 \beta c (b^2/2 + w)$, where w is plasma enthalpy and b is a toroidal magnetic field in the plasma rest frame times $\sqrt{4\pi}$,

After breakout, the flow will accelerate to relativistic velocities. Initially, conical expansion is mostly pressure-driven, even in the strongly magnetized case (in this case, magnetic pressure gradients and hoop stresses balance each other out). This results in dynamics qualitatively similar to the unmagnetized case: the wind plasma accelerates $\Gamma \sim r$ while its density, pressure and temperature decrease $n \sim r^{-3}$, $p \sim r^{-4}$, $T \sim r^{-1}$, so that the magnetization parameter remains approximately constant (Lyutikov and Blandford 2003).

When the temperature falls below ~10–20 keV, most of the pairs annihilate. This suddenly reduces the optical depth to Thomson scattering below unity. (Under certain conditions, photons may remain trapped in the flow. In this case, thermal driving by photon pressure continues, until the thermal photons escape.) As a result the lepto-photonic part of the flow decouples from the magnetic field and σ increases by roughly seven orders of magnitude to $\sigma \sim 10^9$. The thermal radiation from the lepto-photonic component has a rest-frame temperature $T_0 \sim 10-20$ keV times a boost due to the bulk motion. This thermal radiation, which should peak around ~100 keV may put constraints on the initial σ (Lyutikov and Usov 2000, Daigne and Mochkovitch 2002).

7. Relativistic expansion

7.1. Short GRBs and long GRBs after breakout

In the case of long GRBs inflating a bubble inside a star, eventually the bubble will break free of the star forming two axial jets along which Poynting flux will flow until the central source slows down on the timescale $t_s \sim 100$ s. Outside the star, the bubble will expand ultra-relativistically and bi-conically. For short GRBs, presumably associated with merger of neutron stars in a low-density environment, there is no preceding non-relativistic stage, so expansion of the bubble is relativistic from the beginning. In case of relativistic motion, there is no longer any necessity to destroy magnetic flux through ohmic dissipation: the effective load can consist of the performance of work on the expanding blast wave. This is where most of the power that is generated by the central magnetic rotator ends up.

After the bubble has expanded beyond a radius $r_{sh} \sim ct_s \sim 3 \times 10^{12} t_{s,2}$ cm (~10¹⁰ cm for short bursts), the electromagnetic energy will be concentrated within an expanding, electromagnetic shell with thickness ~ r_{sh} and with most of the return current completing along its trailing surface (see figure 1). The global dynamics of this shell and its subsequent expansion are set in place by the electromagnetic conditions at the light cylinder and within the collimation region. An important property of ultra-relativistic outflows is that they are hard to collimate (Chiueh *et al* 1991, Bogovalov 2001), so that any collimation should be achieved close to the source, within a star, where the flow is only mildly relativistic (see, however, Vlahakis and Königl 2003).

Interaction of the magnetic shell with the circumstellar medium proceeds in a similar way to the non-relativistic expansion inside a star: the leading surface of the shell is separated by a CD (which actually becomes a rotational discontinuity if the circumstellar medium is magnetized (Lyutikov 2002)). Outside the CD, an ultra-relativistic shock front forms and propagates into the surrounding circumstellar medium. The expansion is non-spherical. As long as the outflow is ultra-relativistic, the motion of the forward shock is virtually ballistic (Shapiro 1979) and determined by the balance between the magnetic stress at the CD and the ram pressure of the circumstellar medium.

A type of collimation in the case of electromagnetic explosions is somewhat different from the conventional jet models of AGNs. We expect that large Poynting fluxes associated with explosive release of $\sim 10^{51}$ ergs in the case of long GRBs are sufficient to drive a relativistic outflow over a large solid angle, so that during the relativistic stage the resulting cavity is almost spherical, but the Lorentz factor Γ of the CD is a strong function of the polar angle. The angular distribution of magnetic field (and of the Lorentz factor of the expansion) depends on the dynamics of the bubble at the non-relativistic stage and the distribution of the source luminosity.

In the framework of EMM, the outflow is strongly magnetized and *subsonic*, $\sigma > \sigma_{crit}$, despite being strongly relativistic. In this case, the ejecta, in some sense, may be considered as a collection of outgoing fast magnetosonic waves propagating from the source to the CD. Motion of the CD is then determined by the pressure balance between the Poynting flux from the source and the ram pressure of the interstellar medium (ISM). Thus, motion of the CD depends on the source luminosity L(t') at the retarded time t' such that R(t) = t - t'. In addition to forward flux, there is a much weaker, by a factor Γ^2 , reflected flux that propagates backward into the flow information about the circumstellar medium. Interference of forward and backward propagating waves allows us to define a finite Lorentz factor of the ejecta. The distribution of reflected current is determined by the outgoing current and the boundary conditions. At later times, multiple reflections from the CD and the centre become important as well.

7.2. Stages of relativistic expansion of the electromagnetic shell

Relativistic expansion of the magnetized shell may be separated into two stages, which will be called 'early' and 'late', depending on whether or not most of the fast waves emitted by the central source have caught up with the CD and their energy has been given to the circumburst medium. The transition between two stages occurs at the moment which is similar to the deceleration radius in FBM, except that in the case of EMM the shell is decelerating all the time, but with different laws before and after the transition. Keeping with tradition, the transition radius will be called the deceleration radius.

7.2.1. 'Early' stage. At $r > r_{ph}$ the outflow becomes a relativistically expanding shell of thickness $\sim ct_s \sim 3 \times 10^{12}$ cm for long GRBs and $\sim 3 \times 10^{10}$ cm for short GRBs. The shell contains a toroidal magnetic field; the current now detaches from the source and completes along the shell's inner surface. At this stage, the CD is constantly re-energized by the fast-magnetosonic waves propagating from the central source. The average motion of the CD R(t) is determined by the average luminosity at the retarded time t':

$$L_{\Omega}(t') \sim \rho c^3 \Gamma^4 R(t)^2 \beta^3, \tag{8}$$

which for constant luminosity gives $\Gamma \sim (L_{\Omega}/\rho c^3)^{1/4} r^{-1/2}$ (in a constant density medium) or $\Gamma \sim (L_{\Omega}/4\kappa\rho_0 r_0^2 c^3)^{1/4} = \text{const.}$ (in a $\rho(r) = \rho_0 (r_0/r)^2$ wind). If the central source releases most of the current along the axis and the equatorial plane, as argued in subsection 4.2, then $\Gamma \propto 1/\sqrt{\sin \theta}$ at this stage. If source's luminosity varies, this will be reflected in the 'jitter' of the CD. Development of instabilities at the CD, such as the impulsive Kruskal–Schwarzschild instability (Lyutikov and Blandford 2003) may lead to dissipation and particle acceleration. The internal structure of the magnetic shell is a messy mixture of the outgoing waves from the source and the ingoing waves reflected from the CD, similar to a pre-Sedov phase in hydrodynamical explosions. Unlike the case of a hydrodynamic blast wave with energy supply, no internal discontinuities form inside the magnetic shell.

The early stage lasts for $ct_s < r < r_{dec}$, where

$$r_{\rm dec} = (L_{\Omega} t_{\rm s}^2 / \rho c)^{1/4} \sim 3 \times 10^{16} L_{50}^{1/4} t_{\rm s2}^{1/2} n^{-1/4} \,\rm cm \tag{9}$$

for long bursts (in the observer frame this phase lasts $\sim t_{\rm s} \sim 100$ s) and $r_{\rm dec} \sim 3 \times 10^{15}$ cm for short bursts (similarly, in the observer frame this phase lasts $\sim t_{\rm s} \sim 1$ s). Radius $r_{\rm dec}$ (9) is somewhat similar to the deceleration radius in case of FBM; at this moment most energy of the shell is given to the circumburst medium, $L_{\theta}t_{\rm s} \sim E(\theta) \sim \rho c^2 r_{\rm dec}^3 \gamma (r_{\rm dec}, \theta)^2$ (note that here $\Gamma = \Gamma(r_{\rm dec})$, not Γ_0 as in the case of FBM, since there is no formal definition of Γ_0 in case of EMM).

7.2.2. 'Late stage' $(r < r < r_{\rm NR} \equiv (L_{\Omega}t_{\rm s}/\rho c^2)^{1/3})$. At distances $r > r_{\rm dec}$ most of the waves reflected from the CD have propagated throughout the shell, so that all the regions of the shell come into causal contact. Most of the energy of the explosion will reside in the blast wave which will eventually settle down to follow a self-similar expansion. As the expanding shell performs work on the surrounding medium its total energy decreases; the amount of energy that remains in the ejecta shell during the late stage is small, $\sim E_{\Omega}/\Gamma^2$. Most of the energy is still concentrated in a thin shell with $\Delta R \sim R/\Gamma^2$ near the surface of the shell which is moving according to $\Gamma \sim \sqrt{E_{\Omega}/\rho c^2} r^{-3/2}$ (in a constant density medium), or $\Gamma \sim r^{-1/2}$ (in a $\rho \sim r^{-2}$ wind). If the central source releases most of the current along the axis and the equatorial plane, as discussed in the subsection 4.2, then $\Gamma \propto 1/\sin\theta$ at this stage. (Note that at the 'early stage' $\Gamma \propto 1/\sqrt{\sin\theta}$, but no lateral re-distribution of energy is required at the transition since transition between 'early' and 'late' stages occur at different times for different θ .) The energy of the shell decreases only weakly with radius, $dE/dt \sim 1/r$ in constant density and $dE/dt \sim 1/r^3$ in a wind, so that the surface of the shell keeps moving relativistically as long as the preceding shock wave is moving relativistically, until $r \sim (E_{\Omega}/\rho c^2)^{1/3} \sim 10^{18}$ cm—the shock never becomes completely free of the shell (Lyutikov and Blandford 2003). Interestingly, the structure of the magnetic shell (in particular the distribution of energy) resembles at this stage the structure of the hydrodynamical relativistic blast wave (Blandford and McKee 1976). This can be formally understood by noting that for motion perpendicular to the magnetic field dynamical equations for magnetized flow can be reduced to the non-magnetic case, but with a different equation of state (Landau and Lifshits 1982).

'Late stage' of magnetic shell expansion corresponds to the conventional afterglow phase when synchrotron and inverse Compton radiation is emitted throughout the electromagnetic spectrum. The initially aspheric expansion will give the appearance of a jet with the 'achromatic break' occurring when the fastest Lorentz factor of the spine $\Gamma(\theta = 0)$ becomes comparable with the reciprocal of the observer's inclination angle with respect to the

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symmetry axis, $\Gamma(\theta = 0) \sim 1/\theta_{ob}$. When $r > r_{NR} \sim (Lt/\rho c^2)^{1/3} \sim 2 \times 10^{18} L_{50}^{1/3} t_{s2}^{1/3} n^{-1/3}$ cm, the blast wave become non-relativistic and will become more spherically symmetric, while evolving towards a Sedov solution.

8. Production of GRB

By the time the shell radius expands to r_{dec} , most of the electromagnetic Poynting flux from the source will get caught up with the CD and get reflected by it, transferring its momentum to the blast wave. Simultaneously a strong region of magnetic shear is likely to develop at the *outer* part of the CD (Lyutikov 2002).

We propose that the γ -ray-emitting electrons are accelerated near $r_{dec} \sim 10^{15}-10^{16}$ cm (for long bursts) and $\sim 10^{15}$ cm (for short bursts) due to development of electromagnetic current-driven instabilities (conventional model of particle acceleration—acceleration at internal shocks—cannot work in this model since in the limit $\sigma \gg 1$ fast shocks are either weak or do not form at all). The development of current instabilities usually results in enhanced or anomalous plasma resistivity which leads to an efficient dissipation of the magnetic field. The magnetic energy is converted into heat, plasma bulk motion and, most importantly, into high-energy particles which, in turn, are responsible for the production of the prompt γ -ray emission. The conversion of magnetic energy into particles may be very efficient. For example, recent RHESSI observations of the Sun indicate that, in reconnection regions, most of the magnetic energy goes into non-thermal electrons with power-law distribution in energies (Benz and Saint-Hilaire 2003). Though the details of how magnetic dissipation and particle acceleration proceed in Solar flare and GBR outflows are bound to be different, the underlying principles may be similar. We discuss them next.

To illustrate how magnetic dissipation may proceed, we briefly describe the physics underlying the development of the so-called tearing mode. Consider a smooth distribution of electrical current, which can be viewed as a set of many small current wires. Since parallel currents attract, such a system is likely to develop narrow current sublayers where dissipation, which is inversely proportional to the square of magnetic field gradient, becomes high. In addition, high anomalous resistivity, proportional to local current density is likely to develop. Similar to non-relativistic plasmas, in strongly magnetized plasmas, a tearing mode develops on timescales *much shorter than resistive timescale in the bulk* (Lyutikov 2003). The final outcome of the development of the tearing mode is formation of reconnection sites and dissipation of magnetic energy.

Particle acceleration by dissipative magnetic fields may proceed in a number of ways. The best studied non-relativistic example is particle acceleration in reconnection regions either by inductive electric fields outside the current sheet or resistive electric fields inside the current sheets (e.g. Craig and Litvinenko 2002) or formation of shocks in the downstream of reconnection regions (e.g. Blackman and Field 1994). Investigation of the particle acceleration in the relativistic regime of reconnection is only beginning (e.g. Larrabee *et al* 2003). Relativistic reconnection may produce power-law spectra of accelerated particles (Larrabee *et al* 2003, Zenitani and Hoshino 2004). For example, in the relativistic Sweet–Parker reconnection model (Lyutikov and Uzdensky 2003), if one balances linear acceleration inside the reconnection layer by the resistive electric field, $d_T \mathcal{E} \sim e \mathcal{E}c$ with the rate of particle escape (proportional to relativistic gyro-frequency), $d_T \ln N(\mathcal{E}) \sim \omega_B (mc^2/\mathcal{E})$, one finds $N(\mathcal{E}) \sim \mathcal{E}^{-\beta_{in}}$ where \mathcal{E} is the energy of a particle, $N(\mathcal{E})$ is the particle number and β_{in} is the inflow velocity (Zenitani and Hoshino 2004). For relativistic reconnection the inflow velocity can be relativistic (Lyutikov

and Uzdensky 2003), $\beta_{in} \rightarrow 1$. The fact that reconnection models can produce spectra which are prohibitively hard for shock acceleration may serve as a distinctive property of EMMs.

In addition to the acceleration mechanisms which are based on known non-relativistic schemes, it is feasible that acceleration in relativistic, strongly magnetized plasma may proceed through mechanisms that do not have non-relativistic or fluid analogues. Examples of this type of acceleration include particle acceleration through formation of a spectral cascade of nonlinear waves in force-free plasma which transfer energy to progressively larger wave vectors until this energy is taken up in accelerating a population of relativistic electrons and positrons (Thompson and Blaes 1998). Since for $\sigma > 1$ the cascade is likely to be terminated at plasma frequency, which is lower by a factor $\sqrt{\sigma}$ of the cyclotron frequency, the likely emission mechanism in this case is inverse Compton scattering. Another possibility is development of kinetic electromagnetic-type instabilities of the shell surface currents, as proposed by Smolsky and Usov (1996) and in somewhat different form by Liang and Nishimura (2003). Since studies of kinetic properties of strongly magnetized relativistic plasmas are only beginning, it is hard to predict acceleration efficiency and particle spectra. Numerical studies in the coming years will be most important here.

9. Production of afterglows

Except at the early stage (as discussed in subsection 7.2.1), afterglows are generated in a similar way both in the FBM and EMM. As the magnetic shell expands, its energy is gradually transferred to the preceding forward shock wave. Relativistic particles are accelerated in the blast wave producing the observed afterglow in a manner which is similar to that proposed for fluid models, except that the CD itself may be an important source of magnetic flux through impulsive Kruskal–Schwarzschild instability (Lyutikov and Blandford 2003), so that afterglows may result from a mixture of relativistic particles, derived from the shock with magnetic field derived from the shell.

At late times, well beyond r_{dec} (which in the observer frame is nearly coincident with the prompt phase), the temporal behaviour of proper afterglow (as opposed to tails of prompt emission, see subsection 10.3) is determined by the total energy release E_{ω} and not by the form of that energy. As a consequence, late afterglow observations can hardly be used to distinguish between the models. The only property of the source that the forward shock 'remembers' at late times is the angular distribution of the deposited energy $E(\theta)$ (there is little sideways evolution in the relativistic regime (Shapiro 1979)). Thus, the angular distribution of the total energy $E(\theta)$ can be used to distinguish between different models, if a model predicts it.

In case of EMM, the preferred lateral distribution of the magnetic field, energy fluxes and luminosity correspond to the line current, subsection 4.2, so that $L \sim 1/\sin^2 \theta$. This translates into distribution of Lorentz factors of the forward shock

$$\Gamma \sim \left(\frac{E}{\rho_{\rm ex}c^5}\right)^{1/2} \frac{t^{-3/2}}{\sqrt{\theta^2 + \theta_0^2}},\tag{10}$$

where θ_0 is the angular size of the core of the jet. (Its minimal size is the magnetic Debye radius, $r_D = \sqrt{I/2\pi nec}$, which gives $\theta_0 \sim (m^2 c^5 \sigma^2 \Gamma^2 / Le^2)^{1/4} \approx 10^{-3} L_{50}^{-1/4} \sigma_9^{1/2} \Gamma_2^{1/2}$ (Lyutikov and Blandford 2003).) This type of shock has been named 'structured jet' (or universal jet) (Lipunov *et al* 2001, Rossi *et al* 2002), though in our model there is no proper 'jet', but simply

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a non-spherical outflow. The most intense bursts and afterglows in a flux-limited sample will be seen pole-on and can exhibit achromatic breaks when $\Gamma \sim \theta^{-1}$, which might be mistaken for jets.

In conclusion, the observational appearance of GRB afterglows depends mostly on two parameters: (i) explosion energy (more precisely, on the ratio on the explosion energy to circumstellar density) and (ii) the viewing angle that the progenitor's axis is making with the line of sight. This possibility, that all GRBs (and x-ray flashes (XRFs)) are virtually the same but viewed at different angles, resembles the unification scheme of AGNs.

10. Tests of GRB models

10.1. Testing the FBM: reverse shock emission

Perhaps the simplest test of GRB models could come from observations of emission from the reverse shock propagating in the ejecta, which typically falls into the optical range. FBM predicts strong reverse shock emission, so that the absence of nearly contemporaneous optical emission in most GRBs would be a strong argument against FBM. In MHD models with $\sigma > 1$ reverse shock is weak, while in EMM reverse shock is absent altogether.

Since the possible observation of reverse shock emission may play a decisive role in validating a GRB theory, we next discuss briefly properties of the reverse shock expected within a framework of FBM (for more extensive discussion, see, e.g. Sari and Piran 1999, Kobayashi 2000). In the framework of FBM, both reverse shock and internal shocks which produce prompt γ -ray emission originate in the same fluid and have similar properties (e.g. being weakly relativistic). One can naturally expect that the microphysical properties of the accelerating particle, being complicated and poorly understood, are the same for the same type of shocks. Thus, the conventionally introduced quantities like ϵ_B (magnetic field value with respect to equipartition), ϵ_e (electron energy density with respect to equipartition) and γ_{min} (minimum Lorentz factor of accelerated electrons) must be the same for both cases. As a result, for any given burst, observations of the prompt emission can be used to predict the properties of a corresponding optical flash.

The amount of energy dissipated in the reverse shock is comparable to the energy dissipated in the forward shock and to the total GRB energy (Sari and Piran 1999). The principal difference between prompt and optical emitting electrons is the radii of emission and ratio of cooling to expansion timescales. In the framework of FBM, prompt emission is generated at distances $r_{\text{GRB}} \sim 2\Gamma_0^2 \delta tc \sim 10^{12} - 10^{13} \text{ cm} (\delta t \sim 0.01 \text{ s} \text{ is the variability timescale of the central source and,}$ within the framework of FBM, of the prompt emission), while reverse shock emission is typically $generated at distances <math>R \sim 2ct_s\Gamma_0^2 \sim 10^{16} \text{ cm}$ (seen at observer time $t_{\text{obs}} \sim t_s$; we concentrate on a simple so-called 'thick shell' case). Using conventional fireball parameterization for minimum Lorentz factor of accelerated particles $\gamma_{\min} \sim \epsilon_e (m_p/m_e)\Gamma_s$, where Γ_s is the shock Lorentz factor, and parameterizing energy density of magnetic field in the plasma rest frame to ion energy density, $\rho c^2 = L/(4\pi\Gamma_0^2 r^2 c)$ and $b = \sqrt{\epsilon_B}\sqrt{8\pi\rho c^2}$, the ratio of prompt to reverse shock frequencies is

$$\frac{\omega_{\rm GRB}}{\omega_{\rm RS}} \sim \frac{r_{\rm RS}}{r_{\rm GRB}} \left(\frac{\Delta\Gamma}{\Gamma_{\rm RS}}\right)^2 \sim \frac{r_{\rm RS}}{r_{\rm GRB}} \sim \frac{t_{\rm s}}{\delta t},\tag{11}$$

so that the peak of reverse shock emission occurs at $\sim 1-10 \text{ eV}$.

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An important qualitative difference between prompt and optical emitting electrons is that the former are in the fast cooling regime, while the latter are in the slow cooling regime. The radius beyond which optically emitting electrons enter the slow cooling regime,

$$r_{\rm cool} \sim \epsilon_B \epsilon_{\rm e} \frac{(\Gamma_{\rm RS} - 1)Lm_{\rm p}\sigma_{\rm T}}{3\pi c^3 m_{\rm e}^2 \Gamma_0^3} \sim 2 \times 10^{14} L_{50} \,\rm cm.$$
(12)

 $(\Gamma_{\rm RS} - 1 \sim 1 \text{ and } \Gamma_0 = 300 \text{ was assumed})$ is typically smaller that $R_{\rm RS}$. As a result, only a small fraction of energy received by an optical electron $\sim R_{\rm RS}/(\Gamma_0 r_{\rm cool}) \sim 0.02$ is emitted; the rest is lost to adiabatic expansion. For a GRB of $E_{\gamma} \sim 10^{51}$ ergs, optical flash would have $E \sim 10^{49}$ ergs. For a typical GRB burst with fluency $\sim 10^{-6}$ erg cm⁻², and duration of ~ 100 s, this will result in an optical flash of magnitude ~ 12 m. Even if we increased the estimate of $R_{\rm RS}$ and duration of optical flash each by an order of magnitude, the resulting optical flash would have ~ 17 m. On the other hand, the brightest bursts may reach fluences $\sim 10^{-4}$ erg cm⁻², which, according to these estimates, can produce an optical flash of ~ 7 m. In addition, adiabatic cooling results in flux decay $\propto t^{-2}$ and a clear radio signal is expected (e.g. Nakar and Piran 2004). Thus, FBM makes a predictions that all GRBs must have optical flashes in the range 12–17 mm, with some variations of a few magnitudes (both brighter and dimmer) depending on particular properties of each burst.

In the Swift era, not a single GRB has shown the predicted behaviour. This is despite the fast on-board optical telescope and a large number of ground-based robotic telescopes (RAPTOR, ROTSEE, TAROT and others). Reverse shock emission is virtually an unavoidable prediction of the FBM, so that the absence of predicted reverse shocks emission in the Swift era argues against the FBM. Naturally, there is a number of ways that through which optical flashes can be suppressed (e.g. cooling of optically emitting electrons on photons of prompt emission (Beloborodov 2002), 'thin shell case', when the reverse shock emission is spread over longer times, producing a weaker signal (e.g. McMahon *et al* 2005)). A possible explanation of the absence of a clear reverse shock signal is that the ejecta plasma is strongly magnetized. In the case when the energy density of magnetic field dominates the total energy density ($\sigma \ge 1$), reverse shock becomes very inefficient in dissipating flow energy (Kennel and Coroniti 1984).

Irregular optical flashes (like GRB 050525a (Klotz *et al* 2005) and GRB 050904 (Boer *et al* 2005)) may be produced by other mechanisms, like γ -ray pair production in front of the forward shock (Thompson and Madau 2000, Beloborodov 2002), or be a low-energy tail of prompt emission.

10.2. EMM: bright early afterglows

FBMs and EMMs make very different predictions for the properties of early afterglows (see figure 2, Lyutikov 2004). According to EMM, at the early afterglow stage the Lorentz factor and peak frequency are larger (and falling with time) than in the FBM (constant Lorentz factor and peak frequency). *Early afterglows in the EMM are more energetic than in FBM*, figure 3, and can blend with the prompt phase.

10.3. Emission radius of prompt photons and early Swift afterglows

One of the surprising early results from the Swift satellite was the detection of x-ray spikes and breaks in light curves at intermediate times, much longer than the burst duration but well before the conventional jet break (e.g. Chincarini *et al* 2005, Nousek *et al* 2005, Tagliaferri



Figure 2. Evolution of Lorentz factors in the FBMs and EMMs for constant external density. Relative normalization of the curves depends on the viewing angle (in the EMM).

et al 2005). Typical behaviour includes fast–slow–fast decay with transitions near 100–1000 s and ~10⁴ s. This presents a real challenge to GRB models, since if the emission is seen 'head on', within an angle $\theta \leq 1/\gamma$, the radii at which these features should be produced correspond to radii much larger than the deceleration radius. At these times, most of the energy is in the forward shock which should produce a smooth light curve. (Late time injection or specific distribution of Lorentz factors are some of the possibilities that are discussed (e.g. Lazzati and Begelman 2005, Zhang *et al* 2005). For a discussion in the framework of the cannonball model see Dado *et al* (2005).)

The initial fast decaying part of afterglows was argued to be a 'sideways' prompt emission, coming from angles $\theta > 1/\Gamma$ (Kumar and Panaitescu 2000, Barthelmy *et al* 2005). If this interpretation is correct, one can determine emission radii of the prompt emission and compare them with model predictions. (We remind that FBM predicts radii of emission $r_{\rm em} \sim 2\Gamma_0^2 c \delta t \sim 10^{12} - 10^{13}$ cm, while EMM predicts $r_{\rm em} \leq r_{\rm dec} \sim 10^{16}$ cm, equation (9)). If emission is generated



Figure 3. X-ray afterglow in the 2–10 keV energy band for a uniform medium of density $n = 1 \text{ cm}^{-3}$ and injected power $L_{iso} = 10^{52} \text{ erg s}^{-1}$. The source is active for 50 s, the assumed redshift is z = 2.5, fraction of energy in electrons $\epsilon_e = 0.1$, fraction of energy in magnetic field $\epsilon_B = 0.001$. Blue line: EMM, red line: FBM (from Mochkovitch *et al*, in preparation).

at $r_{\rm em}$ and is coming to the observer from large angles, $\theta > 1/\Gamma$, its delay with respect to the start of the prompt pulse is $\Delta t \sim (r_{\rm em}/c)\theta^2/2$. For a typical observer angle $\theta \sim 0.1$ and first break of a light curve at $\Delta t \sim 1000$ s, the implied emission radius is $r_{\rm em} \sim 6 \times 10^{15}$ cm. This is at least two orders of magnitude larger than is assumed in the FBM, but is close to the assumption of the EMM. (To be consistent with FBM and variability on short timescales, the Lorentz factor of the flow should be $\Gamma_0 \sim 3000$, but this would imply that emission is strongly de-boosted, $\Gamma_0 \theta \sim 300 \gg 1$.) Interpretation of light curve breaks at $\sim 10^3$ s as being due to prompt emission seen at large angles, $\theta > 1/\Gamma$, is inconsistent with the FBM.

10.4. Fast variability from large radii

If prompt emission is produced at distances $\sim 10^{15}-10^{16}$ cm, how can fast variability, on timescales as short as milliseconds, be achieved? One possibility is that emission is beamed in the outflow frame, for example due to relativistic motion of 'fundamental emitters' (Lyutikov and Blandford 2003). A possible origin of relativistic motion of 'fundamental emitters' may be the fact that in the case of relativistic reconnection occurring in plasma with $\sigma \gg 1$, the outflowing matter reaches relativistic speeds with $\gamma_{out} \sim \sigma$ (Lyutikov and Uzdensky 2003). Internal synchrotron emission by such jets, or Compton scattering of ambient photons will then be strongly beamed in the frame of the outflow.

Consider an outflow moving with a bulk Lorentz factor Γ with randomly distributed emitters moving with respect to the shell rest frame with a typical Lorentz factor γ_{T} . Highly boosted

emitters, moving towards an observer, have Lorentz factor $\gamma \sim 2\gamma_{\rm T}\Gamma$, so that modest values of $\gamma_{\rm T} \sim 5-10 \ll \Gamma \sim 100-300$ suffice to produce short timescale variability from large distances. As the burst progresses, larger angles and more internal jets producing prompt emission become visible. Most of them will be seen from large angles > $1/\gamma_{\rm T}$ in the bulk frame, producing smooth curves. Occasionally, a jet at large viewing angle, $\theta > 1/\Gamma$, but directed towards an observer will be seen, producing an x-ray flare. One expects a break in the light curve at $\Delta t \sim (r_{\rm em}/c)\theta^2/2$, where θ is a viewing angle (in a structured jet model, this is the angle between the jet axis and direction to the observer). Afterglow should start to blend with prompt emission at later times. In figure 4, we plot an example of a prompt light curve in this model (Lyutikov 2006). The model readily explains many unusual properties of early afterglows: (i) x-ray flares and light curve breaks at late times, much longer than conventional prompt GRB duration (extended source activity is not needed!), (ii) fast variability, (iii) gradual softening of the spectrum and (iv) hardening of a spectrum during x-ray flares (Burrows *et al* 2005).

10.5. Observational implications of the EMM

In this section, we give a short discussion of how the main GRB phenomena are (or may be) explained within the framework of EMM.

- Jet break in afterglow. GRB outflows have large opening angles, but do not have a jet in a proper sense. Outflows are non-isotropic so an achromatic break is inferred when the viewing angle is $\theta_{ob} \sim 1/\gamma$.
- Structured jet. The model predicts and gives a theoretical foundation for the 'structured jet' profile of the external shock.
- XRF flashes. Another testable prediction of the model is that much more numerous XRFs should be observed, which may be coming 'from the sides' of the expanding shell, where the flow is less energetic and the Lorentz boosting is weaker. In addition, the total *bolometric* energy inferred for XRFs (from observations of afterglows before radiative losses become important) should be comparable to the total bolometric energy of γ -ray bursts. Generally, the distributions of parameters of XRFs should continuously match those of GRBs.
- Weak thermal precursor. If a fraction $1/\sigma \sim 0.01-0.1$ of the magnetic energy is dissipated near the source, this should produce a thermal precursor with luminosity $\sim 0.01-0.1$ of the main GRB burst.
- Hard-soft evolution. The trend of GRB spectra to evolve from hard to soft during a pulse is explained as a synchrotron radiation in an expanding flow with magnetic field decreasing with radius $B \propto \sqrt{L}/r$ (later in a pulse emission is produced further out where magnetic field is weaker, so that the peak energy will be lower; this is similar to 'radius-to-frequency mapping' in radio pulsars and AGNs).
- Amati $E_{\text{peak}}-L$ correlation. A correlation between peak energy and total luminosity, $E_{\text{peak}} \sim \sqrt{L}$ (Amati L *et al* 2002) follows from the assumption of a fixed typical emission radii and fixed minimum particle energy since $B \sim \sqrt{L_{\Omega}}$ (see also Lovelace 1976).
- Variability. Variability of the prompt emission reflects the statistical properties of dissipation (and not the source activity as in the FBM). Magnetic fields are nonlinear dissipative dynamical system which often show bursty behaviour with power-law PDF. (For example, solar flares



Figure 4. Prompt emission produced by emitters moving randomly in the bulk frame. Emission is generated within a shell of thickness $t_s c = 3 \times 10^{12}$ cm moving with $\Gamma = 100$ at distance $r_{em} = \Gamma^2 t_s c$ by randomly distributed jets with random orientation moving with random Lorentz factors $1 < \gamma_T < \gamma_{T,max} = 5$. Each emitter is active for random time $0 < t'_{em} < 0.5t_s c\Gamma = t_{pulse,max}$ in its rest frame. Homogeneous jet centred on an observer with opening angle $\theta =$ 0.1, dimensionless parameters $N\pi/(\Gamma\gamma_{T,max}\theta)^2 = 1.2$ (probability of seeing one sub-jet 'head-on' from angles $<1/\Gamma$) and $N(ct_{pulse,max}/2)^2/r_{em}^2\theta^2 t_s c\Gamma = 0.19$ (efficiency of energy conversion), where N is the total number of emitters. Intensity of emission is $\propto \delta^{3+\alpha}$, where δ is a Doppler factor and $\alpha = 0.5$ is a spectral index. As the burst progresses, the average Doppler factor $\delta \approx t_s \Gamma/t$ and the average flux decays as $t^{-(2+\alpha)} = t^{-2.5}$ in accordance with analytical estimates (Fenimore *et al* 1998). Dashed line, expected afterglow signal rising $\propto t^2$, peaking at ~ 100 s and falling off $\propto t^{-1.5}$ with arbitrary normalization (Lyutikov 2006).

show variability on a wide range of temporal scales, down to minutes, which are unrelated to the timescale of 22 years of magnetic field generation in the tachocline.)

• Prompt and afterglow polarization. Claims of high polarization (Coburn and Boggs 2003, Willis *et al* 2005) if confirmed, may provide a decisive test of GRB models (see, however, Rutledge and Fox 2003, Wigger *et al* 2004). The best way to produce polarization in the range $10\% \leq \Pi \leq 60\%$ is through synchrotron emission in large-scale magnetic fields (Lyutikov *et al* 2003). (Larger polarization can only be produced with inverse Compton mechanism, whereas smaller polarization can be produced by small-scale magnetic fields.)

Large-scale field structure in the ejecta emission may also be related to polarization of afterglows if fields from the magnetic shell are mixed in with the shocked circumstellar material. In this case, *the position angle should not change through the afterglow* while if polarization is observed both in prompt and afterglow emission the position angle should be the same. Also, polarization should be most independent of the 'jet break' moment.

11. Conclusion

In this paper, the underlying assumptions for the 'electromagnetic hypothesis' for ultra-relativistic GRB outflows have been outlined. The most striking implications of the electromagnetic hypothesis are that particle acceleration in the sources is due to direct dissipation of electromagnetic energy rather than shocks and that the outflows are cold, electromagnetically dominated flows, at least until they become strongly dissipative.

One of the major drawback of the model is that magnetic dissipation and particle acceleration are very complicated processes, depending crucially on the kinetic and geometric properties of the plasma. This situation may be contrasted with the shock acceleration schemes, where a qualitatively correct result for the spectrum of accelerated particles, *a kinetic property*, can be obtained from simple *macroscopic* considerations (jump conditions). The example of the Solar corona shows that despite being complicated, magnetic dissipation is an effective mean of particle acceleration.

Possible observational tests of the hypothesis have been discussed. In particular, interpretation of early afterglow features as being due to prompt emission seen at large angles, $\theta \ge 1/\Gamma$, allows us to measure radius at which prompt emission has been produced. Large prompt emission radii, $\sim 6 \times 10^{15}$ cm, seem to be inconsistent with the FBM, but close to the prediction of the EMM. Internal relativistic motion of 'fundamental emitters' assumed within the EMM may also explain x-ray flares during early afterglow phases (without the need for long source activity). An important implication of the EMM is that supernova explosions may be magnetically driven as well (Bisnovatyi-Kogan 1971, Leblanc and Wilson 1970, Proga *et al* 2003, Wheeler *et al* 2005).

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Appendix A. Applicability of the fluid approach for a blast wave

In case of extremely high Lorentz factors of the ejecta (which require even higher values of σ than were assumed in this paper), the fluid approximation for interaction of magnetized ejecta with ISM may break down. Consider an interface between the ejecta and the surrounding medium in its rest frame. As a particle from the surrounding medium enters the ejecta, it starts gyrating in the magnetic field. If a fraction $\sigma/(\sigma + 1)$ of the source luminosity L is in the form of magnetic field, then the turn angle in the rest frame in one dynamical time is

$$\omega'_B t_{\exp} \sim \sqrt{\frac{\sigma}{\sigma+1}} \frac{2e\sqrt{\pi L}}{c^{5/2}m_p\Gamma^3}.$$
(A.1)

In order to justify the fluid approximation, this should be larger than unity, which requires

$$\Gamma \leqslant \left(\frac{\sigma}{\sigma+1}\right)^{1/4} \left(\frac{4\pi e^2 L}{c^5 m_p^2}\right)^{1/6} \sim 4 \times 10^4,\tag{A.2}$$

for $\sigma \ge 1$. Thus, for any $\Gamma \le 4 \times 10^4$, an ISM particle can complete a half turn on a timescale short if compared with the expansion timescale. In this case, in a laboratory frame, momentum of the ejecta will be given to the particles almost instantaneously. For larger Lorentz factors, the instantaneous hydrodynamical approximation is not applicable, but if particles are turned by an angle larger than $\sim 1/\Gamma$ (larger than $\sim \pi$ in the observer frame) they will still be carried with the flow. Since the rest-frame magnetic field goes as $\sim 1/(t\Gamma(t))$, approximately linearly with time (for constant Γ), the rotational phase of a particle increases only logarithmically, $\int \omega'_B dt' \propto \ln t$. Thus, it takes a very long time for a particle to complete one gyration and be expelled from the ejecta. In this case, the ejecta will be effectively loaded with ISM particles.

Finally, for very high Lorentz factors,

$$\Gamma \ge \left(\frac{\sigma}{\sigma+1}\right)^{1/4} \frac{\sqrt{e}L^{1/4}}{c^{5/4}\sqrt{m_{\rm p}}} \sim 8 \times 10^6,\tag{A.3}$$

a particle makes a turn of less than $1/\Gamma$ (in the ejecta frame) on a dynamical timescale. In this case, the ejecta just pass through ISM without much interaction and without slowing down.

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