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GRB 990123: reverse and internal shock flashes and late afterglow behaviour

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ABSTRACT

The prompt ($t \lesssim 0.16\,\mathrm{d}$) light curve and initial 9th-magnitude optical flash from GRB 990123 can be attributed to a reverse external shock, or possibly to internal shocks. We discuss the time decay laws and spectral slopes expected under various dynamical regimes, and the constraints imposed on the model by the observations, arguing that they provide strongly suggestive evidence for features beyond those in the simple standard model. The longer term afterglow behaviour is discussed in the context of the forward shock, and it is argued that, if the steepening after 3 d is due to a jet geometry, this is likely to be a result of jet-edge effects, rather than sideways expansion.

Key words: shock waves – cosmology: miscellaneous – gamma-rays: bursts.

1 INTRODUCTION

The observations of GRB 990123 (Akerlof et al. 1999; Galama et al. 1999; Kulkarni et al. 1999) not only pose constraints on the amount of gamma-ray beaming needed from a stellar-mass progenitor in the absence of lensing, but also provide an interesting test of the canonical fireball shock afterglow model. A simultaneous optical flash of 9th magnitude from a burst at cosmological redshifts was discussed more than two years ago by Mészáros & Rees (1997) (models a2, a3; see also Mészáros & Rees 1993), in terms of the reverse shock that accompanies the blast wave. This early optical flash is expected to start at the time of the gamma-ray trigger, and to decay faster than the better known radiation from the forward blast wave, which starts out weaker but dominates the longer duration optical afterglow. A similar prediction based on the reverse shock was made in the light of more recent studies by Sari & Piran (1999a), and more specifically discussed by them in the context of the observations of GRB 990123 (Sari & Piran 1999c).

A different origin for a simultaneous optical flash is possible from internal shocks (Mészáros & Rees 1997: e.g. model b2 of that paper). Internal shock optical flashes are of additional interest because, as pointed out by Fenimore, Ramirez-Ruiz & Wu (1999), the gamma-ray light curve of GRB 990123 (as well as those of several other bursts) appears to be incompatible with the gamma-rays coming from a single external shock, since the gamma-ray pulses in the second half of the burst are not appreciably longer than in the first half. It is, of course, possible that the gamma-rays arise in internal shocks, which are thought to be exempt from such problems, while the optical afterglows may arise from the external and the associated reverse shocks. However, in the light of the need for internal shocks, it is interesting to investigate the implications of the early afterglow observations at various wavelengths including optical, for both external and internal shocks.

We also discuss the longer term behaviour of the afterglow, and the causes for the optical light curve flattening after 0.16 d. The likely interpretation of the light curve after this time is that it is due to the forward shock or blast wave. The discrepancy pointed out between the observed time decay slope and the spectral index within the context of the simple standard model can be resolved by invoking the simplest realistic extensions to this model, and we discuss several specific possibilities (Section 4). We also indicate that the steepening of the light curve to $\propto t^{-1.65}$–$t^{-1.8}$ after about 3 d, if the result of a jet geometry, is more likely to be due to the effects of beginning to see the edge of the jet. This effect occurs before, and its effects fit the steepening better than, the alternative sideways expansion interpretation (Section 5).

2 OPTICAL FLASH FROM REVERSE EXTERNAL SHOCKS

The reverse shock acompanying the forward blast wave gives a prompt optical flash of the right magnitude, with reasonable energy requirements of no more than a few $\times 10^{53}$ erg isotropic (Mészáros & Rees 1997). The time decay constants calculated in that paper for models a2 and a3 were affected by an error, which we correct here; we consider also a more generic prescription for the dynamics and the magnetic field behaviour.

For a general evolution of the bulk Lorentz factor $\Gamma \propto r^{-g}$ with radius, the radius and observer time t are related through $r \sim ct\Gamma^2$, or

$$\Gamma \propto r^{-g} \propto t^{-g/(1+2g)}, r \propto t^{1/(1+2g)}. \quad (1)$$

In the usual Blandford & McKee (1976) impulsive solutions, the 'adiabatic' case is $g = 3/2$, $\Gamma \propto r^{-3/2}$, $r \propto t^{1/4}$; the 'radiative' case is $g = 3$, $\Gamma \propto r^{-3}$, $r \propto t^{1/7}$ and in the similarity limit $g = 7/2$, $\Gamma \propto t^{-7/2}$, $r \propto t^{1/8}$. More general values of g occur if the injection is non-uniform (Rees & Mészáros 1998) or anisotropic, or the external medium is inhomogeneous (Mészáros, Rees & Wijers 1998). Strictly speaking, Γ is the bulk Lorentz factor of the forward shocked material, and may be used also for the contact discontinuity. It is only a very rough approximation for the Lorentz factor of the reverse shock. Assuming that the latter approximation is valid, the comoving width, volume and particle density of the ejecta, after it has been traversed by the reverse shock, evolve with

$$\Delta R \sim r/\Gamma \propto r^{1+g}, \quad V' \sim n'^{-1}_{\mathrm{ej}} \propto r^2 \Delta R \propto r^{3+g}. \quad (2)$$

We consider, in the reverse shocked gas, two possibilities for the comoving magnetic field evolution: one is flux-freezing, $B' \propto V'^{-2/3}$, and the other is that the comoving field in the reverse shocked gas remains in pressure equilibrium with the forward shocked gas, $B' \propto \Gamma$. For flux-freezing (pressure equilibrium) we have then

$$B' \propto r^{-(6+2g)/3} (\text{or} \propto r^{-g}). \quad (3)$$

The energy density ε' and the electron random Lorentz factor γ in the reverse shocked gas follow from $\varepsilon' \propto V'^{-4/3} \propto r^{-(12+4g)/3}$ (or $\propto \Gamma^2 \propto r^{-2g}$) and $\gamma \propto \varepsilon'/n'_{ej} \propto r^{-(3+g)/3} (\propto r^{3-g})$. The synchrotron peak energy in the observer frame is then

$$\nu_{\mathrm{m}} \propto \Gamma B' \gamma^2 \propto r^{-(12+7g)/3} \propto t^{-(12+7g)/(3+6g)}$$
$$(\propto r^{6-4g} \propto t^{(6-4g)/(1+2g)}) \quad (4)$$

in these two magnetic field cases. Considering for simplicity the case where the electron cooling time is long compared with the dynamic expansion time, the comoving synchrotron intensity at the peak frequency is $I'_{\nu_{\mathrm{m}}} \propto n'_{\mathrm{ej}} B' \Delta R \propto r^{-(12+2g)/3} [\propto r^{-(2+g)}]$, and the observer-frame flux is

$$F_{\nu_{\mathrm{m}}} \propto t^2 \Gamma^5 I'_{\nu_{\mathrm{m}}} \propto r^{-(6+5g)/3} \propto t^{-(6+5g)/(3+6g)}$$
$$(\propto r^{-2g} \propto t^{-2g/(1+2g)}). \quad (5)$$

For a photon energy spectral index β ($F_\nu \propto \nu^\beta$) the spectral flux at a given frequency (e.g. optical) expected from the reverse shocked gas is then

$$F_\nu \sim F_{\nu_{\mathrm{m}}} (\nu/\nu_{\mathrm{m}})^\beta \propto F_{\nu_{\mathrm{m}}} \nu_{\mathrm{m}}^{-\beta}$$
$$\propto r^{-[6-12\beta+g(5-7\beta)]/3} \propto t^{-[6-12\beta+g(5-7\beta)]/(3+6g)}$$
$$(\propto r^{-[2g(1-2\beta)+6\beta]} \propto t^{-[2g(1-2\beta)+6\beta]/(1+2g)}), \quad (6)$$

where $\beta = 1/3$ below ν_{m} and $\beta = -(p-1)/2$ above ν_{m} in synchrotron radiation. We have assumed here that the cooling frequency is above the peak frequency, the conditions for the latter appearing to be satisfied in GRB 990123, as pointed out by Sari & Piran (1999c). Under the flux-freezing field behaviour, for an adiabatic case $g = 3/2$ and an electron index $p = 2.5$, the photon spectral index above ν_{m} is $\beta = -3/4$ and we have $F_\nu \propto t^{-81/32} \sim t^{-2.5}$, whereas for $p = 2$ one would have $\beta = -1/2$ and $F_\nu \propto t^{-33/16} \sim t^{-2}$ [while for the similarity case $g = 7/2$, $p = 5/2$ we get, as do Sari & Piran (1999c), $F_\nu \propto t^{-411/192} \sim t^{-2.1}$]. Under the pressure equilibrium field behaviour, for $g = 3/2$, and arbitrary p, we have $F_\nu \propto t^{-3/4}$ (this is a degenerate case where ν_{m} is constant), whereas for $g = 7/2, p = 5/2$ we have $F_\nu \propto t^{-13/8} \sim t^{-1.6}$. The ROTSE observations (Akerlof et al. 1999) give an approximate dependence $\propto t^{-2}$ for about 600 s, which is in rough agreement with the flux-freezing value in either the adiabatic $p = 2$ or similarity $p = 5/2$ case; these two cases cannot be distinguished without spectral information during the ROTSE observations. A different analysis of the earliest data on GRB 990123 (Fruchter et al. 1999) gives the ROTSE slope as $\propto t^{-1.6}$, which is close to the pressure equilibrium similarity solution for $p = 5/2$.

However, note that a decay $\propto t^{-2}$ (or $\propto t^{-1.6}$) can also be obtained in more generic situations than the above, and in particular the spectral slope need not be the only constraint on the decay index, except in the simplest, homogeneous external medium, or single initial Γ (impulsive) model. For an impulsive injection in an inhomogeneous external medium, e.g. $\rho_{\mathrm{ext}} \propto r^{-d}$, one expects in the adiabatic limit $\Gamma \propto r^{-g}$ with $g = (3-d)/2$ (Mészáros, Rees & Wijers 1998). For a more realistic non-uniform injection situation, one expects a range of initial Γ, and in particular for a power-law distribution of Γ in the ejecta (Rees & Mészáros 1998), where the mass injection rate $M(>\Gamma) \propto \Gamma^{-s}$, the effective resulting $\Gamma \propto r^{-g}$ with $g = 3/(2+s)$. The reverse shock relation (6) between photon spectral slope and dynamics could therefore, in the simplest 'post-standard' extensions of the model, depend on e.g. the external density profile, or the mass–energy injection.

3 OPTICAL FLASHES FROM INTERNAL SHOCKS

Internal shocks, just like external shocks, should have a low-energy portion or tail of the spectrum which extends into the optical, starting at the GRB trigger (except possibly in rare cases where the self-absorption frequency extends to the optical). For the same total energy, the flux from the internal shock optical flash is about two orders of magnitude weaker than from the reverse external shock (Mészáros & Rees 1997); however, for beaming factors $\sim 10^{-2}$ as widely considered, they could lead to flashes as bright as 9th magnitude at $z \sim 1$. The simple 'standard internal shock model' (Rees & Mészáros 1994) considers a wind of duration $t_{\mathrm{w}} \sim t_{\mathrm{burst}}$ with $\Gamma \sim$ constant and fluctuations $\Delta\Gamma \sim \Gamma$ over times $t_{\mathrm{v}} < t_{\mathrm{w}}$ producing shocks at a distance $r_{\mathrm{sh}} \sim c t_{\mathrm{v}} \Gamma^2$ for $t \lesssim t_{\mathrm{w}}$.

The most straightforward gamma-ray burst internal shock (e.g. model b1 of Mészáros & Rees 1997) assumes that the gamma-ray break observed around $\sim$100 keV is due to the synchrotron peak energy. The magnetic field required in the internal shocks is not far below equipartition, and this leads to a very short radiative cooling time compared with the expansion time. If the shocks arose from discrete, sharp-edged shells, there would be a very sudden drop of the light curve at all energies (including optical) after internal shocks stop. However, a more realistic situation probably involves smoothly modulated shells, where the shocks weaken gradually after reaching a peak strength. A variety of plausible distributions would then lead to a power-law decay of the optical light.

An alternative internal shock model, which would lead to long decay even in the absence of a gradual fading of the shocks, is obtained in the case where the synchrotron break is at optical energies, and the gamma-ray break is produced by inverse Compton scattered synchrotron photons (model b2 of Mészáros & Rees 1997; see also Papathanassiou & Mészáros 1996). Between the time at which internal shocks (or gamma-rays) stop and the time when deceleration by the external medium starts ($t_{\rm dec} \sim 500[E_{54.6}/(\Omega/4\pi)n_{\rm ext}]^{1/3}\Gamma_{300}^{-8/3}$s), the average bulk Lorentz factor remains approximately unchanged, and the dynamics are described by $\Gamma \propto t^0$, $r \propto t$. Hence the comoving width, volume and particle density of the ejecta evolve as

$$\Delta R \sim r/\Gamma \propto t, V' \sim n'^{-1}_{\rm ej} \sim r^2 \Delta R \propto t^3, \tag{7}$$

and the comoving energy density and electron random Lorentz factor will be

$$\varepsilon' \sim n'\gamma \propto V'^{-4/3} \propto t^{-4}, \quad \gamma \sim \varepsilon'/n' \propto t^{-1}. \tag{8}$$

If the magnetic field is not dynamically dominant, two cases are (i) $B' \propto V'^{-2/3} \propto t^{-2}$, if the field is random, and (ii) $B' \propto B \propto r^{-1} \propto t^{-1}$, if the field is mainly transverse (e.g. inefficient reconnection). Case (i) [(ii)] leads to an observer-frame synchrotron peak flux

$$\nu_{\rm m} \propto \Gamma B' \gamma^2 \propto t^{-4} (\propto t^{-3}). \tag{9}$$

(The first dependence assumes random fields, and the second dependence is for transverse fields.) The comoving intensity is then $I'_{\nu_{\rm m}} \propto n'_{\rm ej} B' \Delta R \propto t^{-4} (\propto t^{-3})$ and

$$F_{\nu_{\rm m}} \propto t^2 \Gamma^5 I'_{\nu_{\rm m}} \propto t^{-2} (\propto t^{-1}), \tag{10}$$

while the flux at a fixed frequency is

$$F_\nu \sim F_{\nu_{\rm m}} \nu_{\rm m}^{(p-1)/2} \sim t^{-2p} [\sim t^{(1-3p)/2}], \tag{11}$$

where one can verify that, at the internal shock radius and above, the optical electrons are in the adiabatic regime, for a wind equipartition parameter $\epsilon_B = 10^{-6}$ and $\gamma \sim 300$. Thus, for $p = 2$, $F_\nu \propto t^{-4} (\propto t^{-5/2})$.

For a magnetically dominated outflow, an interesting third case is (iii) $B' \propto B \propto t^{-1}$, as in (ii) above, but now one might expect the comoving volume and particle density to be dominated by the field evolution (via $B' \propto t^{-1} \propto V'^{-2/3}$), so

$$V' \propto n'^{-1}_{\rm ej} \propto t^{3/2}, \quad \Delta R \sim V'/r^2 \propto t^{-1/2}, \tag{12}$$

$$\varepsilon' \sim n'\gamma \propto V'^{-4/3} \propto t^{-2}, \quad \gamma \sim \varepsilon'/n' \propto t^{-1/2}, \tag{13}$$

$$\nu_{\rm m} \propto \Gamma B' \gamma^2 \propto t^{-2}; \tag{14}$$

the comoving intensity is $I'_{\nu_{\rm m}} \propto n'_{\rm ej} B' \Delta R \propto t^{-3}$ and

$$F_{\nu_{\rm m}} \propto t^2 \Gamma^5 I'_{\nu_{\rm m}} \propto t^{-1}, \tag{15}$$

$$F_\nu \sim F_{\nu_{\rm m}} \nu_{\rm m}^{(p-1)/2} \sim t^{-p}. \tag{16}$$

Adiabatic conditions prevail for $\epsilon_B \sim 1$ and, for example, $\Gamma \gtrsim 300$, $\gamma \sim 100$. For $p = 5/2$ this gives $F_\nu \propto t^{-5/2}$, and for $p = 2$ it gives $F_\nu \propto t^{-2}$.

For the simple internal shock model where the average Γ is constant, model (i) is too steep, but models (ii) and (iii) could fit the observations, with a flat enough electron power law ($p \sim 2$). However, in a realistic model the average Γ of the wind producing internal shocks could vary, and, even for $p > 2$, a decay t^{-2} or $t^{-1.6}$ could be the result, e.g., of an average Γ that increases in time.

4 THE EXTERNAL BLAST WAVE

The standard long-term (as opposed to short-term) afterglow is attributed to the external blast wave (forward shock) evolution, for which the bulk Lorentz factor is described by the same equations (1), and the comoving width and volume of the forward shocked gas are given by equations (2). The comoving density of the forward shocked gas is, however, $n' \sim n_{\rm ext}\Gamma$ and the comoving field is assumed to be some fraction of the equipartition value, $B' \propto \varepsilon_B^{1/2}\Gamma$, while the shocked electron random Lorentzt factor is $\gamma \sim \varepsilon_{\rm e}\Gamma$. As in Mészáros & Rees (1997) (model a1), $\nu_{\rm m} \propto r^{-4g} \propto t^{-4g/(1+2g)}$, for adiabatic electrons $I'_{\nu_{\rm m}} \propto n'B'\Delta R \propto r^{1-g}$, $F_{\nu_{\rm m}} \propto t^2\Gamma^5 I'_{\nu_{\rm m}} \propto r^{3-2g} \propto t^{(3-2g)/(1+2g)}$, and for a spectrum $\propto \nu^\beta$ we have

$$F_\nu \propto r^{3-2g(1-2\beta)} \propto t^{[3-2g(1-2\beta)]/(1+2g)}. \tag{17}$$

Thus in the (impulsive) 'standard model' a simple relation is expected between the time decay index α and spectral slope β, namely $\alpha = [3 - 2g(1 - 2\beta)]/(1 + 2g)$. For example, for the impulsive adiabatic case $g = 3/2$ and one would expect $\alpha = (3/2)\beta$ with $\beta = 1/3, -(p-1)/2$ for synchrotron, or their equivalent for the radiative $g = 3$ or similarity $g = 7/2$ values. Such a simple one-parameter relation between α and β does not appear to hold for the second stage of GRB 990123, where $F_{\rm opt} \propto t^{-1.1}$ (Kulkarni et al. 1999). This, in our view, is a strong indication that 'post-standard' features are present: e.g. a non-homogeneous external medium or an anisotropic outflow (Mészáros, Rees & Wijers 1998). From an observed time decay $F_{\rm opt} \propto t^\alpha$ and an observed spectral slope $F_\nu \propto \nu^\beta$ one can then work backwards to get the effective value of g implied by equation (17) above, and this in turn implies a value of $d = 3 - 2g$ or $s = (3/g) - 2$, thus providing information about the external medium or the injection mechanism. While non-unique, such examples indicate that the time decay index is likely to depend on parameters other than the spectral slope.

5 A JET GEOMETRY? EDGE VERSUS EXPANSION EFFECTS

The temporal decay index after 3 d steepens from about $t^{-1.1}$ to $t^{-1.65\pm0.06}$ (Kulkarni et al. 1999) or $t^{-1.8}$ (Fruchter et al. 1999), at least in the *R* band (*K*-band observations do not show the effect, but this could be due to the higher noise level in this band). This steepening, if real, could be attributed to having a collimated outflow, the effects of which become noticeable after Γ drops sufficiently. If one assumes that the steepening is caused by sideways expansion of the decelerating jet (Rhoads 1997), one might expect a steepening from $t^{-1.1}$ to t^{-p}, where p is the electron index, typically $p = 2.5$, so the change expected would be more than one power of t. However, the edge of the jet begins to be seen when Γ drops below the inverse jet opening angle $1/\theta_{\rm j}$. This occurs well before sideways expansion starts (Panaitescu & Mészáros 1998); the latter is unimportant until the expansion is almost non-relativistic, except for very small angles. So long as $\Gamma > \theta_{\rm j}^{-1}$, the emission that we receive is the same as from a spherically symmetric source, and the effective transverse area is

$A \sim (r_{\parallel}/\Gamma)^2 \propto t^2\Gamma^2$; on the other hand, when $\Gamma < \theta_j^{-1}$ the dependence is $A \sim (r_{\parallel}\theta_j)^2 \propto t^2\Gamma^4$. Note that even after the edge of the jet becomes visible, the outflow is still essentially radial and Γ continues to decay as a power law as before (until sideways expansion sets in). These two additional powers of Γ (e.g. in adiabatic expansion, $\Gamma \propto r^{-3/2} \propto t^{-3/8}$) imply a steepening by $t^{-3/4}$, which matches well the reported steepening by a power of about $t^{-0.7}$. There are of course several other plausible causes for steepening, as discussed by Kulkarni et al. (1999), Mészáros, Rees & Wijers (1998), Rhoads (1997, 1999), and others. On the other hand, since the observed flux arises from an integration over angles which samples gas at different source times, the steepening should set in gradually (Panaitescu & Mészáros 1998), so that it would not be surprising to have initially an intermediate slope, which takes a decade or more in time to steepen to a final asymptotic value. With these caveats in mind, and based on the data mentioned here, the agreement between the observed change in the decay slope and that expected from seeing the edge could well be a signature of the detection of a jet.

6 DISCUSSION

The conclusions that may be drawn from the above are that bright optical flashes, starting at the time of the gamma-ray trigger and extending into the early afterglow, are a straightforward prediction of the simplest afterglow models, as discussed by Mészáros & Rees (1997) and Sari & Piran (1999b,c). A 9th-magnitude optical flash from a gamma-ray burst at redshift $z \sim 1$ can arise from a reverse external shock with total energy $E_{53}(\Omega/4\pi)^{-1} \sim 1$ (Mészáros & Rees 1993, 1997), or from internal shocks with $E_{53}(\Omega/10^{-2}4\pi)^{-1}(t_\gamma/10^2\mathrm{s})^{-1} \sim 1$ (Mészáros & Rees 1997), where Ω is the solid angle into which the radiation is collimated, and t_γ is the gamma-ray burst duration. It is worth noting, however, that, even for a high gamma-ray peak flux or fluence, the magnitude of the optical prompt flash would depend on the magnetic field strength, which in a reverse shock case could be weaker than the field responsible for the gamma-rays (e.g. in internal or forward shocks). Furthermore, as known from $\log N$–$\log P$ fits (e.g. Krumholz, Thorsett & Harrison 1999), in essentially all cosmological models the luminosity function must be broad, with a range of E that can span upwards of 2–3 orders of magnitude. This is compatible with the fact that GRB 990123 is in the top 1 per cent of the BATSE brightness distribution (Galama et al. 1999), and also with the non-detection of similarly bright optical flashes from other bursts (e.g. with the LOTIS instrument: Park et al. 1999).

A steeper time decay law is expected for the early optical flash from reverse or internal shocks, compared with that expected from the forward shock. The decays calculated by Mészáros & Rees (1997) are too steep, compared with the observed t^{-2} behaviour (Akerlof et al. 1999), indicating the need to investigate different assumptions for the magnetic field behaviour. In the present paper we find that an early time decay $\propto t^{-2}$ is naturally explained in the standard (single Γ) afterglow model either by reverse shocks, where the field is in pressure equilibrium with the forward shock or is frozen-in, or by internal shocks in a magnetically dominated outflow. However, such decays may also be easily achieved in more realistic afterglow models without severe restrictions on the electron index, e.g. for bursts in an inhomogeneous external medium, or characterized by a power-law range of Γ.

From the fact that the inferred isotropic equivalent gamma-ray energy is 4×10^{54} erg (Galama et al. 1999; Kulkarni et al. 1999), one infers the need for a collimation of the gamma-rays by at least 10^{-1}, if not 10^{-2}, for any stellar-mass source. For an external reverse shock origin, from the energy alone the optical radiation need not be collimated at all; however, since in the observer frame the Lorentz factor of the reverse shock is initially close to that of the blast wave, the initial optical and gamma-ray collimations should be approximately the same (except, as may be the case, if the jet has an anisotropic Γ). For an internal shock origin of the optical flash, there are stronger grounds for expecting the beaming of the optical to be the same initially as for the gamma-rays.

The flattening of the light curve after 0.16 d to $\propto t^{-1.1}$ is well explained by the optical light expected from the forward shock, conforming to the standard interpretation of the long-term behaviour of afterglows. The discrepancy highlighted by Kulkarni et al. (1999) between the standard model prediction and the observed relation between the time decay index and the spectral slope is, as pointed out in Section 4, likely to be an indication of departures from the simple standard model, which could include, for example, an inhomogeneous external medium, non-uniform injection, or anisotropies. The renewed steepening of the optical light curve after about 3 d to $\propto t^{-1.65}$–$t^{-1.8}$ (Kulkarni et al. 1999; Fruchter et al. 1999) is, as argued in Section 5, likely to be due to seeing the edge of a jet, which gives a better fit for the magnitude of the change of slope than expected from sideways expansion, and also should occur well before the latter begins.

There are two main arguments that militate in favour of the optical flash observed in GRB 990123 being due to a reverse external shock. One is that the optical light curve does not show a good correlation with the gamma-ray light curve (Sari & Piran 1999c), and the model fits generally well. The other is that the flash from internal shocks is weaker than from external reverse shocks, as discussed here. On the other hand, the first ROTSE observations of GRB 990123 (Akerlof & McKay 1999) started 20 s after the gamma-ray burst trigger, and do not appear to have sampled the optical light curve densely enough to establish the degree of correlation with good significance (both curves are compared in Fenimore et al. 1999). While the likelihood of detecting an external reverse shock optical flash appears to be higher at this stage, one needs a faster triggering and a better sampling of the optical data in order to discriminate between reverse and internal shocks. An interesting possibility is that, since the external reverse shock starts somewhat later than the internal shocks, one might initially see a weaker optical flash from the internal shocks, which are overtaken by a stronger reverse external shock radiation after tens of seconds, until the forward shock optical afterglow takes over after 300–1000 s. Especially for the more frequent weaker bursts, the prospect of investigating such features underlines the need for dedicated gamma-ray burst afterglow space missions such as *HETE2* (Ricker et al. 1999) and *Swift* (Gehrels et al. 1999).

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