Implications of the early X-ray afterglow light curves of *Swift* gamma-ray bursts

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ABSTRACT

According to current models, gamma-ray bursts (GRBs) are produced when the energy carried by a relativistic outflow is dissipated and converted into radiation. The efficiency of this process, ϵ_γ, is one of the critical factors in any GRB model. The X-ray afterglow light curves of *Swift* GRBs show an early stage of flattish decay. This has been interpreted as reflecting energy injection. When combined with previous estimates, which have concluded that the kinetic energy of the late ($\gtrsim 10$ h) afterglow is comparable to the energy emitted in γ-rays, this interpretation implies very high values of ϵ_γ, corresponding to $\gtrsim 90$ per cent of the initial energy being converted into γ-rays. Such a high efficiency is hard to reconcile with most models, including in particular the popular internal-shocks model. We re-analyse the derivation of the kinetic energy from the afterglow X-ray flux and re-examine the resulting estimates of the efficiency. We confirm that, if the flattish decay arises from energy injection and the pre-*Swift* broad-band estimates of the kinetic energy are correct, then $\epsilon_\gamma \gtrsim 0.9$. We discuss various issues related to this result, including an alternative interpretation of the light curve in terms of a two-component outflow model, which we apply to the X-ray observations of GRB 050315. We point out, however, that another interpretation of the flattish decay – a variable X-ray afterglow efficiency (e.g. due to a time dependence of afterglow shock microphysical parameters) – is possible. We also show that direct estimates of the kinetic energy from the late X-ray afterglow flux are sensitive to the assumed values of the shock microphysical parameters and suggest that broad-band afterglow fits might have underestimated the kinetic energy (e.g. by overestimating the fraction of electrons that are accelerated to relativistic energies). Either one of these possibilities implies a lower γ-ray efficiency, and their joint effect could conceivably reduce the estimate of the typical ϵ_γ to a value in the range ~ 0.1–0.5.

Key words: radiation mechanisms: non-thermal – ISM: jets and outflows – gamma-rays: bursts – X-rays: individual: GRB 050315.

1 INTRODUCTION

Recent observations by the *Swift* X-ray telescope have provided new information on the early behaviour of the X-ray light curve of long-duration ($\gtrsim 2$ s) gamma-ray burst (GRB) sources. Specifically, it was found (Nousek et al. 2006) that the light curves of these sources have a generic shape consisting of three distinct power-law segments $\propto t^{-\alpha}$: an initial (at $t < t_{break,1}$, with $300\,\text{s} \lesssim t_{break,1} \lesssim 500\,\text{s}$) very steep decline with time t (with a power-law index α_1 in the range $3 \lesssim \alpha_1 \lesssim 5$; see also Bartherlmy et al. 2005; Tagliaferri et al. 2005); a subsequent (at $t_{break,1} < t < t_{break,2}$, with $10^3\,\text{s} \lesssim t_{break,2} \lesssim 10^4\,\text{s}$) very shallow decay ($0.2 \lesssim \alpha_2 \lesssim 0.8$); and a final steepening (at $t > t_{break,2}$) to the canonical power-law behaviour ($1 \lesssim \alpha_3 \lesssim 1.5$) that was known from pre-*Swift* observations.

Nousek et al. (2006) already recognized that these results have direct consequences to the question of the γ-ray emission efficiency in GRB sources. This question is important to our understanding of the basic prompt-emission mechanism. In the currently accepted interpretation (e.g. Piran 1999, 2004), the γ-rays originate in a relativistic jet that is launched from the vicinity of a newly born neutron star or stellar mass black hole. In the simplest picture, a fraction ϵ_γ of the energy injected at the source is given to the prompt radiation, with the remaining

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fraction $(1 - \epsilon_\gamma)$ ending up as kinetic energy of ambient gas that is swept up by a forward shock and then mostly radiated as early-afterglow emission. One attractive mechanism for explaining the prompt-emission characteristics invokes internal shocks that are generated when the ejecta have a non-uniform distribution of Lorentz factors, which results in outflowing 'shells' colliding with each other at large distances from the source.

Pre-*Swift* observations, based on the measurements of the γ-ray fluence and of the 'late' ($\gtrsim 10\,\mathrm{h}$) afterglow emission, have implied (when interpreted in the context of the basic jet model) comparable (and narrowly clustered) values for the radiated γ-ray energy and the kinetic energy feeding the afterglow emission, i.e. $\epsilon_\gamma \sim 0.5$ (e.g. Frail et al. 2001; Panaitescu & Kumar 2001a,b, 2002; Berger, Kulkarni & Frail 2003; Bloom, Frail & Kulkarni 2003; Yost et al. 2003). This result is seemingly problematic for the internal-shocks model, for which an order-of-magnitude lower value for ϵ_γ is a more natural expectation (Kobayashi, Piran & Sari 1997; Daigne & Mochkovitch 1998; Kumar 1999; Guetta, Spada & Waxman 2001). Eichler & Jontof-Hutter (2005) inferred an even higher efficiency in the context of off-beam viewing angles with respect to the gamma-ray emission region. This apparent difficulty could in principle be overcome if the ejected shells have a highly non-uniform distribution of Lorentz factors (e.g. Beloborodov 2000; Kobayashi & Sari 2001). However, to fit the data the shells must also satisfy a number of other restrictive conditions, which reduces the attractiveness of this interpretation (see Section 5). An alternative proposal was made by Peng, Königl & Granot (2005), who argued that if the jet consists of an ultrarelativistic narrow component (from which the prompt emission originates) and a moderately relativistic wide component, with the latter having a higher kinetic energy and the former a higher kinetic energy per unit solid angle, then the wide component would dominate the late-afterglow emission and the γ-ray radiative efficiency of the narrow component could be significantly lower than the value of ϵ_γ inferred under the assumption of a single-component jet. As explained in Peng et al. (2005), this proposal was motivated by observational indications of the presence of two components in the late-afterglow light curves of several GRB sources and by the predictions of certain GRB source models.

The earliest (steepest) segment of the afterglow light curve is most naturally explained as radiation at large angles to our line of sight corresponding either to the prompt emission (Kumar & Panaitescu 2000) or to emission from the reverse shock that is driven into the ejecta (Kobayashi et al. 2005). This implies that the early-afterglow emission is much weaker than what would be expected on the basis of an extrapolation from the late-afterglow data. This behaviour was interpreted by Nousek et al. (2006) as an indication of an even higher γ-ray emission efficiency, typically $\epsilon_\gamma \sim 0.9$. Such a high efficiency could render the internal-shocks model untenable.

Our primary goal is to evaluate the prompt-emission efficiency as accurately and systematically as possible on the basis of current data. For this purpose, we re-derive in Section 2.1 expressions that explicitly relate ϵ_γ to observable quantities. In particular, we express ϵ_γ in terms of the product κf of two parameters, one (κ) encapsulating information that could be obtained by pre-*Swift* observations, and the other (f) representing early-time data obtained in the *Swift* measurements. In Section 2.2, we re-examine the estimates of the kinetic energy during the afterglow phase as inferred from the X-ray flux and present a new general formulation, correcting errors that have propagated in the literature and have generally led to an underestimate of the kinetic energy. We then analyse both pre-*Swift* (Section 2.3) and *Swift* (Section 2.4) data in a uniform manner in the context of this formalism. We argue that the kinetic energy estimates remain subject to considerable uncertainties. In particular, while the simple analysis of a large number of bursts (using 'typical' values for the microphysical parameters) yields rather high values for the kinetic energy and hence a low inferred γ-ray efficiency, a multiwavelength analysis (which determines the microphysical parameters from the fit to the data) of a small subset suggests that the kinetic energy is lower and hence the inferred value of ϵ_γ is higher. Some caveats to this analysis are considered in Section 2.5.

The suggestion that the *Swift* observations imply a higher value of ϵ_γ in comparison with the pre-*Swift* results is based on the interpretation of the flattish segment of the X-ray light curve as reflecting an increase in the kinetic energy of the forward shock during the early stages of the afterglow. In this picture, the kinetic energy just after the prompt-emission phase was significantly lower than the kinetic energy estimated from the later stages of the afterglow (the pre-*Swift* results). One conceivable way of avoiding the need to increase the estimate of ϵ_γ in light of the new *Swift* observations is through a time evolution (specifically, an increase with t) of the X-ray afterglow emission efficiency ϵ_X. Such a behaviour could in principle also account for the flattish segment of the light curve and eliminate the need to invoke an increase in the shock kinetic energy. If, in addition, the value of the afterglow kinetic energy at late times were underestimated by the broad-band fits to pre-*Swift* GRB afterglows, which could be the case if only a fraction $\xi_e < 1$ of the electrons in the afterglow shock were accelerated to relativistic energies with a power-law distribution (see Eichler & Waxman 2005), then the typical afterglow efficiency would be further reduced (to a value as low as $\epsilon_\gamma \sim 0.1$ if $\xi_e \sim 0.1$), which might reconcile the new *Swift* data with the comparatively low efficiencies expected in the internal-shocks model. We discuss these issues in Section 3.

A two-component jet model with the characteristics required for reducing the inferred γ-ray emission efficiency is evidently disfavoured by the *Swift* data, but a model of this type with different parameters could provide an alternative interpretation of the flattish shape of the light curve between $t_{\mathrm{break},1}$ and $t_{\mathrm{break},2}$. We elaborate on these matters in Section 4, where we also present a tentative fit to the X-ray light curve of the *Swift* source GRB 050315 in the context of this scenario.

Our conclusions on the physical implications of the early X-ray light-curve observations of GRB sources are presented in Section 5.

2 ESTIMATING THE γ-RAY EFFICIENCY

The observed isotropic-equivalent luminosity can generally be expressed as $L_{\mathrm{iso}} = \epsilon E_{\mathrm{iso}}/T$, where E_{iso} is the isotropic-equivalent energy of the relevant component, T is the duration of the relevant emission and ϵ is the overall efficiency. This efficiency is a product of several factors: $\epsilon = \epsilon_{\mathrm{dis}}\epsilon_e\epsilon_{\mathrm{rad}}\epsilon_{\mathrm{obs}}$, where a fraction ϵ_{dis} of the total energy is dissipated into internal energy, a fraction ϵ_e of the internal energy

goes into electrons (or positrons) and can in principle be radiated away, a fraction $\epsilon_{\rm rad}$ of the electron energy is actually radiated, and a fraction $\epsilon_{\rm obs}$ of the radiated energy falls within the observed range of photon energies. A fraction $\epsilon_\gamma = \epsilon_{\rm dis,GRB}\epsilon_{\rm e,GRB}\epsilon_{\rm rad,GRB}$ of the total original (isotropic-equivalent) energy $E_{\rm iso,0}$ is radiated away during the prompt emission, $E_{\gamma,\rm iso} = \epsilon_\gamma E_{\rm iso,0}$, while the remaining (kinetic) energy in the γ-ray emitting component of the outflow, $E_{\rm k,iso,0} = (1 - \epsilon_\gamma)E_{\rm iso,0}$, is transferred to the afterglow shock at $t_{\rm dec}$. In the prompt GRB emission $\epsilon_{\rm GRB} = \epsilon_\gamma\epsilon_{\rm obs,GRB}$. In the afterglow $\epsilon_{\rm dis,X} \approx 1$, and therefore $\epsilon_{\rm X} \approx \epsilon_{\rm e,X}\epsilon_{\rm rad,X}\epsilon_{\rm obs,X}$, where we concentrate on the X-ray afterglow.

Using the above expression for ϵ_γ, it becomes clearer why it is difficult for it to assume very high ($\gtrsim 0.9$) values. Whereas $\epsilon_{\rm rad,GRB} \approx 1$ is possible, and even expected if the electrons cool significantly over a dynamical time (as is typically expected in the internal-shocks model), $\epsilon_\gamma \gtrsim 0.9$ requires in addition that both $\epsilon_{\rm dis,GRB} > 0.9$ and $\epsilon_{\rm e,GRB} > 0.9$. It is difficult to achieve $\epsilon_{\rm dis,GRB} > 0.9$ (i.e. dissipate more than 90 per cent of the total energy) in most models for the prompt emission, and in particular in the internal-shocks model. Furthermore, it is not trivial to produce $\epsilon_{\rm e,GRB} > 0.9$ (i.e. more than 90 per cent of the post-shock energy going to electrons) in the internal-shocks model, where the electrons are believed to be accelerated in a shock propagating into a magnetized proton–electron plasma. In particular, this would require that less than 10 per cent of the post-shock energy goes into the protons and the magnetic field, which is difficult since the protons carry most of the energy (in kinetic form) and the magnetic field likely carries a comparable energy in the upstream fluid ahead of the shock. It might conceivably be possible if less than 10 per cent of the energy is in the magnetic field and the protons can somehow transfer their energy to the electrons, which radiate it away.

2.1 Relationship to observed quantities

The X-ray afterglow isotropic-equivalent luminosity, $L_{\rm X,iso}$, can be expressed in terms of the X-ray afterglow flux, $F_{\rm X}$,

$$L_{\rm X,iso}(t) = 4\pi d_{\rm L}^2 (1 + z)^{\beta - \alpha - 1} F_{\rm X}(t) \tag{1}$$

if $F_\nu \propto \nu^{-\beta} t^{-\alpha}$ in the relevant ranges in frequency and time.[1] The efficiency of the X-ray afterglow emission is defined as

$$\epsilon_{\rm X}(t) \equiv \frac{t L_{\rm X,iso}(t)}{E_{\rm k,iso}(t)}, \tag{2}$$

where $E_{\rm k,iso}$ is the isotropic-equivalent kinetic energy in the afterglow shock.

Using the relation $E_{\rm iso,0} = E_{\gamma,\rm iso}/\epsilon_\gamma = E_{\rm k,iso,0}/(1 - \epsilon_\gamma)$, we obtain

$$\frac{\epsilon_\gamma \epsilon_{\rm obs,GRB}}{(1 - \epsilon_\gamma)} = \frac{E_{\gamma,\rm iso}^{\rm obs}}{E_{\rm k,iso,0}} = \kappa f, \quad \kappa \equiv \frac{E_{\gamma,\rm iso}^{\rm obs}}{E_{\rm k,iso}(t_*)}, \quad f \equiv \frac{E_{\rm k,iso}(t_*)}{E_{\rm k,iso,0}}, \tag{3}$$

where t_* can be chosen as a time at which it is convenient to estimate $E_{\rm k,iso}$, and we shall use $t_* = 10$ h (since it is widely used in the literature and is typically greater than $t_{\rm break,2}$, the end of the flattish segment of the X-ray light curve).

An important question is how to estimate f and κ from observations. The most straightforward way of estimating $f = E_{\rm k,iso}(10\,{\rm h})/E_{\rm k,iso,0}$, which has been used by Nousek et al. (2006), is by evaluating the value of the X-ray flux decrement at $t_{\rm dec}$ relative to the extrapolation to $t_{\rm dec}$ of the late-time ($t > t_{\rm break,2}$) flux, and translating this flux ratio into an energy ratio according to the standard afterglow theory. When estimating f in this way we assume that

$$\epsilon_{\rm X} \sim (1 + Y)^{-1}\epsilon_{\rm e}(\nu_{\rm m}/\nu_{\rm X})^{(p-2)/2}, \tag{4}$$

where $\nu_{\rm m}$ and $\nu_{\rm c}$ are the characteristic synchrotron frequency and cooling break frequency, respectively (Sari, Piran & Narayan 1998; Granot & Sari 2002) and Y is the Compton y-parameter. This result can also be obtained under the assumptions of standard afterglow theory, as follows. The overall X-ray afterglow efficiency can be written as $\epsilon_{\rm X} \approx \epsilon_{\rm e}\epsilon_{\rm rad}\epsilon_{\rm obs}$, where $\epsilon_{\rm rad} \approx \min[1, (\nu_{\rm m}/\nu_{\rm c})^{(p-2)/2}]$ and $\epsilon_{\rm obs} \approx (1 + Y)^{-1}\max[(\nu_{\rm m}/\nu_{\rm X})^{(p-2)/2}, (\nu_{\rm c}/\nu_{\rm X})^{(p-2)/2}]$; the factor of $(1 + Y)^{-1}$ is the fraction of the radiated energy in the synchrotron component, and it is present because the synchrotron self-Compton (SSC) component typically does not contribute significantly to $F_{\rm X}(10\,{\rm h})$ but may still dominate the total radiated luminosity. The factor $(1 + Y)^{-1}$ is generally assumed to be small, consistent with the usual inference that the magnetic-to-internal energy ratio in the emission region, ϵ_B, is smaller than $\epsilon_{\rm e}$ (see equation 12).[2] However, one should keep in mind that a different time dependence of $\epsilon_{\rm X}$ (which might occur under less standard assumptions) would modify the value of f accordingly. The value of $\epsilon_{\rm obs,GRB}$ can be estimated by extrapolating the observed part of the spectrum and modelling additional spectral components that might carry considerable energy.

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[1] For the more general case, see equation (1) of Nousek et al. (2006) and the discussion thereafter. Here $L_{\rm X,iso}(t) = \int_{\nu_1}^{\nu_2} {\rm d}\nu\, L_{\nu,\rm iso}(t)$, where both ν and t are measured in the cosmological frame of the GRB; whereas $F_{\rm X}(t) = \int_{\nu_1}^{\nu_2} {\rm d}\nu\, F_\nu(t)$, where both ν and t are measured in the observer frame.

[2] At early times (less than about a day) inverse-Compton radiation is important in cooling the electrons (of Lorentz factor $\gamma_{\rm e,X}$) that emit synchrotron photons of frequency $\nu_{\rm syn}(\gamma_{\rm e,X}) = \nu_X = 10^{18}\nu_{18}$ Hz in the X-ray band. Later the X-ray emitting electrons are within the Klein–Nishina cut-off and are not cooled by inverse-Compton radiation (see Fan & Piran 2006). This can be understood as follows. Using standard estimates for the relevant quantities, we obtain $\gamma_{\rm e,X} = 4.4 \times 10^5 \epsilon_{B,-2}^{-1/4} n_0^{-1/8} E_{\rm k,iso,52}^{-1/8} \nu_{18}^{1/2} t_{10\,{\rm h}}^{3/8}$, where n_0 is the external particle density in units of cm^{-3}. The Klein–Nishina cut-off sets in when $\gamma_{\rm e}h\nu' \approx \gamma_{\rm e}h\nu/\Gamma > m_{\rm e}c^2$, so the SSC cooling of electrons with $\gamma_{\rm e,X}$ will be efficient for as long as the inequality $\gamma_{\rm e,X}h\nu_p < \Gamma m_{\rm e}c^2$ is satisfied; here $\nu_p = \max(\nu_{\rm m}, \nu_{\rm c})$ is the frequency where νF_ν peaks (i.e. where most of the synchrotron energy is radiated). We find $\gamma_{\rm e,X}h\nu_{\rm m}/\Gamma m_{\rm e}c^2 = 4.7 \times 10^{-4}\,[3(p - 2)/(p - 1)]^2\epsilon_{\rm e,-1}^2\epsilon_{B,-2}^{1/4} E_{\rm k,iso,52}^{1/4}\nu_{18}^{1/2}t_{10\,{\rm h}}^{-3/4}$ and $\gamma_{\rm e,X}\,h\nu_{\rm c}/\Gamma m_{\rm e}\,c^2 = 0.6(1 + Y)^{-2}n_0^{-1}\epsilon_{B,-2}^{-7/4}E_{\rm k,iso,54}^{-3/4}\nu_{18}^{1/2}t_{10\,{\rm h}}^{1/4}$, with the numerical coefficient in the latter expression being 0.04 when Y is evaluated for $\epsilon_{\rm e}/\epsilon_B = 10$ using equation (12).

To estimate $\kappa = E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}/E_{\mathrm{k,iso}}(10\,\mathrm{h})$ we calculate $E_{\gamma,\mathrm{iso}}^{\mathrm{obs}} = f_\gamma\, 4\pi d_{\mathrm{L}}^2 (1 + z)^{-1}$ directly from the observed γ-ray fluence, f_γ, and the measured redshift, z. The denominator, $E_{\mathrm{k,iso}}(10\,\mathrm{h})$, has been estimated following Freedman & Waxman (2001) and Kumar (2000) (see also Lloyd-Ronning & Zhang 2004) from $F_{\mathrm{X}}(10\,\mathrm{h})$ using standard afterglow theory. We reconsider this calculation in the next section.

2.2 Estimates of the kinetic energy from the X-ray afterglow observations

Using the results of Granot & Sari (2002) we find that, if $\nu_{\mathrm{X}} > \max(\nu_{\mathrm{m}}, \nu_{\mathrm{c}})$, we can estimate the kinetic energy from the observed X-ray flux. We use

$$E_{\mathrm{k,iso}}(t) = 9.2 \times 10^{52}\, \frac{g(p)}{g(2.2)} (1 + Y)^{4/(p+2)} \epsilon_{e,-1}^{4(1-p)/(p+2)} \epsilon_{B,-2}^{(2-p)/(p+2)} L_{\mathrm{X},46}^{4/(p+2)} t_{10\,\mathrm{h}}^{(3p-2)/(p+2)}\ \mathrm{erg}, \tag{5}$$

where $L_{\mathrm{X,iso}} = L_{\mathrm{X},46}\, 10^{46}\ \mathrm{erg\,s^{-1}}$ is the isotropic-equivalent X-ray luminosity in the range 2–10 keV at a time $10\,t_{10\,\mathrm{h}}$ h, both measured in the cosmological frame of the GRB [corresponding to an observed time of $t = 10(1 + z)t_{10\,\mathrm{h}}$ h and spectral range $2/(1 + z) - 10/(1 + z)$ keV], $\epsilon_{e,-1} = \epsilon_e/0.1$, $\epsilon_{B,-2} = \epsilon_B/0.01$, and

$$g(p) = \left[\frac{(p - 1)^{p-1} \exp(5.89p - 12.7)}{(5^{(p-2)/2} - 1)(p - 2)^{p-2}(p - 0.98)} \right]^{4/(p+2)}. \tag{6}$$

In equation (5), $L_{\mathrm{X,iso}}$ is evaluated through F_{X} using equation (1), and it is assumed that F_{X} is dominated by the synchrotron component. If there is a significant contribution to $F_{\mathrm{X}}(t_*)$ from the SSC component then equation (5) would overestimate $E_{\mathrm{k,iso}}(t_*)$, but the correct estimate could still be obtained if only the synchrotron contribution to $F_{\mathrm{X}}(t_*)$ is used (although in practice it might be difficult and somewhat model-dependent to disentangle the synchrotron and SSC components).

The numerical coefficient in equation (5) is higher than the corresponding coefficient in equation (7) of Lloyd-Ronning & Zhang (2004) by a factor of $\sim$10 for $p = 2.2$ (or $\sim$20 for $p = 2.5$). We track down the difference to the fact that Lloyd-Ronning & Zhang (2004) overestimated the value of ν_{m} in their equation (2) (by a factor of 38 for $p = 2.2$ compared to Granot & Sari 2002), given that their expressions for ν_{c} and $F_{\nu,\max}$ (their equations 3 and 4) are very similar to those in Granot & Sari (2002) and that $F_{\nu > \max(\nu_{\mathrm{m}}, \nu_{\mathrm{c}})} = F_{\nu,\max} \nu_{\mathrm{c}}^{1/2} \nu_{\mathrm{m}}^{(p-1)/2} \nu^{-p/2} \propto \nu_{\mathrm{m}}^{(p-1)/2}$. The numerical coefficient in equation (5) is lower than that in equation (4) of Freedman & Waxman (2001) by a factor of $\sim$7 for $p = 2.2$ (or $\sim$18 for $p = 2.5$), although it agrees with the one obtained by Fan & Piran (2006) to within a factor of 2.[3]

The factor involving $(1 + Y)$ in equation (5) represents inverse-Compton losses by the electrons that emit the observed synchrotron X-rays. The inverse-Compton losses can tap most of the energy of these fast cooling electrons, leaving only a modest fraction of their energy in the synchrotron component. The resulting very-high-energy SSC emission reduces the synchrotron flux in the relevant (X-ray) power-law segment of the spectrum by a factor of $(1 + Y)$ (Sari & Esin 2001; Granot & Sari 2002).[4] Using our fiducial parameters we find $(1 + Y) \approx 3.7$. The factor $(1 + Y)^{4/(p+2)}$ in equation (5) then brings the numerical coefficient in that equation up to 3.2×10^{53} erg for $p = 2.2$ and 5.8×10^{53} erg for $p = 2.5$. Neither Freedman & Waxman (2001) nor Lloyd-Ronning & Zhang (2004) took the effects of inverse-Compton cooling into account. Our inclusion of these effects reduces the discrepancy between the numerical coefficient in equation (5) and that in equation (4) of Freedman & Waxman (2001) to a factor of $\sim$2 ($\sim$5.5) for $p = 2.2$ ($p = 2.5$). It, however, increases the discrepancy between our numerical coefficient and that in equation (7) of Lloyd-Ronning & Zhang (2004) to a factor of $\sim$35–70. This, in turn, leads to large differences between the latter authors' estimates of the γ-ray efficiency and ours.

Using equations (2) and (5), we can express the X-ray efficiency as

$$\epsilon_{\mathrm{X}}(t) = 3.5 \times 10^{-3} \left[\frac{g(p)}{g(2.2)} \right]^{-(p+2)/4} (1 + Y)^{-1} \epsilon_{e,-1}^{p-1} \epsilon_{B,-2}^{(p-2)/4} E_{\mathrm{k,iso},52}^{(p-2)/4} t_{10\,\mathrm{h}}^{-3(p-2)/4}$$

$$= 3.9 \times 10^{-3}\, \frac{g(2.2)}{g(p)} (1 + Y)^{-4/(p+2)} \epsilon_{e,-1}^{4(p-1)/(p+2)} \epsilon_{B,-2}^{(p-2)/(p+2)} L_{\mathrm{X},46}^{(p-2)/(p+2)} t_{10\,\mathrm{h}}^{-2(p-2)/(p+2)}, \tag{7}$$

where $E_{\mathrm{k,iso},52} = E_{\mathrm{k,iso}}(t)/(10^{52}\,\mathrm{erg})$. For our fiducial values (and for $L_{\mathrm{X},46}$ rather than $E_{\mathrm{k,iso},52}$), the numerical coefficient in equation (7) is 1.1×10^{-3} for $p = 2.2$ and 6.3×10^{-4} for $p = 2.5$.

When comparing the γ-ray efficiencies inferred by using our equation (5) and the results of Lloyd-Ronning & Zhang (2004), one should note another important difference between our formulation and theirs. Lloyd-Ronning & Zhang (2004) assumed that the initial afterglow kinetic energy is reduced by a factor $\mathcal{R} \sim 8$ due to radiative losses at early times ($t < 10$ h). This is equivalent to choosing the parameter f in equation (3) to be $\ll 1$, whereas the recent *Swift* observations indicate that it is most likely $\gtrsim 1$, and quite possibly $\gg 1$ (see Section 2.4). As

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[3] Most previous works (e.g. Freedman & Waxman 2001; Lloyd-Ronning & Zhang 2004) first derived the values of the peak flux and spectral break frequencies, each of which typically has an uncertainty of at least a factor of a few, and then used those estimates to calculate the flux normalization in the various power-law segments of the spectrum. In the relevant segment, $\nu > \max(\nu_{\mathrm{m}}, \nu_{\mathrm{c}})$, the flux depends on the values of $F_{\nu,\max}$, ν_{m} and ν_{c}, each of which is not accurately known, so the normalization calculated using the above method is highly uncertain. In contrast, in this work we use the results of Granot & Sari (2002), who calculated the flux normalization directly (i.e. without first estimating the values of the peak flux and spectral break frequencies) using a detailed scheme that employs a minimal number of simplifying assumptions. In the relevant power-law segment of the spectrum the flux normalization calculated in this fashion is rather robust: it relies mainly on the assumption of a power-law electron energy distribution just behind the shock and is rather insensitive to the exact shape of the low-energy cut-off.

[4] This is assuming that the SSC component does not contribute considerably to the observed X-ray flux, which is typically the case at $t_* = 10$ h.

is evident from the expression for the γ-ray efficiency ϵ_γ in equation (3), the low value chosen for f partly compensates for the higher value that Lloyd-Ronning & Zhang (2004) obtain for the parameter $\kappa \propto 1/E_{k,iso}(t_*)$, leading to a lower estimate of ϵ_γ. Without this assumption, the discrepancy between our estimates of ϵ_γ and those of Lloyd-Ronning & Zhang (2004) would directly reflect the large differences in our respective estimates of $E_{k,iso}(t_*)$.

Equations (5)–(7) are valid for $p > 2$, but they can be easily generalized to $p \lesssim 2$ by substituting $\epsilon_e \to \bar{\epsilon}_e(p-1)/(p-2)$, where $\bar{\epsilon}_e = \epsilon_e \gamma_{min}/\langle\gamma_e\rangle$, $\langle\gamma_e\rangle = \int d\gamma_e(dN/d\gamma_e)\gamma_e \left[\int d\gamma_e(dN/d\gamma_e)\right]^{-1}$ is the average electron Lorentz factor, and the electron energy distribution is a power law of index p at low energies above γ_{min}. If the electron energy distribution is a single power law, $dN/d\gamma_e \propto \gamma_e^{-p}$ for $\gamma_{min} < \gamma_e < \gamma_{max}$, then

$$\frac{\epsilon_e}{\bar{\epsilon}_e} = \frac{\langle\gamma_e\rangle}{\gamma_{min}} = \left(\frac{p-1}{p-2}\right)\frac{1-(\gamma_{max}/\gamma_{min})^{2-p}}{1-(\gamma_{max}/\gamma_{min})^{1-p}} = \begin{cases} \approx (p-1)/(p-2) & p > 2, \\ \ln(\gamma_{max}/\gamma_{min}) & p = 2, \\ \approx (\gamma_{max}/\gamma_{min})^{2-p}(p-1)/(2-p) & 1 < p < 2, \\ (\gamma_{max}/\gamma_{min})/\ln(\gamma_{max}/\gamma_{min}) & p = 2, \\ \approx (\gamma_{max}/\gamma_{min})(1-p)/(2-p) & p < 1 \end{cases} \tag{8}$$

This results in slightly modified equations:

$$E_{k,iso}(t) = 1.19 \times 10^{52} \frac{\bar{g}(p)}{\bar{g}(2.2)}(1+Y)^{4/(p+2)}\bar{\epsilon}_{e,-1}^{4(1-p)/(p+2)}\epsilon_{B,-2}^{(2-p)/(p+2)}L_{X,46}^{4/(p+2)}t_{10h}^{(3p-2)/(p+2)} \text{ erg}, \tag{9}$$

$$\bar{g}(p) = \left[\frac{(p-2)\exp(5.89p - 12.7)}{(5^{(p-2)/2}-1)(p-0.98)}\right]^{4/(p+2)}, \tag{10}$$

$$\epsilon_X(t) = 3.01 \times 10^{-2}\left[\frac{\bar{g}(p)}{\bar{g}(2.2)}\right]^{-(p+2)/4}(1+Y)^{-1}\bar{\epsilon}_{e,-1}^{p-1}\bar{\epsilon}_{B,-2}^{(p-2)/4}E_{k,iso,52}^{(p-2)/4}t_{10h}^{-3(p-2)/4}$$
$$= 3.03 \times 10^{-3}\frac{\bar{g}(2.2)}{\bar{g}(p)}(1+Y)^{-4/(p+2)}\bar{\epsilon}_{e,-1}^{4(p-1)/(p+2)}\epsilon_{B,-2}^{(p-2)/(p+2)}L_{X,46}^{(p-2)/(p+2)}t_{10h}^{-2(p-2)/(p+2)}. \tag{11}$$

Note that the numerical coefficient was calculated in Granot & Sari (2002) only for $p = 2.2, 2.5, 3$ and interpolated between these values. Extrapolating that formula to $p \lesssim 2$ could potentially be very inaccurate.

For simplicity we use the expression for Y that is valid in the fast-cooling regime,

$$Y = \frac{(1+4\epsilon_e/\epsilon_B)^{1/2}-1}{2} \approx \begin{cases} (\epsilon_e/\epsilon_B)^{1/2} & \epsilon_e/\epsilon_B \gg 1 \\ \epsilon_e/\epsilon_B & \epsilon_e/\epsilon_B \ll 1 \end{cases}, \tag{12}$$

which is still reasonable at 10 h, even if it is slightly after the transition to slow cooling. More generally, ϵ_e/ϵ_B should be multiplied by $\epsilon_{rad} \approx \min[1, (\nu_m/\nu_c)^{(p-2)/2}]$, where for $p < 2$ and $\nu_c < \nu_{max}$ we have $\epsilon_{rad} \approx 1$.

Before applying these relations to observed bursts, we remark on a common misconception concerning the dependence of $E_{k,iso}$ that is inferred from F_X on ϵ_B and ϵ_e. It has been argued that $E_{k,iso}$ is very insensitive to the exact value of ϵ_B (e.g. Freedman & Waxman 2001; Piran et al. 2001). This follows from the observation (see equation 5) that for $\nu_X > \max(\nu_m, \nu_c)$, $E_{k,iso} \propto (1+Y)^{4/(p+2)}\epsilon_e^{-4(p-1)/(p+2)}\epsilon_B^{-(p-2)/(p+2)}$, which suggests that $E_{k,iso}$ depends very weakly on ϵ_B. However, this holds only in the limit where $(1+Y) \approx 1$, which corresponds to $Y \ll 1$ and $\epsilon_B \gg \epsilon_e$, whereas observations suggest that the opposite limit typically applies, $\epsilon_B \ll \epsilon_e$, in which case $(1+Y) \approx Y \approx (\epsilon_e/\epsilon_B)^{1/2}$ and $E_{k,iso} \propto \epsilon_B^{-p/(p+2)}$. This is a significantly stronger dependence on ϵ_B. (Note that the inferred value of ϵ_B varies by ~ 2 orders of magnitude among different afterglows, from $\sim 10^{-3}$ to ~ 0.1, corresponding to a change of more than an order of magnitude in the estimated value of $E_{k,iso}$.) The dependence of $E_{k,iso}$ on ϵ_e is stronger, $\propto \epsilon_e^{-4(p-1)/(p+2)}$ (i.e. slightly steeper than an inverse linear relation) in the limit $\epsilon_B \gg \epsilon_e$, but only $\propto \epsilon_e^{-2(2p-3)/(p+2)}$ in the more relevant limit of $\epsilon_B \ll \epsilon_e$. In the latter case, the dependence of $E_{k,iso}$ on ϵ_B is rather similar to its dependence on ϵ_e.[5] However, ϵ_e appears to vary much less than ϵ_B among different afterglows, only covering a range of about one order of magnitude (between $\sim 10^{-1.5}$ and $\sim 10^{-0.5}$), which corresponds to a variation in $E_{k,iso}$ by a factor of ~ 4. It is also worth noting that the expression for $E_{k,iso}$ has some (non-trivial) dependence on the value of p (see equation 6).

2.3 Pre-*Swift* GRBs

Table 1 shows the estimated values of κ deduced from the observational properties of 17 pre-*Swift* GRBs with known redshifts, using the samples of Berger et al. (2003) and Bloom et al. (2003). We provide the values $E_{k,iso,10h} = E_{k,iso}(10h)$ and κ for our fiducial parameter values ($\epsilon_e = 0.1$, $\epsilon_B = 0.01$ and $p = 2.2$). The value of $E_{k,iso,10h}$ and therefore of κ depends on the values of the microphysical parameters (ϵ_e, ϵ_B and p) that are not well known. Therefore, we also calculate $E_{k,iso,10h}$ and κ (for those GRBs for which this is possible) using the values of the microphysical parameters inferred from the fits to the broad-band afterglow data that were performed by Panaitescu & Kumar (2002, denoted

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[5] Quite often $Y \sim 1$–2 is inferred, in which case neither of the asymptotic limits $Y \ll 1$ and $Y \gg 1$ is applicable and the dependence of $E_{k,iso}$ on ϵ_B and ϵ_e does not have a power-law form but rather the more complex form given by equations (5) and (12).

Table 1. Estimates of κ for pre-*Swift* GRBs with known redshifts.

GRB #	z	$L_{X,iso,10h}^a$ (10^{46} erg s^{-1})	$E_{\gamma,iso}^{obs,b}$ (10^{52} erg)	$E_{k,iso,10h}^c$ (10^{52} erg)	$E_{k,iso,10h}^{(PK02),d}$ (10^{52} erg)	$E_{k,iso,10h}^{(PK02),e}$ (10^{52} erg)	$E_{k,iso,10h}^{(YO3),f}$ (10^{52} erg)	$E_{k,iso,1d}^{(YO3),g}$ (10^{52} erg)	κ^c	κ^d (PK02)	κ^e (PK02)	κ_{10h}^f (YO3)	κ_{1d}^h (YO3)
970228	0.695	0.682	1.416	22.2	–	–	–	–	0.064	–	–	–	–
970508	0.835	0.374	0.546	12.5	1.31	4.0	1.50	1.6	0.044	0.42^i	0.14^i	0.36	0.34
970828	0.985	1.76	21.98	54.8	–	–	–	–	0.40	–	–	–	–
971214	3.418	8.96	21.05	258	–	–	–	–	0.082	–	–	–	–
980613	1.096	0.536	0.536	17.7	–	–	–	–	0.030	–	–	–	–
980703	0.966	1.02	6.012	32.6	–	–	61.6	13	0.18	–	–	0.098	0.46
990123	1.600	12.83	143.8	364	339	22	–	–	0.40	0.42	6.4	–	–
990510	1.619	8.209	17.64	238	70^i	9.6	–	–	0.074	0.25^i	1.8	–	–
990705	0.840	0.123	25.60	4.35	–	–	–	–	5.89	–	–	–	–
991216	1.02	18.32	53.54	510	–^{j}	9.9	–	–	0.10	–^{j}	5.4	–	–
000210	0.846	0.183	16.93	6.35	–	–	–	–	2.67	–	–	–	–
000926	2.037	7.169	27.97	209	36.7	3.2	272	15	0.13	0.76	8.7	0.10	1.86
010222	1.477	13.79	85.78	389	19^i	16	–	–	0.22	4.5^i	5.4	–	–
011211	2.14	0.886	6.723	28.5	–	–	–	–	0.24	–	–	–	–
020405	0.698	1.729	7.201	53.9	–	–	–	–	0.13	–	–	–	–
020813	1.254	12.12	77.50	344	–	–	–	–	0.23	–	–	–	–
021004	2.323	6.536	5.560	191	–	–	–	–	0.029	–	–	–	–

Notes. Estimates for $E_{k,iso}(10h) = E_{k,iso,10h}$ and $\kappa = E_{\gamma,iso}^{obs}/E_{k,iso,10h}$ for the GRBs with known redshift.

$^a L_{X,iso,10h} = L_{X,iso}(10h)$ from table 2 of Berger et al. (2003); $^b E_{\gamma,iso}^{obs}$ in the 20–2000 keV range form table 2 of Bloom et al. (2003); ^{c}calculated using equation (5) with our fiducial values of the microphysical parameters: $\epsilon_e = 0.1$, $\epsilon_B = 0.01$ and $p = 2.2$; ^{d}substituting the best-fitting values of the microphysical parameters from table 2 of Panaitescu & Kumar (2002) into equation (5); ^{e}the value for $E_{k,iso}(10h)$ is taken to be $\approx 0.5 E_{k,iso,0}$, where $E_{k,iso,0}$ is taken from Panaitescu & Kumar (2002) and the factor of ≈ 0.5 accounts for the average factor by which the energy is reduced relative to $E_{k,iso,0}$ due to radiative losses (Panaitescu, private communication); ^{f}substituting the best-fitting values of the microphysical parameters from table 1 of Yost et al. (2003) into equation (5); $^g E_{k,iso,1d} = E_{k,iso}(1 \, d) = E_{\gamma,iso}^{obs}/\kappa_{1d}$ from table 1 of Yost et al. (2003); $^h \kappa_{1d} = E_{\gamma,iso}^{obs}/E_{k,iso,1d}$; ^{i}these values are for a fit to a uniform external density; PK02 got a significantly better fit to a wind density for which we derive $\kappa = 0.47$ and 0.17 instead of 0.42 and 0.14, respectively; ^{j}in these cases, Panaitescu & Kumar (2002) found $p < 2$ which introduces an uncertainty through the extrapolation that is involved in the expression for the numerical co-efficient in equation (9); for GRB 991216 PK02 found that the X-ray emission is from electrons with $\gamma_e > \gamma_{max}$ where $dN/d\gamma_e \propto \gamma_e^{-q}$ with $p < 2 < q_0$ with $2 < q < q_0$ so that we cannot readily substitute their results into our equations.

Table 2. Estimates of f and κ for *Swift* GRBs with known redshifts.

GRB #	z	$L_{\mathrm{X,iso,10h}}{}^a$ $(10^{46}\,\mathrm{erg\,s^{-1}})$	$E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}{}^{\,b}$ $(10^{52}\,\mathrm{erg})$	$E_{\mathrm{k,iso,10h}}{}^c$ $(10^{52}\,\mathrm{erg})$	p^d	κ^e	$f_{\min}{}^f$	$f_{\max}{}^f$
050126	1.29	0.12	2.2	6.86	3	0.055	–	–
050315	1.949	16	18	699	2.1	0.034	11	29
050318	1.44	0.60	3.9	$\gtrsim 28.4$	2.1	$\lesssim 0.18$	4.2	170
050319	3.24	5.1	12.1	$\gtrsim 118$	2.6	$\lesssim 0.039$	12	76
050401	2.90	9.8	137	433	2.1	0.41	5.6	14
050408	1.236	1.4	2.9	38.5	2.3	0.058	–	–
050416A	0.6535	0.091	0.09	4.51	2.1	0.026	2.2	9.9
050505	4.3	2.3	89	$\gtrsim 59.0$	2.6	$\lesssim 0.58$	19	1800
050525A	0.606	0.12	3.1	$\gtrsim 3.96$	2.4	$\lesssim 0.47$	2.1	5.9
050603	2.821	1.1	126	$\gtrsim 29.6$	2.4	$\lesssim 2.57$	–	–

Notes. The estimates for $f = E_{\mathrm{k,iso,10h}}/E_{\mathrm{k,iso,0}}^{\mathrm{obs}}$ and $\kappa = E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}/E_{\mathrm{k,iso,10h}}$ for the GRBs with known redshift from the Nousek et al. (2006) sample.
$^a L_{\mathrm{X,iso,10h}} = L_{\mathrm{X,iso}}(10\,\mathrm{h})$ in the 2–10 keV range at (both the time and the photon energies measured in the cosmological frame of the GRB) from table 2 of Nousek et al. (2006); $^b E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}$ in the 20–2000 keV range (in the GRB's cosmological frame) form table 2 of Nousek et al. (2006); ^{c} calculated using equation (5) with $\epsilon_\mathrm{e} = 0.1$, $\epsilon_B = 0.01$, and the values of p from this table; ^{d} estimated using the measured spectral slope in the X-rays (Nousek et al. 2006) and attempting to fit it into the range $2 < p < 3$ if allowed within the errors on the spectral slope; $^e \kappa = E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}/E_{\mathrm{k,iso,10h}}$ estimated using the values from this table; $^f f_{\min}$ and $f_{\max}$ are taken from table 3 of Nousek et al. (2006), and are estimated using the measured X-ray flux at $t_{\mathrm{break,1}}$, and the extrapolated X-ray flux at T_{GRB}, respectively (see Nousek et al. 2006 for details).

by PK02, using table 2 therein) and by Yost et al. (2003, denoted by Y03, using table 1 therein). Finally, we compare the values we obtain for $E_{\mathrm{k,iso,10h}}$ using equation (5) to those obtained for $E_{\mathrm{k,iso}}(1\,\mathrm{d}) = E_{\mathrm{k,iso,1d}}$ by Y03 and those obtained for $E_{\mathrm{k,iso}}(10\,\mathrm{h}) \approx 0.5 E_{\mathrm{k,iso,0}}$ by PK02. Also shown in Table 1 are the corresponding values of κ, including $\kappa_{\mathrm{1d}} = E_{\gamma,\mathrm{iso}}/E_{\mathrm{k,iso,1d}}^{Y03}$.

When using the fiducial parameters $\epsilon_\mathrm{e} = 0.1$, $\epsilon_B = 0.01$ and $p = 2.2$ our estimates for $E_{\mathrm{k,iso,10h}}$ are significantly larger than the estimates of Lloyd-Ronning & Zhang (2004), who use the same values. This can be traced to the difference in the numerical factor that appears in equation (5). These relatively large values of $E_{\mathrm{k,iso,10h}}$ lead to a typical value of κ around 0.1–0.2, for which the γ-ray efficiency problem would not be very severe. Similar results were obtained by Fan & Piran (2006), whose estimates for $E_{\mathrm{k,iso,10h}}$ are within a factor of 2 of those presented here.

The situation is different when we use the values of the microphysical parameters from the pre-*Swift* afterglow fits. In these cases the values of $E_{\mathrm{k,iso,10h}}$ are typically lower, resulting in higher estimates for κ. The estimates of the kinetic energy from the PK02 fits to the afterglow data (rather than from using equation 5) are generally the lowest.[6] The corresponding values of κ are close to unity. These results reflect the pre-*Swift* inference that there exists a rough equality between the isotropic-equivalent γ-ray and late (10 h) kinetic energies. The comparison between the detailed calculations and those based on equation (5) may suggest that we might be better off adopting different fiducial parameters (e.g. $\epsilon_\mathrm{e} = 0.3$, $\epsilon_B = 0.08$ and $p = 2.2$). We also note that the values of $E_{\mathrm{k,iso,10h}}$ obtained by substituting the values of the microphysical parameters from the afterglow fits into equation (5) are generally higher (corresponding to lower values of κ) compared to those derived directly by those fits. The 'typical' values of the microphysical parameters inferred from the fits are roughly $\epsilon_\mathrm{e} \approx 0.3$, $\epsilon_B \approx 0.03$ and $p \approx 2.2$.

The values of κ obtained here are crucial to the overall estimate of ϵ_γ and to the origin of the 'high-efficiency problem.' Considering Table 1, one should proceed with care in view of the large dispersion in the different estimates of $E_{\mathrm{k,iso,10h}}$. An example is a factor of 4.7 between the independent estimates of PK02 and of Y03 for GRB 000926 (and a factor of 2.5 in the opposite direction for GRB 970508). The dispersion is even greater between the estimates of $E_{\mathrm{k,iso}}$ from the afterglow fits and the values obtained using equation (5) for the same values of the microphysical parameters – a factor of 18 for Y03 (and 11 for PK02) for GRB 000926.

2.4 *Swift* GRBs

The new result found by *Swift* is the appearance of a rapid decay followed by a shallow decline phase in the X-ray afterglow. For seven out of the 10 *Swift* GRBs considered here there was a clear observation of the two breaks in the light curve, at $t_{\mathrm{break,1}}$ and $t_{\mathrm{break,2}}$ (the beginning and end, respectively, of the flattish segment of the X-ray light curve). If we interpret the shallow decline as arising from an additional injection of energy into the afterglow shock (see Section 1) then the ratio of the X-ray fluxes at the end and at the beginning of this phase can be used to estimate f (the corresponding ratio of the kinetic energies) for these bursts (Nousek et al. 2006). A lower limit, $f_{\min}$, is obtained using the flux decrement at $t_{\mathrm{break,1}}$ relative to the extrapolation back to that time of the late-time ($t > t_{\mathrm{break,2}}$) flux. An approximate upper limit, $f_{\max}$, is

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[6] In PK02 the values of $E_{\mathrm{k,iso}}(10\,\mathrm{h})$ are typically a factor of ~ 2 smaller than $E_{\mathrm{k,iso,0}}$ due to radiative losses at early times (Panaitescu, private communication; no energy injection is assumed in that work).

obtained by assuming that the flat part of the emission from the forward shock starts at $t_{\rm dec} \sim T_{\rm GRB}$ and is simply buried underneath the tail of the prompt emission at $t < t_{\rm break,1}$. While formally $f_{\rm min} < f \lesssim f_{\rm max}$, it is reasonable, in the context of the basic jet model, that the assumption made to calculate $f_{\rm max}$ is basically applicable, so that $f \sim f_{\rm max}$. Under this assumption one infers $f \gtrsim 10$, and in some cases even a much larger value of $f (\sim 10^2 - 10^3)$.

The results for f need to be combined with an estimate of κ. Table 2 shows the values of $E_{\rm k,iso,10h}$ and κ for the 10 *Swift* GRBs with known redshifts from the Nousek et al. (2006) sample, estimated using equation (5) with $\epsilon_{\rm e} = 0.1$ and $\epsilon_B = 0.01$. The value of p was derived using the measured spectral slope in the X-ray band (attempting to fit it into the range $2 < p < 3$ if allowed within the measurement errors). Lacking any broad-band fits to *Swift* bursts, this is the best direct evidence that we have so far from these data. The resulting values of κ are similar to those from the pre-*Swift* era (see Table 1). With the exception of GRB 050401, for which $\kappa = 0.41$, and a few other bursts for which we only have upper limits, the inferred values of κ are less than 0.1. If this is the correct value of κ then, using

$$\epsilon_\gamma = \left(1 + \frac{\epsilon_{\rm obs,GRB}}{\kappa f}\right)^{-1} \tag{13}$$

(see equation 3), we find that with $f \sim 10$ the overall γ-ray efficiency is not larger than ~ 0.5 (assuming $\epsilon_{\rm obs,GRB} \sim 1$). A similar conclusion was reached by Fan & Piran (2006).

One may question this conclusion in view of the fact that, in pre-*Swift* bursts, broad-band analyses of the afterglow data generally resulted in a significantly lower values of $E_{\rm k,iso,10h}$, and correspondingly higher values of κ, compared to those obtained from equation (5) with the same fiducial values of the microphysical parameters ($\epsilon_{\rm e} = 0.1$ and $\epsilon_B = 0.01$; see Table 1). Furthermore, the choice of the fiducial values of the microphysical parameters is somewhat arbitrary, and it affects the resulting values of $E_{\rm k,iso,10h}$ and κ. It is reasonable to expect that the values of the microphysical parameters that would have been inferred from a broad-band fit to the afterglow data of the *Swift* bursts would have led to higher estimates of $E_{\rm k,iso,10h}$ and κ that were similar to those derived for the pre-*Swift* GRBs. The latter values, however, vary among the different estimates, from as high as ~ 5–8 to as low as ~ 0.1–0.3 (see Table 1). In light of this, one may adopt a 'typical' value of $\kappa \sim 1$, keeping in mind that there is an uncertainty of almost an order of magnitude in this value.

The adoption of this higher value of κ (~ 1) for the *Swift* GRBs, similar to the values inferred from broad-band modelling of pre-*Swift* sources, together with the interpretation of the shallow decay phase as arising from energy injection (and hence $f \sim 10$) leads to the conclusion that typically $\epsilon_\gamma \sim 0.9$, and in some cases ϵ_γ is even as high as ~ 99 per cent (or, equivalently, $1 - \epsilon_\gamma \approx \epsilon_{\rm obs,GRB}/\kappa f$ is as low as $\sim 10^{-3} - 10^{-2}$). Such a high γ-ray efficiency would be extremely hard to produce in the internal-shocks model (see Section 1).

If, on the other hand, $\kappa \sim 0.1$ and there is energy injection (i.e. $f \sim 10$), or if $\kappa \sim 1$ and there is no energy injection (i.e. $f = 1$; see Section 3), then this would imply a significantly lower typical γ-ray efficiency, $\epsilon_\gamma \sim 0.5$, although in some cases ϵ_γ would still be as high as ~ 90 per cent (or, equivalently, $1 - \epsilon_\gamma \approx \epsilon_{\rm obs,GRB}/\kappa f$ would still be as low as $\sim 10^{-2} - 10^{-1}$). Even the latter, more moderate, requirements on the γ-ray efficiency are not easily satisfied in the internal-shocks model, although they might possibly still be accommodated in this scenario (Kobayashi et al. 1997; Kumar 1999; Guetta et al. 2001). Finally, if $\kappa \sim 0.1$ and the shallow decline does not arise from energy injection (i.e. $f = 1$) but, say, from varying afterglow efficiency (as discussed in Section 3), then the γ-ray efficiency would typically be $\epsilon_\gamma \sim 0.1$. In the latter case, the results are fully consistent with the predictions of the internal-shocks model.

2.5 Some caveats

The discussion so far was relevant to the power-law segment of the spectrum where $\nu_{\rm X} > \max(\nu_{\rm m}, \nu_{\rm c})$. If, instead, $\nu_{\rm m} < \nu_{\rm X} < \nu_{\rm c}$, we have only a lower limit on the value of $E_{\rm k,iso}$ from this consideration (i.e. from equation 5) that corresponds to an upper limit on the value of κ, the true value being smaller than this upper bound by a factor of $[\nu_{\rm c}(10\,{\rm h})/\nu_{\rm X}]^{2/(p+2)}$ (which, however, is not typically expected to be very large).

It is possible that there is a contribution to $F_{\rm X}(t_{\rm dec})$ from a component of the outflow that is not along the line of sight, or for some other reason did not contribute to the observed γ-ray emission. In this case equations (3) and (5) would provide a lower limit on ϵ_γ, rather than directly determining its value. If both the prompt γ-ray emission and $F_{\rm X}(t_{\rm dec})$ were dominated by emission from angles $\theta > 1/\Gamma$ relative to the line of sight (where Γ is the Lorentz factor of the outflow), then again equations (3) and (5) would provide a lower limit on ϵ_γ, since the beaming of radiation away from the line of sight is expected to be either comparable or somewhat smaller during the afterglow emission at $t_{\rm dec}$ compared to the prompt GRB.

It is also possible that some fraction ϵ_* of the original energy $E_{\rm iso,0}$ ended up in a totally different form, i.e. was not radiated during the prompt emission and did not end up in the kinetic energy of the afterglow shock. This could occur, for example, if along some directions a forward shock is not formed (or at least not formed efficiently) for a very pure Poynting-flux outflow. In such a case, some of the energy (potentially even most of the energy) could be carried out to very large distances (in principle out to infinity) in electromagnetic form (such as low-frequency electromagnetic waves).[7] Alternatively, a good fraction of the energy might be carried away in high-energy cosmic rays and neutrinos (Waxman 1995; Waxman & Bahcall 1997) and thus would not contribute to the kinetic energy of the afterglow shock. In this case we have $E_{\rm k,iso,0} = (1 - \epsilon_\gamma - \epsilon_*) E_{\rm iso,0}$ and we need to make the substitution $(1 - \epsilon_\gamma) \rightarrow (1 - \epsilon_\gamma - \epsilon_*)$ everywhere, so the estimate (13) for

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[7] Both the formation of a forward shock and the ability of energy to escape to infinity in electromagnetic form have not yet been fully worked out (e.g. Melatos & Melrose 1996; Smolsky & Usov 2000).

ϵ_γ (assuming $\epsilon_* = 0$) should be multiplied by $(1 - \epsilon_*)$,

$$\epsilon_\gamma = (1 - \epsilon_*)\left(1 + \frac{\epsilon_{\mathrm{obs,GRB}}}{\kappa f}\right)^{-1}. \tag{14}$$

This means that $\epsilon_\gamma < 1 - \epsilon_*$ (even for $\kappa f \gg \epsilon_{\mathrm{obs,GRB}}$), and therefore $\epsilon_* \gtrsim 0.5$ would imply $\epsilon_\gamma \lesssim 0.5$. Thus, even when one infers $\epsilon_\gamma/(1 - \epsilon_\gamma) \gg 1$ and hence $1 - \epsilon_\gamma \ll 1$ under the usual assumption that $\epsilon_* = 0$ (or at least $\epsilon_* \ll 1 - \epsilon_\gamma$), the intrinsic γ-ray efficiency might still be significantly smaller, and is $\epsilon_\gamma \lesssim 0.5$ for $\epsilon_* \gtrsim 0.5$. Note that, in order for GRBs to produce the highest energy cosmic rays, their energy should be comparable to that of the prompt γ-rays (Waxman 1995, 2004), i.e. $\epsilon_\gamma \lesssim \epsilon_*$ (the inequality arising since there might be other forms of energy that escape the prompt-emission site without being directly detected), and therefore $\epsilon_\gamma \lesssim 0.5$.

One should, however, keep in mind that high-energy cosmic rays and neutrinos must tap the same dissipated energy that also powers the prompt γ-ray emission. Therefore, for the same observed energy in γ-rays and inferred kinetic energy in the afterglow, in addition to increasing the required total initial energy $E_{\mathrm{iso},0}$ by a factor of $(1 - \epsilon_*)^{-1}$, these particles would also increase the required dissipated energy $\epsilon_{\mathrm{dis,GRB}}E_{\mathrm{iso},0}$ and the fraction $\epsilon_{\mathrm{dis,GRB}}$ of the dissipated energy that ends up in γ-rays.[8] In other words, $\epsilon_{\mathrm{dis,GRB}} \geqslant \epsilon_\gamma + \epsilon_* = (E_{\gamma,\mathrm{iso}} + E_{*,\mathrm{iso}})/(E_{\mathrm{k,iso},0} + E_{\gamma,\mathrm{iso}} + E_{*,\mathrm{iso}}) > E_{\gamma,\mathrm{iso}}/(E_{\mathrm{k,iso},0} + E_{\gamma,\mathrm{iso}})$, where $E_{*,\mathrm{iso}} = \epsilon_* E_{\mathrm{iso},0}$ is the (isotropic-equivalent) energy in cosmic rays and neutrinos, and $E_{\mathrm{k,iso},0}$ and $E_{\gamma,\mathrm{iso}}$ are determined (at least in principle) by observations. Clearly, $E_{*,\mathrm{iso}}$ increases the lower limit on $\epsilon_{\mathrm{dis,GRB}}$. Nevertheless, the fact that $E_{*,\mathrm{iso}}$ may reduce ϵ_γ to $\lesssim 0.5$ even for $\kappa f \gg 1$ makes it possible to have $\epsilon_{\mathrm{e,GRB}} \lesssim 0.5$, which should be easier to accommodate for shock acceleration in a proton–electron plasma. Still, dissipating and getting rid of almost all of the energy (through radiation, cosmic rays, neutrinos, etc.) and leaving only a small fraction of the original energy in the kinetic energy of the forward shock, as is required for $\kappa f \gg 1$, is not an easy task for any model of the prompt emission.

3 EFFICIENCY OF THE X-RAY AFTERGLOW EMISSION

As noted in Section 1, one of the new features discovered by *Swift* is the early shallow decline phase: $F_X \propto t^{-\alpha}$ with $0.2 \lesssim \alpha \lesssim 0.8$ for $t_{\mathrm{break},1} \lesssim t \lesssim t_{\mathrm{break},2}$. During this phase, $tF_X(t)$ increases with time. Using the definition of $\epsilon_X(t)$ (equation 2) and the relation between $L_{X,\mathrm{iso}}$ and F_X (equation 1), we find that

$$\frac{\epsilon_X(t)E_{\mathrm{k,iso}}(t)}{tF_X(t)} = 4\pi d_L^2(1 + z)^{\beta - \alpha - 1} \tag{15}$$

is constant in time. Note that if ν and t in the expression for $F_X(t)$ were referred to the GRB rest frame rather than to the observer frame then the factor $(1 + z)^{\beta - \alpha - 1}$ on the right-hand side of equation (15) would be eliminated, and with it any potential (weak) time dependence resulting from a possible temporal variation of α or β.

Now, if the observed frequencies satisfy $\nu_X > \max(\nu_m, \nu_c)$, and $p > 2$, then equation (4) is applicable. Furthermore, if the afterglow shock evolves according to the adiabatic self-similar solution of Blandford & McKee (1976) and the fractions of the post-shock internal energy in electrons (ϵ_e) and in magnetic field (ϵ_B) are constant in time, then $\nu_m \propto t^{-3/2}$. Under these circumstances, $\epsilon_X \propto t^{-3(p-2)/4}$ decreases (slowly) with time. As can be seen from equation (15), $\epsilon_X(t)E_{\mathrm{k,iso}}(t) \propto tF_X(t)$. Therefore, the observed rise in $tF_X(t)$ implies a similar rise in $\epsilon_X(t)E_{\mathrm{k,iso}}(t)$. Given the expected decrease of $\epsilon_X(t)$ with t for $p > 2$, the slowly decaying portion of the light curve has been attributed by several researchers to an increase in $E_{\mathrm{k,iso}}(t)$, i.e. to some sort of energy injection into the forward shock (e.g. Nousek et al. 2006; Panaitescu et al. 2006; Zhang et al. 2006; Granot & Kumar 2006).

It is, however, conceivable that the rise in $tF_X(t)$ corresponds, at least in part, to an increase of $\epsilon_X(t)$ with time. One way in which this could be brought about is if p were < 2 (assuming ϵ_e and ϵ_B remain constant).[9] In this case $N(\gamma_e) \propto \gamma_e^{-p}$ for $\gamma_{\min} < \gamma_e < \gamma_{\max}$ and $\epsilon_X \sim (1 + Y)^{-1}\epsilon_e(\nu_{\max}/\nu_X)^{(p-2)/2}$ [the same as equation (4) for $p < 2$, but with $\nu_{\max}$ replacing ν_m], where $\nu_{\max} = \nu_{\mathrm{syn}}(\gamma_{\max}) \propto \gamma B\gamma_{\max}^2 \propto \gamma^4 \rho_{\mathrm{ext}}^{1/2} \propto t^{-3/2}$ [where B is the comoving magnetic field amplitude and ρ_{ext} is the external density; the same as the scaling of $\nu_m = \nu_{\mathrm{syn}}(\gamma_{\min})$ for $p > 2$], so $\epsilon_X \propto t^{3(2-p)/4}$. This time dependence is the same as for $p > 2$, but in this case ϵ_X increases with time whereas for $p > 2$ it decreases. Similarly, $F_X \propto t^{-(3p-2)/4}$, just as for $p > 2$, except that for $p < 2$ this corresponds to a decay rate flatter than t^{-1}. One possible difference between the two cases is that for $p < 2$ and $\nu_c < \nu_{\max}$ radiative losses are not always negligible since most of the energy in the electrons is at $\gamma_e \sim \gamma_{\max}$ and is therefore radiated away. Thus, unless $\epsilon_e \ll 1$, radiative losses could be significant and would tend to steepen the flux decay rate and make it harder to achieve a flattish light curve. We also note that the X-ray spectral slope for $p < 2$ [assuming $\nu_X > \max(\nu_c, \nu_m)$] is $\beta_X = p/2 < 1$, which in many cases is inconsistent with the observed value (Nousek et al. 2006), so this explanation of the shallow decay of F_X might only apply to a subset of the sources (see fig. 8 of Nousek et al. 2006).

An alternative possibility for $\epsilon_X(t)$ to increase with t is for either one (or both) of the microphysical parameters ϵ_e and ϵ_B to increase with time. Using the dependence of ν_m on these parameters (e.g. Sari et al. 1998), we find that, for $p > 2$, $\epsilon_X \propto \epsilon_e^{(p-1)}\epsilon_B^{(p-2)/4}$ when $\epsilon_e \ll \epsilon_B$ and $\epsilon_X \propto \epsilon_e^{(p-3/2)}\epsilon_B^{p/4}$ when $\epsilon_e \gg \epsilon_B$. This applies when parametrizing in terms of $E_{\mathrm{k,iso}}$ (which is not measured directly), whereas a

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[8] In the case of a very pure Poynting flux, the escaping energy $E_{*,\mathrm{iso}}$ is the fraction that did not dissipate. Therefore, while $E_{\mathrm{iso},0}$ increases by a factor of $(1 - \epsilon_*)^{-1}$, the dissipated energy $\epsilon_{\mathrm{dis,GRB}} E_{\mathrm{iso},0}$ remains unchanged (assuming other efficiencies do not change), and thus $\epsilon_{\mathrm{dis,GRB}}$ decreases by a factor of $(1 - \epsilon_*)^{-1}$.

[9] It is in principle possible that ϵ_X could increase with time on account of its dependence on Y (see equation 4) even if ϵ_e and ϵ_B remained constant and p were > 2, given that Y decreases with time in the slow-cooling regime ($\nu_m < \nu_c$) that is relevant for $p > 2$. However, one can show that, to be relevant during the early-afterglow phase, this would require unrealistically high values of p.

parametrization in terms of $L_{\rm X,iso}$ (which is measured directly) yields $\epsilon_{\rm X} \propto \epsilon_{\rm e}^{4(p-1)/(p+2)}\, \epsilon_B^{(p-2)/(p+2)}$ when $\epsilon_{\rm e} \ll \epsilon_B$ and $\epsilon_{\rm X} \propto \epsilon_{\rm e}^{2(2p-3)/(p+2)}\, \epsilon_B^{p/(p+2)}$ when $\epsilon_{\rm e} \gg \epsilon_B$ (see equation 2). Table 3 of Nousek et al. (2006) provides the values of $\Delta\alpha$ – the change in the temporal decay index across the break in the light curve at $t_{\rm break,2}$ – for nine *Swift* GRBs in which it could be measured reliably, and shows that typically $0.5 \lesssim \Delta\alpha \lesssim 1$. In our context, if $\epsilon_{\rm e} \propto t^{\alpha_{\rm e}}$ and $\epsilon_B \propto t^{\alpha_B}$ at $t < t_{\rm break,2}$, then attributing the flattish decay phase to a growth of $\epsilon_{\rm e}$ and/or ϵ_B with time requires (in the limit $\epsilon_{\rm e} \gg \epsilon_B$) that $\Delta\alpha = 2(2p-3)\alpha_{\rm e}/(p+2) + \alpha_B p/(p+2)$. For $2 < p < 3$, $0.5 < p/(p+2) < 0.6$ and $0.5 < 2(2p-3)/(p+2) < 1.2$. Therefore, for $p \approx 2$, $\alpha_{\rm e} + \alpha_B \approx 2\Delta\alpha \sim 1$–$2$ and for $\Delta\alpha \sim 0.5$, a roughly linear growth with time of either $\epsilon_{\rm e}$ or ϵ_B (or of their product) is required. For $p \sim 2.6$ and $\Delta\alpha \sim 1$, a linear growth of $\epsilon_{\rm e}$ and a constant ϵ_B would work (i.e. $\alpha_{\rm e} \approx 1$ and $\alpha_B = 0$). If $Y \sim 1$ (rather than $Y \gg 1$ or $Y \ll 1$), the dependence of $\epsilon_{\rm X}$ on $\epsilon_{\rm e}$ and ϵ_B is no longer a simple power law, requiring a similarly non-trivial dependence of $\epsilon_{\rm e}$ and/or ϵ_B on the observed time t (insofar as $F_{\rm X}$ is indeed a pure power law in t during the flattish decay phase). A physical scenario will need to account both for this behaviour and for why the time-evolution of the microphysical parameters effectively terminates at $t_{\rm break,2}$.

The magnetic energy parameter ϵ_B could reflect either the structure of the ambient magnetic field or post-shock field-amplification processes. In the former case, an increase of ϵ_B with time could be caused by an increase of the magnetization parameter $\sigma = B_{\rm ext}^2/4\pi\,\rho_{\rm ext}\,c^2$ of the ambient gas with distance from the source, which might occur in certain GRB progenitor models (e.g. Königl & Granot 2002). In the latter case one cannot at present identify a natural reason for ϵ_B to increase during the early-afterglow phase, but future theoretical advances (see, e.g. Medvedev et al. 2005) might make it possible to study the evolution of shock-generated magnetic fields over time-scales that are long enough to address this question. The value of $\epsilon_{\rm e}$ might also be linked to the changing shock parameters (in particular, the shock Lorentz factor). However, in this case, again, our current level of understanding does not allow us to make a specific prediction.

The possible increase of the afterglow radiation efficiency with time during the early phases of the X-ray light curve may also help to lower the estimate of the γ-ray radiative efficiency and thereby alleviate the constraints on the internal-shocks model. If at early times ($t < t_*$) $\epsilon_{\rm X}$ increased with t, then (by equation 15) $E_{\rm k,iso,0}$ would be underestimated, and therefore the parameter f and the value of ϵ_γ would be overestimated. The prompt emission from internal shocks could in principle be observable even if the very early radiation from the external (forward) shock is weak because of a low value of ϵ_B or of a possible suppression of $\epsilon_{\rm e}$ when the shock Lorentz factor is high. This is because the value of ϵ_B in the internal shocks might be determined by a comparatively strong magnetic field advected from the central source (e.g. Spruit, Daigne & Drenkhahn 2001) and because (in contrast to the forward shock at the deceleration time, whose Lorentz factor is $\gtrsim 10^2$) the internal shocks are only mildly relativistic. If the afterglow emission efficiency recovers its commonly assumed value (equation 4) at $t \gtrsim t_{\rm break,2}$ then one could in principle have $f \sim 1$, $\kappa \sim 1$, and $\epsilon_\gamma \sim 0.5$. In this case, $E_{\rm k,iso}$ remains constant while the flattish flux decay is caused by an increase of $\epsilon_{\rm X}$ with time. However, if even at late times the afterglow efficiency is only a fraction $\delta \ll 1$ of its standard value [where δ might, for example, correspond to the fraction $\xi_{\rm e}$ of the electrons behind the forward shock that are accelerated to relativistic energies with a power-law distribution, which could conceivably be $\ll 1$ (Eichler & Waxman 2005)] then $\kappa \sim \delta\kappa_{\rm standard} \sim \delta \ll 1$ and $\epsilon_\gamma \sim \delta$ (for $\epsilon_{\rm obs,GRB} \sim 1$). In this case $E_{\rm k,iso}$ is again constant with time and the flattish decay phase is due to $\epsilon_{\rm X}$ increasing with time; however, the value of $\epsilon_{\rm X}$ at late times is δ times its standard value (given by equation 4) and the true kinetic energy in the afterglow shock is a factor $\delta^{-1} \gg 1$ larger than the usual estimate of $\sim 10^{51}$ erg.

An alternative to the explanation of the flattish segment of the light curve in terms of a prolonged energy injection into the forward shock of a basic jet model is the possibility (in a generalized jet model) of distinct *spatial* components, some of which only contribute to the afterglow emission at later times. This situation could arise if (1) some of the ejecta have lower initial Lorentz factors that result in longer deceleration times, so only a small fraction of the injected energy is transferred to the shocked external medium early on, or if (2) radiation from components that do not move along our line of sight is strongly beamed away from us at early times, becoming visible only later on when deceleration causes the respective beaming cones to widen. Case (1) is exemplified by the two-component jet model considered in the next section, whereas case (2) might be realized in the 'patchy shell' model of GRB sources (Kumar & Piran 2000) and in the 'anisotropic afterglow efficiency' scenario outlined by Eichler & Granot (2006).[10]

4 THE TWO-COMPONENT JET MODEL: A CASE STUDY

As pointed out in Section 1, the two-component jet model was originally invoked by Peng et al. (2005) as a possible way of alleviating the pre-*Swift* constraints on the γ-ray emission efficiency. In this section, we interpret the *Swift* results in the context of this scenario, using again the parameters κ and f introduced in Section 2 and affirming some of the conclusions reached in that section. In this discussion, we assume that the standard afterglow theory applies and that the microphysical parameters do not change with time. This formulation is used to demonstrate that the recent observations are inconsistent with parameter values of the two-component jet model that could lead to a lower inferred magnitude for ϵ_γ, reinforcing the result that, in the context of the standard afterglow theory, the *Swift* measurements have tightened the constraints on the prompt-emission efficiency. We go on to show, however, that the two-component model can nevertheless account for the X-ray afterglow light curve of GRB sources, including the flattish early-time segment.

The generic two-component jet model consists of a narrow and initially highly relativistic (conical) outflow from which the prompt emission originates, and a moderately relativistic flow that decelerates at a significantly later time and occupies a wider (coaxial) cone. The

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[10] Case (2) might also be relevant to the two-component jet model if the line of sight to the observer lies within the solid angle subtended by the narrow component but outside the inner edge of the wide component.

narrow component is the source of the prompt emission (which is observed as a GRB if the observer's line of sight lies within, or very close to, its opening solid angle), whereas the wide component only makes a contribution to the afterglow emission (which becomes observable after it decelerates). Letting E_i, $E_{k,i}$, θ_i and η_i stand for the total initial energy, initial kinetic energy, opening half-angle and Lorentz factor of the two components (with i = n, w corresponding to *narrow* and *wide*, respectively), we have

$$E_{\gamma,\text{iso}} = \epsilon_\gamma E_{n,\text{iso}} = \epsilon_\gamma E_n \frac{2}{\theta_n^2}. \tag{16}$$

The parameter κ defined in equation (3) can be expressed in this context by the relation

$$E_{\gamma,\text{iso}}^{\text{obs}} \frac{\theta_w^2}{2} = \epsilon_{\text{obs,GRB}} \epsilon_\gamma E_n \left(\frac{\theta_w}{\theta_n}\right)^2 \approx \kappa(E_{k,n} + E_{k,w}), \tag{17}$$

where $E_{k,n} = (1 - \epsilon_\gamma)E_n$ and $E_{k,w} \approx E_w$. Equation (17) incorporates the fact that the kinetic energy responsible for the late-afterglow emission is determined empirically by assuming a jet of half-opening angle θ_w, or equivalently that the isotropic-equivalent kinetic energy inferred from late-time afterglow corresponds approximately to the total kinetic energy over the fraction of the total solid angle occupied by the wide component.

The energy that determines the early-afterglow phase is

$$E_{\text{early,iso}} = E_{k,n} \frac{2}{\theta_n^2} = E_{k,n,\text{iso}}, \tag{18}$$

whereas the late-afterglow phase is determined by

$$E_{\text{late,iso}} = (E_{k,n} + E_{k,w}) \frac{2}{\theta_w^2}. \tag{19}$$

The evidence from the *Swift* observations that we do not see the early-afterglow emission above the rapidly decaying tail of the prompt emission, and that even when it shows up it is rather weak, implies that the early-isotropic kinetic energy $E_{\text{early,iso}}$ is a factor $f \sim 10$ smaller than the late-isotropic kinetic energy $E_{\text{late,iso}}$, where we reintroduced the parameter f defined in equation (3). This implies that

$$\frac{E_{k,n}}{\theta_n^2} = \frac{E_{k,n} + E_{k,w}}{f\theta_w^2} \approx \frac{E_{k,w}}{f\theta_w^2}, \tag{20}$$

which in turn implies that

$$E_{k,n} < E_{k,w}. \tag{21}$$

Hence, we can delete the term involving $E_{k,n}$ from the right-hand side of equation (17) and obtain

$$\frac{\epsilon_\gamma \epsilon_{\text{obs,GRB}}}{(1 - \epsilon_\gamma)} = \kappa \frac{E_{k,w}}{E_{k,n}} \left(\frac{\theta_n}{\theta_w}\right)^2 = \kappa \frac{E_{k,w,\text{iso}}}{E_{k,n,\text{iso}}}. \tag{22}$$

Expressing now $E_{k,w}$ in terms of $E_{k,n}$ using this last equation and substituting into equation (20) (which becomes simply $f \approx E_{k,w,\text{iso}}/E_{k,n,\text{iso}}$ when one neglects $E_{k,n}$ on the r.h.s.), we rediscover the first relation in equation (3), which can be expressed in the form of equation (13) to yield $\epsilon_\gamma \sim 0.9$ for $f \sim 10$ and $(\kappa/\epsilon_{\text{obs,GRB}}) \sim 1$.

Equation (22) with $\epsilon_\gamma \epsilon_{\text{obs,GRB}} \approx 1$ implies that $\kappa E_{k,w}/\theta_w^2 \approx E_n/\theta_n^2$, and hence, given that $E_{k,w} \approx E_w$ and taking κ to be ~ 1, that the narrow and wide jet components initially have comparable isotropic-equivalent energies. The ratio of true energies of the two components is initially $E_w/E_n \approx \kappa^{-1} (\theta_w/\theta_n)^2 \sim (4\text{--}9)\kappa^{-1}$ for reasonable ratios of the opening half-angles. This ratio further increases by a factor of $(1 - \epsilon_\gamma)^{-1} \approx \kappa f$ (see equation 13) during the prompt-emission phase.

As argued by Peng et al. (2005), this model could reduce the inferred value of ϵ_γ if $E_{k,w}/E_{k,n} > 1$ *and* $E_{k,w,\text{iso}}/E_{k,n,\text{iso}} < 1$. However, the *Swift* results, as expressed by equation (20), demonstrate that the latter ratio is equal to $f \sim 10$, and hence that this possibility is not realized in practice. Peng et al. (2005) also discussed the ability of this scenario to account for certain 'bumps' in the late-afterglow light curve of several pre-*Swift* GRB sources. We now show that this model can similarly account for the early-time flattening of the X-ray light curve. The required model parameters are, however, distinctly different from those considered by Peng et al. (2005); in fact, the fits that we obtain reinforce the conclusion that $E_{k,w,\text{iso}}/E_{k,n,\text{iso}}$ ($\approx f$) must be $\gg 1$.

The Lorentz factor of the narrow jet component does not directly enter into our modelling of the light curve, although the usual 'compactness' arguments for the prompt emission (e.g. Lithwick & Sari 2001) imply that its value is $\eta_n \gtrsim 10^2$. Our interpretation of the flattish segment of the light curve in the context of this scenario is that it largely corresponds to the emergence of the wide component around its deceleration time $t_{\text{dec,w}} \propto (E_{k,w,\text{iso}}/n_{\text{ext}})^{1/3}\eta_w^{-8/3}$, where n_{ext} is the particle density of the ambient medium at the deceleration radius. Owing to the strong dependence of t_{dec} on η, we can constrain the value of η_w within a factor of 2 or so.

Fig. 1 shows a tentative fit to the X-ray light curve of GRB 050315 with the two-component jet model, whereas Fig. 2 demonstrates the effect of modifying the model parameters. The extended flat segment of the light curve together with the requirement that the contribution from the narrow component at $t/(1 + z) \sim 200\,\text{s}$ does not overproduce the observed flux imply $f \approx E_{k,w,\text{iso}}/E_{k,n,\text{iso}} \gtrsim 30$. This suggests that $f \sim f_{\text{max}} \sim 30$ (see Table 2) and that both f_{min} and f_{max} might have been somewhat underestimated for this event [since the fit to the data should produce $f \gtrsim f_{\text{min}}$, suggesting that $f_{\text{min}} \sim 30$ for GRB 050315, which is higher than the value of $f_{\text{min}} = 11$ derived by Nousek et al. (2006) and shown in Table 2. For our fiducial parameter values ($\epsilon_e = 0.1$, $\epsilon_B = 0.01$, and $p \approx 2.0\text{--}2.1$; see Table 2 and Fig. 1) we find (using equation 5) $\kappa \approx 0.034$ (see Table 2). For this value of κ the product κf is ~ 1, which implies (from equation 13 with $\epsilon_{\text{obs,GRB}} \sim 1$) that
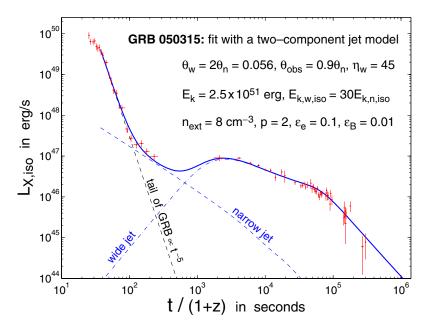
Figure 1. Tentative fit to the X-ray light curve of GRB 050315 (from Nousek et al. 2006) with the two-component jet model. The numerical code used to calculate the light curve is essentially model 1 of Granot & Kumar (2003), which neglects the lateral spreading of the jet and includes SSC emission. In addition to the total light curve (thick solid line), we also show the separate contributions of the different components: the tail of the prompt emission ($\propto t^{-5}$), the narrow outflow and the wide outflow. Here $E_k = E_{k,w} + E_{k,n}$ is the total kinetic energy of the two components. The narrow and wide components occupy the non-overlapping ranges $\theta < \theta_n$ and $\theta_n < \theta < \theta_w$, respectively, in the polar angle θ (measured from the symmetry axis); θ_{obs} is the viewing angle with respect to this axis.
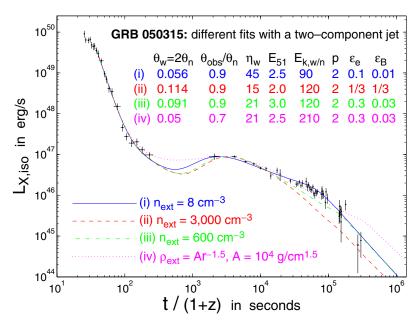

Figure 2. The effects of varying the two-component jet model parameters with respect to those of the fit shown in Fig. 1 (reproduced in this figure by the solid line). Here $E_{51} = E_k/(10^{51}\,\mathrm{erg})$ (with $E_k = E_{k,w} + E_{k,n}$), whereas $E_{k,w/n}$ denotes the ratio $E_{k,w}/E_{k,n}$ of the kinetic energies of the wide and narrow components.

$\epsilon_\gamma \sim 0.5$. The fit to the data shown in Fig. 1 incorporates an SSC component[11] whose contribution in this case turns out to be comparable to that of the synchrotron component at $t_* = 10\,\mathrm{h}$, resulting in a decrease by a factor of $\sim$2 in the estimate of $E_{k,iso}(10\,\mathrm{h})$ (and in a corresponding increase in the estimate of κ) in comparison with the values inferred by using equation (5) (which only takes account of the synchrotron contribution). The actual numerical difference between these two estimates is, in fact, slightly larger (a factor of $\sim$2.5), reflecting the fact that

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[11] This SSC component arises from inverse-Compton scattering of low-energy photons that places the upscattered photons in the X-ray range. It should not be confused with the inverse-Compton cooling of electrons that emits synchroton X-rays, which is discussed in Section 2.2.

the fit employs a numerical scheme that is not identical to equation (5). All in all, the model fit shown in Fig. 1 yields $f \approx E_{\mathrm{k,w,iso}}/E_{\mathrm{k,n,iso}} = 30$, $\kappa \approx E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}/E_{\mathrm{k,w,iso}} \approx 0.086$ and $\kappa f \approx E_{\gamma,\mathrm{iso}}^{\mathrm{obs}}/E_{\mathrm{k,n,iso}} \approx 2.6$, which implies $\epsilon_\gamma \approx 0.72$ for $\epsilon_{\mathrm{obs,GRB}} = 1$.

One realization of a two-component relativistic outflow of the type considered here is an initially neutron-rich, hydromagnetically accelerated jet (Vlahakis, Peng & Königl 2003). In this picture, the narrow component consists of the originally injected protons that are gradually accelerated to η_n; the injected neutrons decouple from the original proton component when the jet Lorentz factor reaches η_w and eventually decay to form a distinct (wide) proton component. In the illustrative solution presented in Vlahakis et al. (2003), $\eta_w \approx 15$. The *dashed curve* in Fig. 2 demonstrates that, to be consistent with the data, this value of η_w requires a very high external density, $n_{\mathrm{ext}} \sim 10^{3.5}$ cm^{-3} (although even for this value the fit is not as good as for the parameters adopted in Fig. 1). This can be understood from the parameter dependence of $t_{\mathrm{dec,w}}/(1 + z)$, for which the model fit implies a value of $\sim 2 \times 10^3$ s. To reproduce the observed flux, the value of $E_{\mathrm{k,w,iso}}$ cannot be too low. In fact, in the model fit represented by the dashed curve we have adopted the lowest possible value of this quantity, corresponding to an equipartition between the electron and magnetic field energy densities. With the values of $t_{\mathrm{dec,w}}$ and $E_{\mathrm{k,w,iso}}$ thus fixed, the inferred external density becomes very sensitive to the value of the Lorentz factor ($n_{\mathrm{ext}} \propto \eta_\mathrm{w}^{-8}$). Intermediate options, with a somewhat higher Lorentz factor ($\eta_\mathrm{w} = 21$) are also shown in Fig. 2, both for the uniform-density case used in the previous fits (*dash–dotted curve*) and for an $r^{-1.5}$ density profile (where r is the distance to the source; *dotted curve*). We interpret the break in the light curve of GRB 050315 at $t/(1 + z) \sim 1$ d as the jet break time of the wide component, $t_{\mathrm{jet,w}}/(1 + z)$. Since the jet break time is particularly sensitive to the value of the opening half-angle, $t_{\mathrm{jet,w}} \propto (E_{\mathrm{k,w,iso}}/n_{\mathrm{ext}})^{1/3}\theta_\mathrm{w}^{8/3}$, this allows us to constrain the value of θ_w. Given that $t_{\mathrm{jet,w}}/t_{\mathrm{dec,w}} \propto (\eta_\mathrm{w}\theta_\mathrm{,w})^{8/3}$, it is seen that any reduction in the fitted value of η_w requires an increase in θ_w by a similar factor.

As was already noted in Peng et al. (2005), another realization of this type of an outflow is potentially provided by a relativistic, baryon-poor jet, which is driven electromagnetically along disc-anchored magnetic field lines that thread the horizon of a rotating black hole, and which is 'contaminated' by neutron diffusion from a baryon-rich disc wind. In the original version of this scenario, which was proposed by Levinson & Eichler (1993, 2003; see also van Putten & Levinson 2003) and recently studied numerically by McKinney (2005a,b), the narrow and wide components correspond to the baryon-poor and baryon-rich outflows, respectively. However, the predicted Lorentz factor of the disc wind is too low to be consistent with the rather high value of η_w inferred from our model fits. An alternative possibility is that the wide component corresponds to a hydromagnetically accelerated baryon-rich disc outflow of the type modelled by Vlahakis & Königl (2003a,b), whereas the narrow component corresponds (as in the Levinson & Eichler picture) to a baryon-poor Blandford & Znajek (1977)-type outflow.[12] It should be noted that in either one of the above two-component jet realizations the γ-ray emitting component is associated with an initially Poynting-dominated outflow. This could in principle make it possible to account for the relatively high emission efficiency that the *Swift* results seem to imply (see Section 5).

5 SUMMARY AND CONCLUSIONS

We have shown that the γ-ray efficiency implied by the *Swift* observations is model-dependent and can vary over a wide range (from typical values of $\epsilon_\gamma \sim 0.9$ or higher to $\epsilon_\gamma \sim 0.1$ or lower) depending on the adopted model assumptions. The γ-ray efficiency has been expressed in terms of observable quantities (see equations 3 and 13), namely κ and f, where κ relates the γ-ray emission to the late-time afterglow emission (and was therefore available in the pre-*Swift* era) and f relates the early- and late-time afterglow energies (and therefore became available only with the launching *Swift*). We have shown that there is a large uncertainty in the values of both κ and f, which translates into a corresponding uncertainty in the value of the γ-ray efficiency, ϵ_γ. In the following discussion, we make the conservative assumption that we observe most of the emitted energy in γ-rays ($\epsilon_{\mathrm{obs,GRB}} \approx 1$); if a significant fraction of the total radiated energy is emitted outside of the observed photon energy range (e.g. at higher energies) then this would increase the required value of ϵ_γ (see equation 13).

The kinetic energy of the afterglow shock has been estimated in the pre-*Swift* era using broad-band afterglow fits for a small number of GRB sources that had the best available broad-band afterglow data, yielding a typical value of $\kappa \sim 1$ with a large scatter of almost an order of magnitude (see Section 2.3). Substituting the values of the microphysical parameters (ϵ_e, ϵ_B and p) that were derived from these fits into our equations for $E_{\mathrm{k,iso}}$ generally yielded somewhat lower values of κ. Finally, using our equations with fiducial values of the microphysical parameters ($\epsilon_\mathrm{e} = 0.1$, $\epsilon_B = 0.01$ and $p = 2.2$) gives a typical value of $\kappa \sim 0.1$–0.2, both for the pre-*Swift* and the *Swift* GRB samples that we use, with a moderate scatter. Specifically, for our pre-*Swift* (*Swift*) sample, $\langle \log_{10} \kappa \rangle = -0.75 \, (-0.82)$ corresponding to $\kappa = 0.18 \, (0.15)$ and $\sigma_{\log_{10} \kappa} = 0.60 \, (0.63)$. Obviously, the choice of fiducial values for the microphysical parameters is somewhat arbitrary and affects the resulting value of κ. Higher values of the microphysical parameters ϵ_e and ϵ_B (e.g. $\epsilon_\mathrm{e} \approx 0.3$, $\epsilon_B \approx 0.08$) are required in order to obtain an average value of $\log \kappa$, using our equations, similar to that derived from pre-*Swift* broad-band afterglow fits. Altogether, there is almost an order-of-magnitude uncertainty in the typical value of κ (which ranges from ~ 0.1–0.2 to ~ 1).

Even if we adopt the high typical values of κ (~ 1) inferred from pre-*Swift* afterglow fits, it is important to keep in mind that these values have been estimated on the basis of the standard assumptions of afterglow theory. Changing these assumptions would modify the inferred value of κ for the same fits. For example, as pointed out by Eichler & Waxman (2005), if only a fraction $\xi_\mathrm{e} < 1$ of the electrons are accelerated to relativistic energies with a power-law distribution, then there is a degeneracy where the same observable quantities are obtained for $\epsilon_\mathrm{e} \to \xi_\mathrm{e}\epsilon_\mathrm{e}$, $\epsilon_B \to \xi_\mathrm{e}\epsilon_B$,

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[12] It is in principle also conceivable that the baryon-poor outflow could develop an internal structure that would correspond to the two outflow components considered here, but there are at present no quantitative results to support this conjecture.

$n_{ext} \rightarrow \xi_e^{-1} n_{ext}$ and $E_{k,iso} \rightarrow \xi_e^{-1} E_{k,iso}$. Since this increases the inferred value of $E_{k,iso}$ by a factor of ξ_e^{-1}, κ is reduced by the same factor in comparison with the estimate from the standard theory (which uses $\xi_e = 1$).

An alternative way of reducing the inferred value of ϵ_γ was proposed by Peng et al. (2005) in the context of the pre-*Swift* observations. Specifically, they considered a two-componet outflow model with parameters that effectively corresponded to the parameter *f* having a value <1. This parameter choice appears to be inconsistent with *Swift*'s detection of an early flattish decay phase in the X-ray light curve, which, when interpreted in the context of the standard afterglow theory as arising from a gradual increase with time of $E_{k,iso}$, typically implies $f \sim$ 10 (and, in some cases, values of *f* that are as high as $\sim 10^2$–10^3). It is worth noting in this connection that the existence of a two-component GRB jet model can be plausibly expected on various theoretical grounds and has been suggested independently on the basis of fits to several pre-*Swift* afterglows (see discussion in Peng et al. 2005). Furthermore, the *Swift* observations by no means rule out this model, although they can be used to constrain its parameters. We have illustrated this fact through the fit to the X-ray light curve of GRB 050315 that we performed in Section 4 within the framework of this model. This fit has yielded plausible ranges for the kinetic energies and opening angles of the two components as well for as the Lorentz factor of the dominant (wide) component. A key conclusion from this fit is that the kinetic energy of the wide component is much larger than that of the narrow one ($E_{k,w}/E_{k,n} \sim 10^2$). Physically, the narrow and wide components could conceivably correspond to a baryon-poor, black hole driven outflow (Levinson & Eichler 1993, 2003) and a baryon-rich, disc-driven outflow (Vlahakis & Königl 2003a,b), respectively, although this remains to be demonstrated. We also note that the two-component jet parameters derived in Section 4 were based on standard assumptions; they could change if the underlying assumptions (involving, for example, the values and time constancy of the microphysical parameters) were altered.

In conjunction with the $\kappa \sim 1$ estimates of the standard broad-band afterglow fits, values of $f \gtrsim 10$ imply γ-ray radiative efficiencies $\epsilon_\gamma \gtrsim 0.9$. Such high efficiencies would be essentially impossible to achieve in any scheme, such as the internal-shocks model, that tapped the kinetic energy of the outflow for radiation. An alternative possibility that has been discussed in the literature is the direct transfer of Poynting flux (which evidently is also a major contributor to the flow acceleration – e.g. Drenkahn & Spruit 2002; Vlahakis & Königl 2003a) into non-thermal radiation (e.g. Thompson 1994; Usov 1994). It is at present unclear how to assess the efficiency of this process. There are two generic possibilities: dissipative fronts and magnetic reconnection sites. The first option corresponds to overtaking collisions of magnetically dominated relativistic streams and is not expected to result in high radiative efficiencies (e.g. Levinson & van Putten 1997; Romanova & Lovelace 1997). The second case would require magnetic field orientation reversals and would most naturally arise in a pulsar-type outflow from a rapidly rotating neutron star (e.g. Spruit et al. 2001). In this case it is, however, still unclear whether radiative efficiencies $\gtrsim 0.9$ could be attained even under the most favourable assumptions about the field reconnection rate (Drenkahn & Spruit 2002), and it has in fact been suggested that the reconnection rate might be self-limiting (Lyubarsky & Kirk 2001).

An 'intermediate' situation could prevail if $f \sim 10$ but $\kappa \sim 0.1$, reflecting the possibility that κ was overestimated by the pre-*Swift* afterglow fits (perhaps because some of the assumptions of the standard theory do not hold – e.g. $\xi_e \sim 0.1$ rather than $\xi_e = 1$). Alternatively, κ could be ~ 1 but $f = 1$, corresponding to the early flattish decay phase reflecting an increase with time of the X-ray afterglow efficiency ϵ_X (due, e.g. to *p* being <2 or to an increase with time of ϵ_e or ϵ_B; see Section 3) rather than an early increase in $E_{k,iso}$. In either one of these cases the inferred γ-ray radiative efficiency would be reduced to $\epsilon_\gamma \sim 0.5$. Although this value is less extreme than the estimate discussed in the preceding paragraph, it is worth noting that it is still fairly restrictive for the internal-shocks model, in which it could potentially be attained only if *all* of the following conditions (already summarized in Peng et al. 2005) are satisfied: (1) the ratio between the maximum and minimum initial Lorentz factors of the ejected shells is large enough ($\gtrsim 10$); (2) the distribution of initial Lorentz factors is sufficiently non-uniform; (3) the shells are approximately of equal mass and their number is large enough ($\gtrsim 30$) and (4) the fraction of the dissipated energy that is deposited in electrons and then radiated away is sufficiently high ($\epsilon_{e,GRB} \epsilon_{rad,GRB} \gtrsim 0.5$), with a similar constraint applying to the fraction $\epsilon_{obs,GRB}$ of the radiated energy that is emitted as the observed γ-rays (see Beloborodov 2000; Kobayashi & Sari 2001). Only if *both* $f \sim 1$ and $\kappa \sim 0.1$ were satisfied (which could occur, for example, if ϵ_e or ϵ_B initially increased with time and ξ_e were ~ 0.1) would the inferred value of ϵ_γ drop to ~ 0.1 and be compatible with the values that are expected to arise under less-constrained circumstances in the internal-shocks model.

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