Mon. Not. R. Astron. Soc. **382,** L72–L76 (2007)

doi:10.1111/j.1745-3933.2007.00392.x

Re-born fireballs in gamma-ray bursts

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Accepted 2007 September 10. Received 2007 September 4; in original form 2007 July 19

ABSTRACT

We consider the interaction between a relativistic fireball and material assumed to be still located just outside the progenitor star. Only a small fraction of the expected mass is sufficient to decelerate the fireball efficiently, leading to dissipation of most of its kinetic energy. Since the scattering optical depths are still large at distances comparable to the progenitor radius, the dissipated energy is trapped in the system, accelerating it to relativistic velocities. The process resembles the birth of another fireball at radii $R \sim 10^{11}$ cm, not far from the transparency radius, and with starting bulk Lorentz factors $\Gamma_c \sim 10$. As seen in the observer frame, this 're-generated' fireball appears collimated within an angle $\theta_j = 1/\Gamma_c$. If the central engine works intermittently, the funnel can, at least partially, refill and the process can repeat itself. We discuss how this idea can help to solve some open issues of the more conventional internal shock scenario for interpreting gamma-ray burst properties.

Key words: radiation mechanisms: general – gamma-rays: bursts – X-rays: general.

1 INTRODUCTION

The internal/external shock scenario (see e.g. Piran 2004; Mészáros 2006) is currently the leading model to explain the complex phenomenology of gamma-ray burst (GRB) prompt and afterglow emission. Despite the fact that it can account for many observed characteristics, there are a few open issues and difficulties that this model cannot solve, or can accommodate only with some important modifications.

Here we recall some problems of the standard scenario and mention some ideas already put forward to account for them.

(i) *Efficiency I: high* Γ*-contrast.* In internal shocks only the relative kinetic energy of the two colliding shells can be dissipated. Thus 'dynamical' efficiencies of only a few per cent can be achieved for colliding shells whose Lorentz factors Γ differ by a factor of order unity. Such efficiency has to include energy dissipated into randomizing protons, amplifying (or even generating) magnetic fields and accelerating emitting leptons. As the emitted radiation is produced only by the latter component, it corresponds to just a fraction of the dynamical efficiency. This problem, pointed out by Kumar (1999) among others, can be solved by postulating contrasts in Γ much exceeding 100 (Beloborodov 2000; Kobayashi & Sari 2001). In these cases the typical Lorentz factor of GRBs should thus largely exceed the 'canonical' value ∼100. In this case it is difficult to understand how the value of the peak energy of the prompt spectrum does not change wildly.

(ii) *Efficiency II: afterglow/prompt power ratio.* A related inconsistency concerns the observed ratios of the bolometric fluence originating in the afterglow to that in the prompt phase. Since external shocks are dynamically more efficient than internal ones, such a ratio is expected to exceed 1, contrary to current estimates. The problem has been exacerbated by recent observations by the *Swift* satellite, showing that the X-ray afterglow light curve seen after a few hours – thought to be smoothly connected with the end of the prompt phase – comprises a steep early phase. As a consequence, the total afterglow energy is less than postulated before. Willingale et al. (2007), parametrizing the behaviour of the *Swift* GRB X-ray light curves, derived an average X-ray afterglow-to-prompt fluence ratio of around 10 per cent (see also Zhang et al. 2007). Furthermore, the very same origin of the early X-ray radiation as produced by external shocks is questioned, since its behaviour is different from the optical one (e.g. Panaitescu et al. 2006). If the X-ray emission does not originate in the afterglow phase, this further reduces the above ratio.

(iii) *Spectral energy correlations.* Correlations have been found between (1) the energy where most of the prompt power is emitted ($E_{\rm peak}$) and the isotropic prompt bolometric energy $E_{\gamma,\rm iso}$ (Amati relation, Amati et al. 2002; Amati 2006), and (2) $E_{\rm peak}$ and the collimation-corrected energy E_γ (Ghirlanda relation). The slope of the former correlation is $E_{\rm peak} \propto E_{\gamma,\rm iso}^{1/2}$ while the slope of the latter depends on the radial profile of the circumburst density. For a homogeneous-density medium $E_{\rm peak} \propto E_\gamma^{0.7}$ (Ghirlanda, Ghisellini & Lazzati 2004), while for a wind-like profile in density ($\propto r^{-2}$) the correlation is linear: $E_{\rm peak} \propto E_\gamma$ (Nava et al. 2006; Ghirlanda et al. 2007a). If linear, the relation is Lorentz-invariant and indicates that

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different GRBs roughly emit the same number of photons at the peak (i.e. $E_\gamma/E_{\rm peak} \sim$ constant).

Note that the derivation of E_γ requires not only information on the jet break time $t_{\rm j}$, but also a model relating $t_{\rm j}$ with the collimation angle $\theta_{\rm j}$, which in turn depends on the circumburst density value, the profile and the radiative efficiency η (i.e. $E_\gamma = \eta E_{\rm kin}$). The phenomenological connection between the three observables $E_{\gamma,\rm iso}, E_{\rm peak}$ and $t_{\rm j}$, as found by Liang & Zhang (2005), is instead model-independent. It is of the form $E_{\rm iso} \propto E^a_{\rm peak}\, t_{\rm j}^{-b}$, which for $b \sim -1$ is consistent with the Ghirlanda relation (in both the homogeneous and wind cases; see Nava et al. 2006). A further tight phenomenological relation appears to link three prompt emission quantities: the isotropic peak luminosity $L_{\rm iso}, E_{\rm peak}$ and the time interval $T_{0.45}$ during which the emission is above a certain level (Firmani et al. 2006). All these correlations were not predicted by the internal/external shock scenario, and can only be reconciled with it as long as specific dependences of the bulk Lorentz factor upon $E_{\rm iso}$ are satisfied (see table 1 in Zhang & Mészáros 2002).

The above issues motivate the search for alternatives or for substantial modifications of the standard model. The efficiency problem and the existence of the spectral energy relations prompted Thompson (2006, T06 hereafter) and Thompson, Mészáros & Rees (2007, T07 hereafter) to suggest that, besides internal shocks, dissipation might also occur because of the interaction between the fireball and the walls of the funnel in the star through which it propagates (see Section 2). This hypothesis also introduces a typical scale to the problem, namely the radius of the progenitor star ($R_* \sim 10^{10}$–10^{11} cm): shear instabilities within R_* can reconvert a significant fraction of bulk kinetic energy into heat. Since this dissipation occurs up to R_* (i.e. not far from the transparency radius), the increased internal energy can only partially reconvert into bulk motion via adiabatic expansion, increasing the efficiency. Similarly, studies of magnetized fireballs (e.g. Drenkhahn & Spruit 2002; Giannios & Spruit 2007) have shown that dissipation of magnetic energy through reconnection can also contribute to increase the radiation content of the fireball at relatively large radii.

Along the above-mentioned lines, in this Letter we propose a further possible way in which a large fraction of the fireball bulk energy can be dissipated at distances $R \sim R_*$. This assumes that at these distances the fireball collides with some mass which is (nearly) at rest: a small fraction of the mass swept up in the funnel left along the fireball propagation axis is sufficient to lead to efficient dissipation. Hereafter such mass will be referred to as 'IDM' ('intervening debris of the cocoon material'). The collision of an expanding pair-electromagnetic pulse with a shell of baryonic matter has been investigated by Ruffini et al. (2000), but only in a spherically symmetric case.

Our treatment of the collision is simplified, in order to allow an analytical and simple description. We assume the IDM to be at rest and homogeneous in density and the fireball to have a Lorentz factor $\Gamma \gg 1$. The interaction is described in 'steps', while in reality it will be continuous in time. A complete treatment of the dynamics and emission properties of our model requires numerical simulations [of the kind presented by Morsony, Lazzati & Begelman (2007), introducing some erratic behaviour of the injected jet energy]. Interestingly, the model predicts that jet properties depend on the polar angle (like in a structured jet, see Rossi, Lazzati & Rees 2002), and it naturally implies a connection among the observed spectral energy correlations.

2 SHEAR-DRIVEN INSTABILITIES AND DISSIPATION OF BULK KINETIC ENERGY

As mentioned, T06 and T07 proposed a model in which the efficiency of dissipation of kinetic energy into radiation is enhanced with respect to the internal shock scenario and the spectral energy correlations, in particular the Amati one, can be accounted for. At the same time, in their scenario synchrotron emission could play a minor role, the radiation field being dominated by thermalized high-energy photons or by the inverse Compton process.

For what follows it is useful to summarize their main arguments here. Consider a fireball that at some distance $R_0 \lesssim R_*$ from the central engine is moving relativistically with a bulk Lorentz factor Γ_0. The fireball is initially propagating inside the funnel of the progenitor star. T06 and T07 assume that a large fraction of the energy dissipated at R_0 is thermalized into blackbody radiation of luminosity

$$L_{\rm BB,iso} = 4\pi R_0^2 \Gamma_0^2 \sigma T_0'^4 = 4\pi \frac{R_0^2}{\Gamma_0^2} \sigma T_0^4, \quad (1)$$

where T_0' and $T_0 = \Gamma_0 T_0'$ are the temperatures at R_0 in the comoving and observing frame, respectively. The collimation-corrected luminosity is $L_{\rm BB} = (1 - \cos\theta_{\rm j}) L_{\rm BB,iso}$ which, for small semi-aperture angles $\theta_{\rm j}$ of the jetted fireball (assumed conical), gives

$$\theta_{\rm j}^2 \sim \frac{2L_{\rm BB}}{L_{\rm BB,iso}}. \quad (2)$$

A key assumption of the model is that $\Gamma_0 \sim 1/\theta_{\rm j}$. The argument behind it is that if $\Gamma_0 \gg 1/\theta_{\rm j}$, shear-driven instabilities do not have time to grow (in the comoving frame), while in the opposite case the flow mixes easily with the heavier material and decelerates to $\Gamma \sim 1$. Then, assuming $\Gamma_0 = 1/\theta_{\rm j}$ and substituting it in equation (1),

$$L_{\rm BB,iso} \sim 8\pi R_0^2 \frac{L_{\rm BB}}{L_{\rm BB,iso}} \sigma {T_0}^4. \quad (3)$$

Setting $E_{\rm BB,iso} = L_{\rm BB,iso}\, t_{\rm burst}$ and $E_{\rm BB} = L_{\rm BB} t_{\rm burst}$, where $t_{\rm burst}$ is the duration of the prompt emission, gives

$$E_{\rm peak} \propto T_0 \propto E_{\rm BB,iso}^{1/2} E_{\rm BB}^{-1/4} t_{\rm burst}^{-1/4}. \quad (4)$$

This corresponds to the Amati relation if $E_{\rm BB}$ is similar in different bursts and the dispersion of GRB durations is also limited. It should be noted that a relation similar to the Amati one can be also recovered by adopting $E_{\rm BB} \propto E^a_{\rm peak}$, as suggested by the Ghirlanda relation. For instance, for $a = 1$ (wind case),

$$E_{\rm peak} \propto E_{\rm BB,iso}^{2/5} t_{\rm burst}^{-1/5}. \quad (5)$$

For the derivation of equations (4) and (5) a key assumption is the dependence on temperature of the blackbody law, which leads to both a slope *and* a normalization similar to those characterizing the Amati relation.

We can ask what happens if, instead of a blackbody, one assumes that the spectrum is a cut-off power law. This question is particularly relevant since the burst spectrum is rarely described by a pure blackbody (even if some bursts are, see Ghirlanda, Celotti & Ghisellini 2003), and also the blackbody plus power-law model (Ryde 2005) faces severe problems, even considering time-resolved spectra (Ghirlanda et al. 2007b).

How is the above derivation modified if, instead of blackbody emission, the spectrum is best described by a cut-off power law? Consider then a spectrum described in the comoving frame by $L'_{\gamma,\rm iso}(E') \propto E'^{-\beta} \exp(-E'/E'_0)$ and approximate the observed isotropic bolometric luminosity as

$$L_{\rm iso} \propto \Gamma_0^2 \left(\frac{E_{\rm peak}}{\Gamma_0}\right)^{1-\beta} \propto \left(\frac{L_\gamma}{L_{\gamma,\rm iso}}\right)^{-(1+\beta)/2} E_{\rm peak}^{1-\beta}. \quad (6)$$

This leads to

$$E_{\rm peak} \propto E_{\rm iso}^{1/2} E_{\gamma}^{(1+\beta)/(2-2\beta)} t_{\rm burst}^{-1/(1-\beta)}. \qquad (7)$$

Therefore the dependence of the Amati relation can be recovered even for cut-off power-law spectra, but the normalization in this case is not determined. Note also that for $\beta = -3$ (Wien spectrum), equations (7) and (4) have the same dependences.

3 FIREBALL–IDM COLLISION

To excavate a funnel inside a progenitor star of mass $M_* = 10 M_{*,1}$ solar masses, the 'proto-jet' has to push out the mass that did not fall into the newly born black hole. This means a fraction of $(1 - \cos\theta_{\rm f})$ $M_* = 0.1\, M_{*,1}\theta^2_{\rm f,-1}$, where $\theta_{\rm f} = 0.1\theta_{\rm f,-1}$ is the funnel opening angle. This mass expands sideways as the proto-jet breaks out at the surface of the progenitor, forming a cocoon (see also Ramirez-Ruiz, Celotti & Rees 2002). We must expect, however, that after the break-out the region in front of the funnel will not be perfectly cleared of mass. To be negligible, the mass $M_{\rm c}$ left as IDM should be $\ll E_0/(\Gamma_0\, c^2) \sim 5 \times 10^{-7}\,{\rm M}_{\odot}$: less than one in a million particles should remain there.

If the jetted fireball is not continuous, this mass may be still there at the moment of arrival of the new fireball pulse. Also, even if $M_{\rm c}$ is 10^{-3}–10^{-4} the excavated mass, the IDM can have important dynamical effects on it. As the bulk velocity and energy content of the IDM can be neglected in comparison with those of the coming fireball, the IDM will be approximated as initially at rest and cold.

An interesting aspect of this scenario concerns multi-peaked bursts, especially when pulses in the prompt emission are separated by quiescent periods. If the central engine works at a reduced rate during quiescence, material from the walls of the funnel, previously in pressure equilibrium with the jet, will tend to refill the funnel again (see also Wang & Mészáros 2007). This requires a time $T_{\rm f} \sim \theta_{\rm f} R/(\beta_{\rm s} c) \sim 3 \times 10^{-4} (R/R_{0,7})\theta_{\rm f,-1}/\beta_{\rm s,-1}$ s, where $\beta_{\rm s} c$ is the sound speed, and $R_0 = 10^7 R_{0,7}$ cm is the radius at the base of the funnel. Then the amount of mass that should be pushed out again depends upon the quiescent time, but after $\sim$1 s, the funnel is completely closed, and the process repeats itself.

3.1 Results: dynamics and dissipation

In the following we derive the main characteristics of the system formed by a fireball impacting against the IDM. Consider a fireball with energy E_0, mass M_0 and bulk Lorentz factor $\Gamma_0 = E_0/(M_0 c^2)$, impacting against a mass $M_{\rm c}$, initially at rest.

The energy and momentum conservation laws read

$$M_0\Gamma_0 + M_{\rm c} = \Gamma_{\rm c}(M_0 + M_{\rm c} + \epsilon'/c^2),$$
$$M_0\Gamma_0\beta_0 = \Gamma_{\rm c}\beta_{\rm c}(M_0 + M_{\rm c} + \epsilon'/c^2), \qquad (8)$$

where ϵ' is the dissipated energy measured in the frame moving at $\beta_{\rm c} c$. Solving for $\beta_{\rm c}$ and ϵ' gives

$$\beta_{\rm c} = \beta_0 \frac{M_0\Gamma_0}{M_0\Gamma_0 + M_{\rm c}} \equiv \frac{\beta_0}{1+x}; \quad x \equiv \frac{M_{\rm c} c^2}{M_0\Gamma_0 c^2}, \qquad (9)$$

$$\epsilon' = \frac{E_0}{\Gamma_{\rm c}}\left(1 + x - x\Gamma_{\rm c} - \frac{\Gamma_{\rm c}}{\Gamma_0}\right), \qquad (10)$$

$$\epsilon \equiv \Gamma_{\rm c}\epsilon'. \qquad (11)$$

The Lorentz factor $\Gamma_{\rm c}$ is shown in Fig. 1 as a function of $M_{\rm c}$ for four values of E_0 (top panel) and as a function of E_0 for three values of $M_{\rm c}$ (bottom panel). One can see that for the process to be interesting (i.e. $\Gamma_{\rm c}$ significantly smaller than Γ_0) and to avoid 'over-loading' of baryons (too small $\Gamma_{\rm c}$), $M_{\rm c}$ is required to be in specific ranges, which depend on E_0. These ranges, however, encompass almost two orders of magnitude.

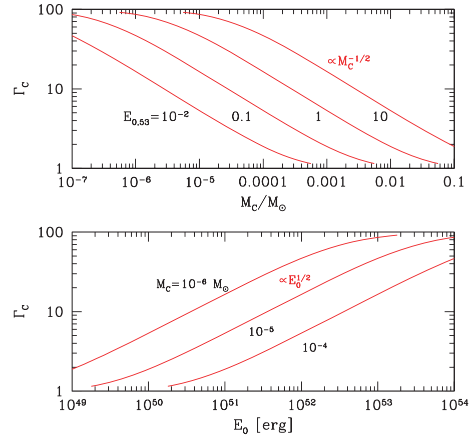

Figure 1. The Lorentz factor $\Gamma_{\rm c}$ of the IDM + fireball system as a function of $M_{\rm c}$ for different values of E_0 (top panel) and as a function of E_0 for selected values of $M_{\rm c}$ (bottom panel).

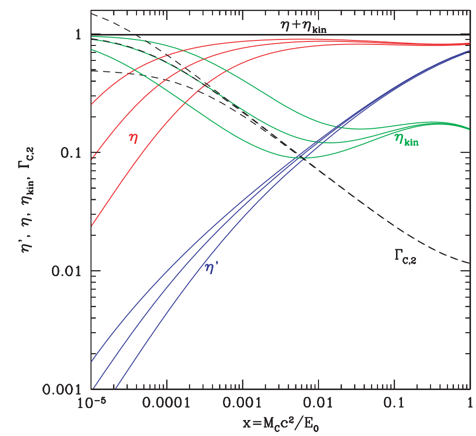

Figure 2. The Lorentz factor $\Gamma_{\rm c,2} \equiv \Gamma_{\rm c}/100$, the dynamical efficiency η' and η (in the comoving and observer frame, respectively) as a function of x. Each quantity is calculated for $\Gamma_0 = 50$, 100 and 200 (from top to bottom).

$\Gamma_{\rm c}$ as a function of x is reported in Fig. 2. The power-law dependence $\Gamma_{\rm c} \propto x^{-1/2}$ can be derived directly from equation (9), since $\Gamma_{\rm c}^{-2} = 1 - \beta_{\rm c}^2 \propto 2x$ for $x \ll 1$ and $\beta_0 \to 1$. $\Gamma_{\rm c}$ is limited to $\sim$10 even for $x \sim 4 \times 10^{-3}$, corresponding to $M_{\rm c} = 2.2 \times 10^{-5} E_{0,52}\,{\rm M}_{\odot}$.

In Fig. 2 we show $\eta' \equiv \epsilon'/E_0$ and $\eta \equiv \epsilon/E_0$ as a function of x. For clarity also the fraction of the initial kinetic energy preserved after the collision

$$\eta_{\rm kin} \equiv \frac{(\Gamma_{\rm c} - 1)(M_0 + M_{\rm c})c^2}{E_0} = (\Gamma_{\rm c} - 1)\left(x + \frac{1}{\Gamma_0}\right) \qquad (12)$$

is plotted. The sum $\eta + \eta_{\rm kin}$ is unity by definition.

Note that since x is a ratio, $M_{\rm c}$ and E_0 can be taken as either 'isotropic' or 'real' (collimation-corrected) values.

3.2 Evolution of the fireball + IDM system

After the collision/dissipation phase the fireball + IDM is expected to be optically thick: the Thomson scattering optical depth of the IDM material is

$$\tau_c = \sigma_c n \Delta R = \frac{\sigma_T M_c}{4\pi R_c^2 m_p} = 3.2 \times 10^3 \frac{M_{c,-6}}{R_{c,11}^2}, \quad (13)$$

where $M_c = 10^{-6} M_{c,-6}$ solar masses, and $R_c = 10^{11} R_{c,11}$ cm. The optical depth of the fireball just before the collision is of order

$$\tau_0 = \frac{\sigma_T E_0}{4\pi R^2 \Gamma_0 m_p c^2} = 1.8 \times 10^5 \frac{E_{0,52}}{R_{c,11}^2}, \quad (14)$$

where M_c and $E_0 = 10^{52} E_{0,52}$ erg are here isotropic quantities. These large optical depths imply that the radiation produced following the collision is trapped inside the fireball + IDM system, which will expand because of the internal pressure. In the frame moving with Γ_c, the expansion is isotropic. In this frame some final Γ' will be reached. As seen in the observer frame, the expansion is highly asymmetric, and the geometry of the system resembles a cone, with semi-aperture angle given by (Barbiellini, Celotti & Longo 2003)

$$\tan\theta_j = \frac{\beta_\perp}{\beta_\parallel} = \frac{\beta' \sin\theta'}{\Gamma_c(\beta'\cos\theta' + \beta_c)} \rightarrow \theta_j \sim \frac{1}{\Gamma_c}, \quad (15)$$

where the last equality assumed $\theta' = 90°$ ($\beta' \sim 1$ and $\beta_c \sim 1$). Therefore the aperture angle of the re-born fireball is related to Γ_c, independently of the initial Γ-factor of the fireball *before* the collision or *after* the expansion. Note that the initial aperture angle of the fireball is irrelevant, as it is the aperture angle of the funnel of the progenitor star.

The fireball is not collimated in a perfect cone, and mass and energy propagate also outside θ_j. Since in the frame moving with Γ_c it expands isotropically, $M'(\Omega') = M'/(4\pi)$ is approximately constant. Therefore in the observer frame

$$M(\theta) = \frac{M'}{4\pi}\frac{\mathrm{d}\cos\theta'}{\mathrm{d}\cos\theta}. \quad (16)$$

Also the resulting Lorentz factor Γ is angle-dependent: from the relativistic composition of velocity (e.g. Rybicki & Lightman 1979),

$$\beta_\parallel = \beta\cos\theta = \frac{\beta'\cos\theta' + \beta_c}{(1 + \beta_c\beta'\cos\theta')},$$
$$\beta_\perp = \beta\sin\theta = \frac{\beta'\sin\theta'}{\Gamma_c(1 + \beta_c\beta'\cos\theta')},$$
$$\beta(\theta) = \left(\beta_\parallel^2 + \beta_\perp^2\right)^{1/2},$$
$$\Gamma(\theta) = \left[1 - \beta^2(\theta)\right]^{-1/2} = \Gamma'\Gamma_c(1 + \beta'\beta_c\cos\theta'). \quad (17)$$

The observed Γ-factor is constant up to angles slightly smaller than $1/\Gamma_c$, and decreases as θ^{-2} above.

As a consequence of the angular dependence of mass and bulk Lorentz factor, also the energy depends on θ as

$$E(\theta) = \Gamma(\theta)M(\theta)c^2. \quad (18)$$

Such dependences of mass, Γ and E on polar angle are illustrated in Fig. 3. The jet is structured and well approximated by a top-hat jet: the energy profile is nearly constant within an angle slightly smaller than $1/\Gamma_c$, and at larger angles decreases approximately as a steep power law $E(\theta) \propto \theta^{-11/2}$, since $M(\theta) \propto \theta^{-7/2}$ and $\Gamma(\theta) \propto \theta^{-2}$ for $\theta \gg 1/\Gamma_c$. This particular behaviour gives rise to an afterglow light curve indistinguishable from a top-hat jet (see Rossi et al. 2004).

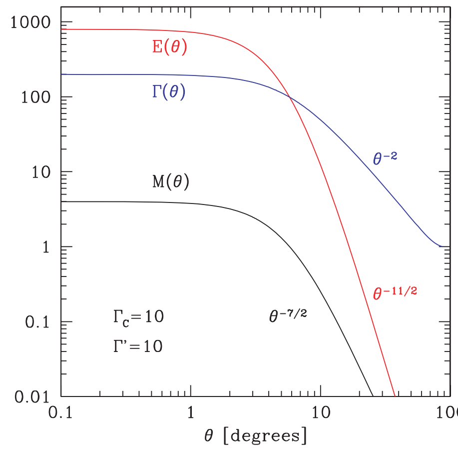

Figure 3. The mass, energy and bulk Lorentz factor (observer frame) as functions of the angle from the jet axis θ, for $\Gamma_c = 10$ and $\Gamma' = 10$. All the three quantities are constant up to $\theta \simeq 1/\Gamma_c$ and then decrease (approximately) as power laws.

4 DISCUSSION AND CONCLUSIONS

The most appealing feature of the proposed model is the high efficiency in re-converting the fireball kinetic energy into internal pressure at a radius comparable to the radius of the progenitor star, i.e. on a scale not far from the transparency one. Also the observed energetics of the internal radiation will be large, since the system becomes transparent during (or slightly after) the expansion/acceleration phase, similarly to the standard fireball models of initially high entropy, where the fireball becomes transparent before coasting. Our model therefore increases the parameter space of high-efficiency regimes. Our model is also similar to the model proposed by T06 and T07 and to those in which the dissipation of an energetically important magnetic field occurs at large radii (see e.g. Giannios & Spruit 2007). The efficiency of the energy re-conversion for the fireball–IDM collision is of the order of 50–80 per cent for large ranges in the mass of the IDM and energy of the fireball (see Fig. 2).

With respect to the idea proposed by T07, summarized in Section 2, in our scenario there is no requirement on any specific value for the fireball bulk Lorentz factor Γ_0 prior to its collision with the IDM. Note also that in the T07 model, a 'standard' fireball (i.e. not magnetic) moving with $\Gamma_0 \sim 1/\theta_j$ can dissipate part of its kinetic energy, but it cannot reach final Γ-factors larger than Γ_0, which is bound to be small for typical θ_j. In our case the final Γ can instead be large (even if always smaller than Γ_0). The jet angle is $\sim 1/\Gamma_c$ even for large values of the initial Γ_0 and final Γ. This is a result of our model, and not an assumption.

The re-born fireball is structured and the $M(\theta)$ and $\Gamma(\theta)$ behaviours imply that $E(\theta)$ depends on θ as steep power laws. Despite the angle dependence of the energy, the jet should produce an afterglow indistinguishable from a top-hat jet. Clearly our description of the fireball–IDM interaction is extremely simplified, aimed at building a physical intuition based on the analytical treatment. More realistic situations should be studied via numerical simulations, but we would like to comment on two aspects. (i) Even if the IDM were initially at rest, as soon as the fireball started depositing a fraction of energy and momentum, the IDM would begin to move. The whole process would take long enough that towards the end it would be probably better described as the interaction with a moving IDM, with a consequent loss of efficiency. (ii) In a 'continuous' (non-intermittent) scenario, the IDM would predominantly interact with the fireball edge, causing only a partial dissipation of its energy and the formation of a 'fast spine–slow layer' structure. In this case the determination of the relevant jet opening angle (i.e. within which most of the energy is concentrated) requires a more accurate numerical treatment.

Despite these caveats, it is still interesting to consider whether the model can account for the spectral energy correlations, or at least highlight the relations between them. Consider the Ghirlanda correlation in the wind case, $E_{\rm peak} \propto E_\gamma$, which can be rewritten as $E_{\rm peak} \propto \theta_{\rm j}^2 E_{\gamma,\rm iso}$ for small $\theta_{\rm j}$. The requirement that also $E_{\rm peak} \propto E_{\gamma,\rm iso}^{1/2}$ (Amati relation) leads to

$$E_\gamma \theta_{\rm j}^2 = \text{constant}, \qquad (19)$$

i.e. more energetic bursts are more collimated, as predicted in our model. The above condition (equation 19) can be quantitatively satisfied if (i) the mass against which the fireball collides is similar in different GRBs, namely $\Gamma_{\rm c} \sim 1/\theta_{\rm j} \sim E_0^{1/2}$ (see Fig. 1), and (ii) the prompt emission luminosity E_γ is also a constant fraction of E_0 for different GRBs. Equation (19) does not explain the Amati or Ghirlanda relation, although it offers some physical meaning to the required connection between the two. Note that, in the standard internal shock model, one can recover the Amati relation if $\Gamma \sim$ constant (see Zhang & Mészáros 2002).

In this Letter we have not discussed the characteristics of the spectrum predicted in our scenario (Nava et al., in preparation). In general terms, the most effective radiation process would be 'dynamical Compton', or Fermi 'acceleration' of photons, as discussed by Gruzinov & Mészáros (2000) in the context of internal shocks. The large number of photons per proton in the fireball (corresponding to the 'fossil' radiation that has accelerated the fireball itself in the first place) implies that a significant fraction of the internal energy following the fireball–IDM collision *directly* energizes photons, i.e. photons amplify their energy by interacting with leptons with bulk momentum not yet randomized in the shock. The process is analogous to particle acceleration in shocks and gives rise to high-energy photons, conserving their number.

Since this occurs at $R \sim 10^{11}$ cm, i.e. slightly above the progenitor star, the re-born fireball will become transparent during or just after the acceleration phase (at a transparency radius $R_\tau \sim 3 \times 10^{12}$ cm, equation 13). This ensures that photons do not have time to lose energy via adiabatic expansion, once again leading to large efficiencies.

The time-scale for the refilling of the funnel during quiescent phases of the central engine can be short enough that the process repeats itself. The estimated refilling time-scale is of the order of a second; bursts with shorter or longer 'quiescent' phases should then show different properties. If the fireball–IDM collision is the dominant process for the dissipation (and emission), more energetic spikes are expected to follow longer quiescent phases (see Ramirez-Ruiz & Merloni 2001).

ACKNOWLEDGMENTS

A 2005 PRIN–INAF grant is acknowledged for partial funding. We thank the anonymous referee for useful criticism.

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