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VIEWING ANGLE AND ENVIRONMENT EFFECTS IN GAMMA-RAY BURSTS: SOURCES OF AFTERGLOW DIVERSITY

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ABSTRACT

We discuss the afterglows from the evolution of both spherical and anisotropic fireballs decelerating in an inhomogeneous external medium. We consider both the radiative and adiabatic evolution regimes and analyze the physical conditions under which these regimes can be used. Afterglows may be expected to differ widely among themselves, depending on the angular anisotropy of the fireball and the properties of the environment. They may be entirely absent or may be detected without a corresponding γ-ray event. A tabulation of different representative light curves is presented, covering a wide range of behaviors that resemble what is currently observed in the γ-ray bursts (GRBs) GRB 970228, GRB 970508, and other objects.

Subject headings: gamma rays: bursts — radiation mechanisms: nonthermal

1. INTRODUCTION

The discovery of the afterglows of γ-ray bursts (GRBs) provides new information that can constrain the models used to explain these objects. Significant interest was aroused by the fact that several of the features reported in the first GRB detected over timescales of greater than days at X-ray (X) and optical (O) wavelengths, i.e., GRB 970228 (Costa et al. 1997), agreed quite well with theoretical expectations from the simplest relativistic fireball afterglow models published in advance of the observations (Mészáros & Rees 1997a; see also Vietri 1997a). A number of theoretical papers were stimulated by this and subsequent observations (see, e.g., Tavani 1997; Waxman 1997a; Reichart 1997; Wijers, Rees, & Mészáros 1997 among others), and interest continued to grow as new observations provided controversial evidence for the distance scale and the possible host (Sahu et al. 1997). New evidence and new puzzles were added when the optical counterpart to the second discovered afterglow (GRB 970508) was claimed to have a cosmological redshift (Metzger et al. 1997), as well as a radio counterpart (Frail et al. 1997; Taylor et al. 1997) and X/O light curves showing a rise and decline (Djorgovski et al. 1997; Fruchter et al. 1997). Some bursts, however, were detected only in X but not O (e.g., GRB 970828), while some that would have been expected to be seen in X or O were not (e.g., GRB 970111).

Additional structure on the light curves has also emerged from a continued analysis of some of these objects down to the faintest flux levels. The large variety of behaviors exhibited by afterglows, while clearly compatible with relativistic fireball models, poses new challenges of interpretation (e.g., Waxman 1997b; Vietri 1997b; Katz & Piran 1997; Rhoads 1997; Paczyński 1997). Some of the questions at the forefront of attention include the effect of the external medium, the degree to which afterglows may be considered to be isotropic events, and the effects of the radiative efficiency on the evolution of the remnant. We address all three of these issues here. We also clarify some of the issues that have been recently raised about the dynamical effects of different radiative efficiency regimes. We then discuss the possible variety of afterglow behavior that is expected from isotropic or anisotropic fireballs expanding in a medium that may be inhomogeneous, either because of external gradients or because of expansion in an irregular cavity. We apply these models to interpret some of the salient observational features of several GRB afterglows and discuss their possible use for predicting detection rates of X/O/radio (R) afterglows undetected in γ-rays, as well as some possible reasons for the nondetection of afterglows in GRBs.

2. EXPANSION DYNAMICS AND RADIATIVE EFFICIENCY

In some bursts (e.g., GRB 970508) the afterglow seems to contain a significant amount of energy compared with the typical (isotropic) estimate of $E \sim 10^{51}$ ergs s^{-1}. This led Vietri (1997b) to suggest that the afterglow must remain radiatively efficient for approximately weeks after the burst and must evolve with $\Gamma \propto r^{-3}$. This regime was also considered by Katz & Piran (1997), who refer to previous relativistic fireball models as radiating only a small fraction of the total kinetic energy and go on to consider instantaneously cooling fireballs. It is important to discuss in more detail what conditions are necessary for the radiative efficiency having an effect on the dynamic evolution of an expanding cloud.

The classical fireball models, with "isotropic equivalent" energies of $E \sim 10^{51}$ ergs and bulk Lorentz factors $\Gamma \sim 10^2$–10^3 have, in fact, generally been taken to be in the radiative stage during the γ-ray event, i.e., radiative efficiency near unity, meaning that energy on the order of the initial total kinetic energy of the protons is radiated in the observer-frame expansion time (Rees & Mészáros 1992 and subsequent papers). However, for some parameters the bulk of this energy can appear at energies other than MeV (Mészáros, Rees, & Papathanassiou 1994; Mészáros & Rees 1994). This high efficiency in the initial deceleration shock can occur if the electrons are assumed *in the shock* to be heated to $\gamma_e \sim \xi_e(m_p/m_e)\Gamma$, with ξ_e being reasonably close to unity. Experimental evidence from interplanetary collisionless shocks indicates that this could be the case. In such fireballs, the electrons are likely to retain high radiative efficiency for some time after the GRB, and the most important energetically are the newly shocked electrons near the

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downward evolving peak of initial postshock energy $\gamma_e \sim \xi_e(m_p/m_e)\Gamma(t)$. This peak is where most of the electron internal energy generally is. In the regime where the peak electrons have high radiative efficiency, *if throughout the entire remnant volume the protons are able to establish (and remain in) equipartition with the electrons*, then the remnant evolves with $\Gamma \propto r^{-3}$ in a homogeneous external medium (e.g., Blandford & McKee 1976; Vietri 1997a; Katz & Piran 1997). This follows simply from the momentum conservation law, if we can assume that the radiative losses also tap the proton and magnetic energy (strong coupling) and that the radiative timescale is shorter than the dynamic timescale so that energy conservation cannot be used. This is the classical "snowplow" approximation of supernova remnants. The alternative regime is the one where the radiative losses do not tap the proton and magnetic energy, only the electron energy (weak coupling), and/or the radiative cooling timescale is longer than the dynamic timescale. In this case, possibly after an initial short cooling of the electrons, one can assume energy conservation (most of the energy is in the protons and/or magnetic fields), and one has $\Gamma \propto r^{-3/2}$ in a homogeneous medium (e.g., Blandford & McKee 1976; Paczyński & Rhoads 1993; Katz 1994; Mészáros & Rees 1997a).

The (strong-coupling) radiative regime and the adiabatic regime can be generalized to the case where the fireball moves into an inhomogeneous external medium. We consider a spherical fireball of energy E and bulk Lorentz factor Γ, which are independent of angle θ, expanding into an external medium of density $n(r) \propto r^{-d}$, which is also independent of the angle, where r is the distance from the center of the explosion. The shocked gas evolves according to the conservation law

$$\Gamma^{1+A} r^3 n \propto \Gamma^{1+A} r^{3-d} \propto \text{constant} , \quad (1)$$

where $A = 1(0)$ corresponds to the energy (momentum) conservation, i.e., to the adiabatic (radiative) regimes. These regimes must be understood in a global or dynamic sense as applying to the entire remnant, i.e., baryons, magnetic fields, electrons, etc., or at any rate to its dynamically dominant constituents. Since the observer-frame or detector time t must satisfy $r \propto ct\Gamma^2$, we have

$$\Gamma \propto r^{-(3-d)/(1+A)} \propto t^{-(3-d)/(7+A-2d)} , \quad r \propto t^{(1+A)/(7+A-2d)} . \quad (2)$$

If $d = 0$ and a remnant starts out in the strong-coupling radiative regime ($A = 0$), then $\Gamma \propto r^{-3}$. However, after the expansion has proceeded for some time, eventually the cooling time of the electrons at the peak of the distribution becomes longer than the expansion time, at which point energy conservation (adiabatic approximation) becomes valid (eq. [1] with $A = 1$) leading to $\Gamma \propto r^{-3/2}$.

The power law of equation (2) is valid as long as the remnant is relativistic, i.e., until the time $t_{\rm nr}/t_o \sim \Gamma_o^{(7+A-2d)/(3-d)}$, where $t_o \sim t_\gamma$ is the duration of the GRB itself and Γ_o is the initial Lorentz during the burst. For a homogeneous medium $d = 0$ and $t_{\rm nr}/t_o \sim \Gamma_o^{(7+A)/3}$, e.g., for an adiabatic remnant $A = 1$ with $\Gamma_o = 10^3$ and $t_\gamma = 1$ s, one has $t_{\rm nr} \sim 1$ yr. The subsequent behavior in the nonrelativistic stage is described by Wijers et al. (1997). Shorter times for reaching the nonrelativistic stage are possible for a radiative remnant or for a radiative stage followed by an adiabatic one, while longer times can occur for longer t_γ or for expansion into a medium whose density decreases with distance ($0 \leq d \leq 3$).

However, the conditions under which the remnant dynamics may be considered adiabatic or radiative is far from unambiguous and is crucially dependent on poorly known questions about postshock energy exchange between protons and electrons. Electrons cool quickly compared with protons, and it remains an unsolved question, of importance also in other areas of astrophysics, whether *behind the shocks*, after the electrons have cooled, the protons remain hot (i.e., a two-temperature plasma, as in hot torus models of active galactic nuclei [AGNs]) or whether they tend toward some degree of equipartition with the cooled electrons by virtue of unknown fast energy exchange mechanisms. The collisionless shock transition is the only place where one is guaranteed fast varying chaotic electric and magnetic fields that can lead to quick relaxation, and it is known that an interplay between protons and magnetic fields can occur there, which can lead to values close to equipartition. Such equipartition is also inferred from measurements in the ISM. However, in fast-evolving flows such as those in GRBs, it is unclear whether such equipartition can occur anywhere except possibly near the shock transition itself.

If behind the shocks the protons are unable to quickly readjust to the electron losses, then even if the electrons are radiative ($a = 1$) the protons can be adiabatic ($A = 1$) leading to $\Gamma \propto r^{-3/2}$ in a homogeneous external medium. Of course, if the electrons continue to reach strict equipartition in the shock, then some of the proton energy is continually fed to electrons, and this can depress the afterglow after 1 week by a small factor relative to the ideal adiabatic case (Sari 1997). A related problem is that in order for the remnant to evolve with $\Gamma \propto r^{-3}$ (again for a homogeneous external medium), the magnetic energy should also decrease, since if the latter were conserved it would soon dominate the total energy density and the remnant would evolve as a polytrope with adiabatic index 4/3, which leads to $\Gamma \propto r^{-3/2}$. Thus both the proton and the magnetic energy need to be transferred on a fast timescale to the electrons in order to ensure $\Gamma \propto r^{-3}$.

An additional factor that might contribute to a steepening of the decay of Γ are other energy losses, e.g., such as energy loss from the escape of accelerated nonthermal protons from the shell, if these carry substantial energy. So far, there are neither detailed simulations nor experimental evidence concerning this in GRBs. In the absence of such losses, or in the absence of a quick energy exchange between postshock protons plus magnetic fields and the electrons, one can have a situation where the electrons are "radiatively efficient" but the dynamics of the remnant expansion follows an "adiabatic" law. This occurs if the electron cooling time is less than the expansion time. The shocked electrons can radiate up to half of the proton energy in the shocks, but the protons and the magnetic fields retain at least half, and this would be enough to ensure a quasi-adiabatic dynamic evolution of the remnant with $\Gamma \propto r^{-3/2}$. The latter is also true when the electrons are adiabatic. For an inhomogeneous medium that decays with radius, both of these decay laws would be flatter.

There is no difficulty in treating the weak-coupling case, where the electrons are radiatively efficient but the dynamics is adiabatic (§ 3). This regime is physically as plausible, if not more, as the strong-coupling regime, where protons and fields exchange energy with electrons on a fast timescale. For reasons of simplicity, in §§ 3 and 4 and in the rest of the paper we will assume that magnetic fields are near equipartition with the protons, which ensures a simple expression for the electron radiative efficiency in terms of only the synchrotron cooling time and the expansion time (the situation where inverse Compton [IC] cooling is important only introduces some extra changes in the way the synchrotron efficiency is defined).

3. SPHERICAL INHOMOGENEOUS MODELS

As in the previous section, we consider a spherical fireball of energy E and bulk Lorentz factor Γ, independent of angle θ, expanding into an external medium of density $n(r) \propto r^{-d}$. In the simplest afterglow model one considers the time evolution of the radiation from the external medium shocked by the blast wave as it slows down. Denoting quantities in the comoving frame of the shocked fluid with primes, in the postshock region the density is $n' \sim n\Gamma$, the mean proton and electron random Lorentz factors are $\gamma_p \sim \Gamma$ and $\gamma_e \sim \xi_e(m_p/m_e)\Gamma$ [where $(m_e/m_p) \lesssim \xi_e \lesssim 1$ is the fraction of the electron equipartition energy relative to protons], the turbulently generated magnetic field (assumed to build up to a fraction $\xi_B \leq 1$ of the field in equipartition with the random proton energy) is $B' \propto \xi_B n^{1/2}\Gamma$, and the peak of the electron synchrotron spectrum is at comoving frequency $\nu'_m \propto B'\gamma_e^2 \propto \xi_B \xi_e^2 n^{1/2}\Gamma^3 \propto r^{-d/2}\Gamma^3$. The corresponding observer peak frequency is

$$\nu_m \propto \xi_B \xi_e^2 n^{1/2}\Gamma^4 \propto t^{-[12-(d/2)(7-A)]/(7+A-2d)} . \quad (3)$$

The synchrotron radiative efficiency at ν_m is $e_{\mathrm{sy},m} \sim t'^{-1}_{\mathrm{sy},m}/(t'^{-1}_{\mathrm{sy},m} + t'^{-1}_{\mathrm{ex}} + t'^{-1}_{\mathrm{other}})$, where $t'_{\mathrm{sy},m} \propto 1/(\xi_e \xi_B^2 \Gamma^3 n)$ is the comoving synchrotron cooling time at ν_m, t'_{ex} is the comoving expansion time or adiabatic cooling time, and t'_{other} is any other loss mechanism, e.g., IC, if important. In the limit where only synchrotron and/or adiabatic losses are important we may write $e_{\mathrm{sy},m} \sim (t'_{\mathrm{ex}}/t'_{\mathrm{sy},m})^a$, which is unity in the electron radiative ($a = 0$) regime and ≤ 1 in the electron adiabatic ($a = 1$) regime. We have $e_{\mathrm{sy},m} \sim (\xi_B^2 \xi_e r n \Gamma^2)^a \propto (\xi_B^2 \xi_e r^{1-d}\Gamma^2)^a$.

First, we assume (§ 2) that protons and magnetic fields are strongly coupled to electrons, so if the electrons are radiative the entire remnant is radiative and the index $A = a$. The comoving synchrotron intensity at the comoving peak frequency is $I'_{\nu'_m} \propto n'_e(P'_{\mathrm{sy}}/\nu'_m)ct'_{\mathrm{min}} \sim n'_e(\gamma_e m_e c^2/\nu'_m)ce_{\mathrm{sy},m} \propto \xi_B^{2a-1}\xi_e^{a-1} r^{a-d(a+1/2)}\Gamma^{2a-1}$, where $t'_{\mathrm{min}} \sim t'_{\mathrm{sy}}\, e_{\mathrm{sy},m}$ is the shortest of the possible cooling times (synchrotron or adiabatic in the above approximation). The flux from the relativistically expanding source at observer frequency ν_m is $F_{\nu_m} \propto t^2\Gamma^5 I'_{\nu_m}$, or

$$F_{\nu_m} \propto \nu_m^{-[2(1-a)-(d/2)(1+a)]/[12-(d/2)(7-a)]} \propto t^{[2(1-a)-(d/2)(1+a)]/(7+a-2d)} , \quad (4)$$

scaling with $\xi_e^{a-1}\xi_B^{2a-1}$. If the expansion is in the radiative $a = 0$ regime, $F_{\nu_m} \propto \nu_m^{-[2-(d/2)]/[12-(7/2)d]} \propto t^{(4-d)/(14-4d)}$ increases in time for any $d < 3$, and $F_{\nu_m} \propto \nu_m^{-1/6} \propto t^{2/7}$ in a homogeneous medium with $d = 0$ (Vietri 1997b obtains a different scaling by taking the shocked gas comoving width $\Delta R'$ in I'_{ν_m} as path length, which, however, is equal to ct'_{min} only for adiabatic expansion; e.g., Mészáros & Rees 1997a, where $t'_{\mathrm{min}} \sim t'_{\mathrm{ex}}$). In the adiabatic $a = 1$ regime, $F_{\nu_m} \propto \nu_m^{d/(12-3d)} \propto t^{-d/(8-2d)}$, which is a constant independent of ν_m and of time for adiabatic expansion in a homogeneous $d = 0$ medium (Mészáros & Rees 1997a; Katz 1994), but F_{ν_m} decreases in time for an inhomogeneous medium with $0 < d < 3$ (for $d \geq 3$ the fireball encounters most of the external mass near its initial radius). For a power-law spectrum $F_\nu \propto \nu^\alpha$, the flux at a fixed detector frequency ν_D is

$$F_D = F_{\nu_m}(\nu_D/\nu_m)^\alpha \propto t^{\{2(1-a)+12\alpha-(d/2)[(1+a)+\alpha(7-a)]\}/(7+a-2d)} , \quad (5)$$

which scales as $\xi_e^{-1+a-2\alpha}\xi_B^{-(1+\alpha)+2a}$. Equation (5) is valid for the strong electron-proton coupling regime. For a typical synchrotron spectrum $\alpha \simeq \frac{1}{3}\,(-1)$ below (above) the break frequency $\nu_m(t)$, as the latter decreases in time, the flux from radiative $a = 0$ models in the detector frequency band ν_D at times for which $\nu_D < \nu_m$, $\nu_D > \nu_m$ is $F_D \propto t^{6/7}$, $t^{-10/7}$ in a homogeneous $d = 0$ medium, and, for $\nu_D < \nu_m$, $\nu_D > \nu_m$ the flux is $F_D \propto t^{8/9}$, $t^{4/3}$ in an inhomogeneous $d = 2$ medium. Adiabatic $a = 1$ models give $F_D \propto t^{1/2}$, $t^{-3/2}$ in a homogeneous medium and $F_D \propto t^0$, t^{-2} in an inhomogeneous $d = 2$ medium. Other values can be calculated from equation (5) for different α before and after the break. Tables 1 and 2 give several examples in the next to last column for

TABLE 1

TIME EXPONENTS OF SYNCHROTRON PEAK FLUX AND PEAK FREQUENCY

a	d	F_{ν_m}, ν_m	$j = 3k$	$j = 2k$	$j = k$	$j = 0, k = 1$	$j = -k$	$j = -2k$	$j = -3k$	Isotr $A = a$	Isotr $A = 1$
						F_{ν_m}					
0......	0	$F_{\nu_m} \propto \nu_m^q$	1/6	0	−1/6	−1/3	−1/2	−2/3	−5/6	−1/6	−1/3
0......	2	$F_{\nu_m} \propto \nu_m^q$	1/3	0	−1/5	−1/3	−3/7	−1/2	−5/9	−1/5	−1/3
1......	0	$F_{\nu_m} \propto \nu_m^q$	3/4	1/2	1/4	0	−1/4	−1/2	−3/4	0	0
1......	2	$F_{\nu_m} \propto \nu_m^q$	2/3	1/2	2/5	1/3	2/7	1/4	2/9	1/3	1/3
						ν_m					
0......	0	$\nu_m \propto t^p$	−12/5	−2	−12/7	−3/2	−4/3	−6/5	−12/11	−12/7	−3/2
0......	2	$\nu_m \propto t^p$	−3	−2	−5/3	−3/2	−7/5	−4/3	−9/7	−5/3	−3/2
1......	0	$\nu_m \propto t^p$	−12/5	−2	−12/7	−3/2	−4/3	−6/5	−12/11	−3/2	−3/2
1......	2	$\nu_m \propto t^p$	−3	−2	−5/3	−3/2	−7/5	−4/3	−9/7	−3/2	−3/2

NOTE.—Exponents of the power-law dependence of the synchrotron peak flux F_{ν_m} as a function of the time-varying synchrotron peak ν_m (*top four rows*) and exponents of the time dependence of ν_m on observer (detector) time t (*bottom four rows*). The first column indicates whether the electrons are in the radiative ($a = 0$) or adiabatic ($a = 1$) regime, and the second column indicates the value of the exponent of the external medium density dependence on radius $n \propto r^{-d}$, $d = 0$ being homogeneous. The fourth through tenth columns give the exponents for the anisotropic model $E \propto \theta^{-j}$, $\Gamma \propto \theta^{-k}$ of § 4 and various values of j and k. The last two columns on the right give the corresponding exponents for the isotropic models of § 3, the first being for the strong electron-proton coupling $A = a$ case and the second for the weak-coupling case $A = 1$. For $A = 0(1)$ the remnant as a whole is dynamically radiative (adiabatic).

TABLE 2

TIME EXPONENTS OF FLUX OBSERVED IN A GIVEN BAND

			$F_D \propto t^w$								
a	d	α	$j = 3k$	$j = 2k$	$j = k$	$j = 0, k = 1$	$j = -k$	$j = -2k$	$j = -3k$	Isotr $A = a$	Isotr $A = 1$
0......	0	1/3	2/5	2/3	6/7	1	10/9	6/5	14/11	6/7	1
0......	0	−1	−14/5	−2	−10/7	−1	−2/3	−2/5	−2/11	−10/7	−1
0......	2	1/3	0	2/3	8/9	1	16/15	10/9	8/7	8/9	1
0......	2	−1	−4	−2	−4/3	−1	−4/5	−2/3	−4/7	−4/3	−1
1......	0	1/3	−1	−1/3	1/7	1/2	7/9	1	13/11	1/2	1/2
1......	0	−1	−21/5	−3	−15/7	−3/2	−1	−3/5	−3/11	−3/2	−3/2
1......	2	1/3	−1	−1/3	−1/9	0	1/15	1/9	1/7	0	0
1......	2	−1	−5	−3	−7/3	−2	−9/5	−5/3	−11/7	−2	−2

NOTE.—Exponents of the power-law dependence of the synchrotron flux in a given detector band ν_D as a function of observer (detector) time t. The first column indicates whether the electrons are radiative ($a = 0$) or adiabatic ($a = 1$); the second column gives the exponent of the external medium density dependence on radius $n \propto r^{-d}$, $d = 0$ being homogeneous; the third column gives the value of the spectral index α of the spectrum, $F_\nu \propto \nu^\alpha$, which may be $\frac{1}{3}$ below the peak ν_m, and -1 above (although these indices can vary around these representative values). As the peak ν_m passes down through the observation band ν_D, the index α changes from the positive value left of the peak to the negative to the right of the peak. The fourth through tenth columns give the time exponents of the $F_D \propto t^w$ in the detector band for the anisotropic model $E \propto \theta^{-j}$, $\Gamma \propto \theta^{-k}$ of § 4 and various values of j and k. The last two columns on the right give the corresponding exponents for the isotropic models of § 3, the first being for the strong-coupling case $A = a$ and the second for the weak coupling $A = 1$. For $A = 0(1)$ the remnant as a whole is dynamically radiative (adiabatic).

the isotropic strong-coupling case.

A different regime is obtained if one assumes that the protons and magnetic fields are *not* strongly coupled to the electrons behind the shocks. In this case, the dynamics of the remnant as a whole is controlled by the index A in equations (1) and (2), and as long as $\Gamma \gg 1$, the remnant is adiabatic with $A = 1$ and $\Gamma \propto r^{-3/2}$ for a homogeneous medium, whether the electrons are radiative or not, i.e., independent of a or not (see § 2). More generally, in this case $\Gamma \propto t^{-(3-d)/(8-2d)}$, $r \propto t^{2/(8-2d)}$, and

$$\nu_m \propto \xi_B \xi_e^2 t^{-(12-3d)/(8-2d)} . \quad (6)$$

In the synchrotron efficiency one must keep a separate index $a \neq A$ to account for the electrons being radiative or not. We then have $I'_{\nu_m} \propto \xi_B^{2a-1} \xi_e^{a-1} r^{a(1-d)-d/2} \Gamma^{2a-1}$, and $F_{\nu_m} \sim t^2 \Gamma^5 I'_{\nu_m}$ is

$$F_{\nu_m} \propto \nu_m^{-[4(1-a)-d]/(12-3d)}$$
$$\propto \xi_B^{2a-1} \xi_e^{a-1} t^{[4(1-a)-d]/(8-2d)} . \quad (7)$$

In a given fixed detector band ν_D, one observes for a typical spectral shape $F_\nu \propto \nu^\alpha$ a time-dependent flux $F_D \propto F_{\nu_m} \nu_m^{-\alpha}$ or

$$F_D \propto \xi_B^{(2a-1-\alpha)} \xi_e^{a-1-2\alpha} t^{[4(1-a)-d+\alpha(12-3d)]/(8-2d)} . \quad (8)$$

This is valid in the relativistic expansion regime. After a remnant becomes nonrelativistic (see eq. [2] above), the flux in a homogeneous $d = 0$, $a = 1$ medium would steepen to $F_D \propto t^{(3+15\alpha)/5} \propto t^{-12/5}$ for $\alpha = -1$ (Wijers et al. 1997). The variety of time behaviors possible for relativistic isotropic models in both the strong and weak-coupling cases, depending on whether they are radiative or adiabatic, is shown in the last two columns of Tables 1 and 2.

The time behavior given by equations (5) and (8) can be complicated by at least two effects. One is that unless observations start after the peak electrons are adiabatic, at some subsequent time the value of a in equations (4), (5), (7), and (8) switches from 0 to 1 as the peak electrons become adiabatic. The second is that if the detector frequency ν_D is greater than ν_m, the flow (being controlled by particles radiating at ν_m) can already be adiabatic, while the smaller number of higher energy particles radiating at ν_D may still radiate efficiently. The frequency and time power-law dependences of the flux before and after electron radiative inefficiency occurs are different. For a steady injection of accelerated electrons, the self-consistent electron energy power-law index p in the presence of fast synchrotron losses is one power steeper than in the injected spectrum, and the self-consistent synchrotron power-law index $\alpha = (p - 1)/2$ is a half power steeper than for the adiabatic (negligible loss) case. If the lowest energy electrons near the peak $\gamma_e \sim \xi_e(m_p/m_e)\Gamma$ are radiatively efficient, all electrons above that are as well. As the remnant evolves, the first electrons to become inefficient are the lowest energy ones [in the peak corresponding to $\nu_m(t)$], and a flattening break by half power in the photon spectrum at frequency $\nu_b(t) > \nu_m(t)$ moves to frequencies increasingly higher than $\nu_m(t)$. The electron Lorentz factor at which the synchrotron time just equals the expansion time $r/c\Gamma$ is $\gamma_b \propto r^{[2-A(1-d)]/(1+A)}$, and the corresponding "adiabatic" photon frequency is

$$\nu_b \propto B' \gamma_b^2 \Gamma \propto n^{1/2} \Gamma^2 \gamma_b^2$$
$$\propto r^{-[2-(3/2)d]} \propto t^{-[2-(3/2)d](1+A)/(7+A-d)} , \quad (9)$$

which can either decrease in time for $d < 4/3$ (including a homogeneous medium with $d = 0$) or increase for $d > 4/3$ [although it always increases with respect to $\nu_m(t)$]. For an external medium whose density drops with radius faster than $d \gtrsim 4/3$, if initially $\nu_b > \nu_D$, it will always remain so, and the spectral index remains adiabatic without change (until a much higher cutoff is reached, at which the acceleration becomes inefficient and the spectrum drops off exponentially). However, for a homogeneous medium or one with $d < 4/3$, if initially $\nu_b > \nu_D$, the photon spectral index will at some later time steepen by $\frac{1}{2}$ as ν_b sweeps through the observing band ν_D and the observed spectrum transitions from the adiabatic to the radiative regime.

4. ANISOTROPIC INHOMOGENEOUS MODELS

The observed afterglow temporal decays are conventionally fitted by power laws, and it is interesting to explore how the decay slopes would depend on the angular dependence of the dynamically relevant quantities of a fireball. To that effect, we consider anisotropic relativistic outflows where both the energy per unit solid angle and the bulk Lorentz factor depend on the angle θ as power laws (at least over some range of angles), and we also consider the external density distribution to depend on radius as a power law; i.e.,

$$E \propto \theta^{-j}\,, \quad \Gamma \propto \theta^{-k}\,, \quad n \propto r^{-d}\,. \tag{10}$$

If there is a well-defined jet, the normalizations of E and Γ may be different for material inside and outside the jet opening angle θ_o. At each angle the outflow starts converting a significant fraction of its bulk kinetic energy into radiation when an external blast wave develops at the angle-dependent deceleration radius $r_d \propto [E/n(r_d)\Gamma^2]^{1/3}$ at an angle-dependent observer-frame (detector) time $t \sim r/c\Gamma^2$. The θ-dependence of equation (10) implies that the deceleration blast wave at different angles occurs at

$$\Gamma \propto t^{-k(3-d)/(8k-j-2dk)}\,, \quad r \propto t^{(2k-j)/(8k-j-2dk)}\,,$$
$$E \propto t^{-j(3-d)/(8k-j-2dk)}\,, \quad \theta \propto t^{(3-d)/(8k-j-2dk)}\,. \tag{11}$$

Depending on the normalization of equation (10) and causality considerations, the radiation from the blast waves occurring at increasing θ at successive times t can dominate the afterglow evolution (as opposed to the decay of E and Γ along the same θ as a function of time). For instance, if the event has been detected at γ-rays and there is a jet of opening angle θ_o, the observer is presumably within angles $\lesssim \theta_o$ from the axis. For $\Gamma = \Gamma_o(\theta/\theta_o)^{-k}$, in order for subsequent blast waves at $r = r_d$ from $\theta > \theta_o$ to be observed at times $t > t_o$, one needs $\Gamma^{-1} \gtrsim \theta$ to be satisfied, that is $\theta/\theta_o > (\theta_o \Gamma_o)^{1/(k-1)}$, or

$$t/t_o \gtrsim (\theta_o \Gamma_o)^{(8k-j-2dk)/[(3-d)(k-1)]}\,. \tag{12}$$

For values of $k < 1$, the blast waves are detectable at all angles, but for $k > 1$ there are ranges of k, j for which $\theta_o \Gamma_o$ is limited to values $\lesssim 5$–10 in order to detect the blast wave at reasonable t/t_o. In this case, the initial part of the afterglow may be due to the evolution in time of the gas responsible for the burst initially observed, until such a time when the causality condition is satisfied for gas at larger angles and the newly shocked gas at increasing angles can become dominant in providing the observed flux. This second case introduces additional complexities and will not be discussed here, since even the simpler case first mentioned above will serve to illustrate the point that great variety can be expected in the temporal behavior of afterglows.

For the conditions where the afterglow is dominated by the newly shocked gas at increasing angles, the observer-frame peak frequency of the synchrotron radiation spectrum from the blast waves (eq. [11]) coming from increasingly larger θ at increasing times t is

$$\nu_m \propto \xi_B \xi_e^2 n^{1/2}\Gamma^4 \propto t^{-[12k-d(3k+j/2)]/(8k-j-2dk)}\,. \tag{13}$$

The observer-frame intensity at this peak frequency is

$$I_{\nu_m} = \Gamma^3 I'_{\nu'_m} \sim \Gamma^2(E e_{\mathrm{sy},m}/4\pi r^2 t \nu_m) \propto \xi_B^{-1}\,\xi_e^{-2}\,e_{\mathrm{sy},m}$$
$$\times\, t^{-d(k-j/2)/(8k-j-2dk)}\,,$$

where, as in § 3, the synchrotron efficiency is $e_{\mathrm{sy},m} \sim (t'_{\mathrm{ex}}/t'_{\mathrm{sy},m})^a \propto (\xi_B^2\,\xi_e\, r^{1-d}\Gamma^2)^a$ if synchrotron and adiabatic cooling are the two most important energy-loss mechanisms (or its generalization if IC or other effects need to be included). We then have $e_{\mathrm{sy},m} \propto t^{-a[4k+j(1-d)]/(8k-j-2dk)}$, where $a = 1(0)$ if the peak electrons in the deceleration blast wave at the angle corresponding to detector time t are adiabatic (radiative). (In this model, the dominant radiation is produced at the initial deceleration blast wave for that θ, so a does not enter in the dynamics, only in the radiative efficiency of the initial blast wave.) The flux observed at ν_m from the deceleration blasts at increasing θ is $F_{\nu_m} \sim t^2\Gamma^2 I_{\nu_m}$, or

$$F_{\nu_m} \propto \nu_m^{-\{4k-2j-d(k-j/2)-a[4k+j(1-d)]\}/[12k-d(3k+j/2)]}$$
$$\propto t^{\{4k-2j-d(k-j/2)-a[4k+j(1-d)]\}/(8k-j-2dkW}\,, \tag{14}$$

which scales with $\xi_B^{2a-1}\xi_e^{a-1}$. Depending on the normalization of equation (10), the flux equation (14) from increasing θ values can dominate the flux given by equations (4) or (7). (In other cases, one can approximate the evolution as the superposition of isotropic blast waves from individual θ, which could in some cases be dominated by that of the central jet region.) At a fixed detector frequency ν_D, the observed flux corresponding to equation (14) is

$$F_D \sim F_{\nu_m}(\nu_D/\nu_m)^\alpha$$
$$\propto t^{\{k(4+12\alpha)-2j-d[k(1+3\alpha)-(j/2)(1-\alpha)]-a[4k+j(1-d)]\}/(8k-j-2dk)}\,, \tag{15}$$

which scales with $\xi_B^{2a-1-\alpha}\xi_e^{a-1-2\alpha}$. For characteristic spectra with α positive (negative) below (above) the break, this leads to detected fluxes that initially rise in time and then decay. However, a variety of behaviors are possible, including some where after the break passes through the detector window, the flux continues to grow at a slower rate or saturates. Note that the scaling of equation (14) with ν_m allows both for F_{ν_m} to decrease or to increase as ν_m decreases in time, both in the radiative and adiabatic cases, depending on the values of j and k, which characterize the angular dependence of E and Γ.

5. DISCUSSION

5.1. *Dynamics, Cooling, and Decay Law*

In the simplest model where the GRB and the afterglow both arise from an external shock (case a1 of Mészáros & Rees 1997), to zeroth order the GRB γ-ray flux should lie near the backward extrapolation of the afterglow, provided the basic conditions have not changed and the same radiation mechanisms are responsible for both. This is clearly a rough approximation, since it is likely that the GRB is initially radiatively efficient and becomes radiatively inefficient at some later stage. The observations of GRB 970228, especially following the *Hubble Space Telescope* (*HST*) observations of 1997 September (Fruchter et al. 1997), indicate an optical flux decaying as an approximately constant power law in time. There may be shorter timescale variations superposed on this overall long-term behavior (e.g., Galama et al. 1997), which could be either due to a difference in calibration or a real wiggle in the decay.

At late times (e.g., 6 months in GRB 970228), the remnant is most likely adiabatic. One question that can arise is whether a simple external shock afterglow model whose dynamics is manifestly "adiabatic" can be radiatively efficient enough to produce the initial relatively high X-ray and optical afterglow luminosity. For an afterglow luminosity, $L_{\mathrm{X,O}} \lesssim L_\gamma$; as observed from our discussion in § 2, this is not a problem. At each radius, the electrons can radiate up to half of the total newly shocked proton energy randomized in the shock transition, as long as the electron cooling time is shorter than the expansion time. The electrons can be radiatively efficient even when the dynamics of the remnant (i.e., the shell of hot protons and magnetic fields behind the shock that provide most of the mass and inertia) follows an "adiabatic" law $\Gamma \propto r^{-3/2}$ for a homogeneous medium. Arguments were presented in § 2 about why the latter behavior may be more likely than a faster evolution with $\Gamma \propto r^{-3}$, a conclusion that is also supported by comparison with observations relating to the size of the GRB 970228 remnant (Waxman, Kulkarni, & Frail 1997). Note also that, from equation (2), for expansion in a medium whose density decreases, Γ could drop even more shallowly than the above for a homogeneous medium. As shown by Tavani (1997), Waxman (1997a), Wijers et al. (1997), Reichart (1997), and others, in GRB 970228 the initial γ-ray flux and the overall X/O afterglow behavior are in good agreement with a simple external shock where the observations started after $\nu_m < \nu_D$, without any substantial changes of slope during the observed decay phase. However, as discussed below, there are various mechanisms capable of producing changes in the decay slope, which could be responsible for some of the reported departures from a simple power-law behavior.

5.2. *Intensity Offsets between Afterglows and Main Burst*

In the case of GRB 970228, a backward extrapolation of the afterglow flux shows some hints of undershooting the γ-ray flux. If this were real, a slight undershooting could be due to the GRB being radiative initially (this is expected especially during the initial deceleration shock), with the afterglow dynamics becoming adiabatic soon afterward. This would lead to a flattening of the spectral slope (e.g., second row, last two columns of Table 2).

Another possibility (Mészáros & Rees 1997a; Wijers et al. 1997; Katz & Piran 1997) is that the γ-rays could have originated in an internal shock. Internal shocks leave essentially no afterglow, yet they should be followed (eventually) by external shocks. An internal shock leaves unused anywhere from $\sim$20% to $\gtrsim$90% of the total kinetic energy (Rees & Mészáros 1994; Kobayashi, Piran, & Sari 1997). The leftover energy is liberated in the external shock, most of whose initial radiation can come out at GeV energies because initially IC losses dominate over synchrotron losses; the initial synchrotron MeV radiation would then be typically low below the BATSE threshold (Mészáros & Rees 1994), but later on it becomes dominant over IC (Waxman 1997b). Thus at later stages, the afterglow from the synchrotron peak can have a flux level whose back extrapolation might overshoot the γ-ray flux. Since external shocks are generally smoother (at most 3–5 pulses; Panaitescu & Mészáros 1998), while internal shocks may be very variable (Rees & Mészáros 1994; Kobayashi et al. 1997), a natural conclusion (also reached independently by Piran 1997) is that since afterglows appear to arise from external shocks, a burst where the γ-ray light curve (or at least the last γ-ray pulse of the light curve) is relatively smooth has a better chance of leaving behind a visible afterglow at lower frequencies.

A reason for offsets may also be (§ 4) that the relativistic outflow has an angle dependence in either the energy, the bulk Lorentz factor, or both. The GRB itself may be due, for instance, to a high-Γ ejecta that shocks first, while the afterglow could be dominated by slower material ejected at larger angles relative to the observer, which shock later and produce softer radiation, but which could carry a substantial or even larger fraction of the total energy. If detected this would generally be as an upward offset of the afterglow relative to the GRB, since otherwise the afterglow would be dominated by the evolution of the same material that gave rise to the GRB.

5.3. *Rise and Decay, Late Rises, Bumps*

An initial afterglow flux rise followed by a decay is a direct consequence of the simplest afterglow model (a1 of Mészáros & Rees 1997a) and is true also of any generic peaked spectrum from an expanding cloud where the peak energy decreases in time faster than the peak flux. Estimates for expanding clouds were made by Paczyński & Rhoads (1993) and Katz (1994) based on simplified radiation models. In the more detailed model (a1) of Mészáros & Rees (1997a), if the slope α at frequencies below ν_m is positive, i.e., $F_\nu \propto \nu^\alpha$, the initial rise is $F_D \propto t^{3\alpha/2}$ for $\nu_D < \nu_m$ in a homogeneous medium (Wijers et al. 1997). For an "average" GRB spectrum, $\alpha \sim 0$ below the break (Band et al. 1993), which implies that F_{ν_D} is initially constant (Mészáros & Rees 1997a). However, $\alpha \sim 0$ is only the average value; there are many GRBs with $\alpha > 0$ below the break (this is also the case for an ideal synchrotron spectrum $\alpha \sim \frac{1}{3}$ below the break, e.g., Mészáros et al. 1994), and in such cases one obtains an initial power-law increase in F_D. After the frequency ν_m of the peak has dropped below the observing frequency ν_D, if the spectrum above the peak is $F_\nu \propto \nu^\beta$ with $\beta < 0$, the flux observed at ν_D starts to decay $\propto t^{-3\beta/2}$. This occurs after a time $t_D/t_\gamma \sim (\nu_\gamma/\nu_D)^{2/3}$ (for an adiabatic remnant $\Gamma \propto t^{-3/8}$), which for observations in the R-band is $t_{\rm opt} \sim 10^4 t_\gamma$. A maximum at 1.5 days, such as in GRB 970508, is therefore compatible with the observed $t_\gamma \sim 5 \times 10^1$ s.

However, there are other mechanisms that can give rise to optical fluxes rising at times that could be even later than the above. One is an anisotropic outflow (§ 4), where the GRBs and X-rays come from material close to the axis oriented near the line of sight, and further off-axis material with a slower Γ starts to decelerate after 1.5 days, say, or has been slowed down enough that its light cone includes the line of sight after 1.5 days. Since beaming does not change the slope, if the spectrum is a pure power law then its slope would remain constant. While a constant spectral slope is the simplest assumption, it is not a necessary one. It is conceivable, for instance, that the slope of accelerated electrons (and consequently of the synchrotron spectrum) could change as the bulk Lorentz factor and the shock strength changes, or it might depend on other effects associated with the angular anisotropy of the outflow. An interesting result is the indication that the optical light curve of GRB 970508 may have been steady or even decreasing (Pedersen 1997) before the 1.5 day rise phase preceding the maximum. This may be simply explained by an anisotropic model such as the one described in § 4, where the indices j and k change to give a light-curve transition in this sense or, even more simply, by a bimodal model, where one has a central jet associated with the γ-ray event whose tail is just seen to decay, followed by the emission of a slower Γ outflow over much wider angles outside the jet, which is responsible for the main part of the afterglow. Another possibility for a late rise in the optical would be if a lower Γ shell catches up with the main shock front with a comparable energy after $t \sim 1.5$ days, so we see then the emission from the onset of deceleration of this late shell. Very prolonged optical decays with a shallow power law are possible in anisotropic models such as those in § 4 (e.g., Table 2, the sixth or seventh row).

An alternative explanation for a late turn-on of an afterglow may be that the GRB occurs inside a very low density cavity inflated by the pulsar activity of one of the neutron stars in the progenitor binary. The shock, at least over a range of directions, would not arise until the ejecta hits the wall of the cavity, and this could take a time on the order of weeks, the ensuing shock being spread over a dynamic timescale sufficiently short to produce a large flux per unit time. A characteristic feature of pulsar cavities is that they are usually asymmetric and irregular in shape, often being elongated because of the proper motion of the energizing source. One would naturally expect a wide variety of time histories for the afterglows arising from the impact of a (possibly anisotropic) ejecta on an irregularly shaped cavity wall whose dimension (depending on direction respect to the line of sight) may vary considerably.

5.4. *Peak Flux Level Evolution*

In some observed afterglows, the flux level F_ν at lower energies is, at least initially, significant relative to that of the maximum γ-ray flux (even if νF_ν is smaller). One prediction of the simplest model is that the maximum value of the afterglow flux in every band, F_{ν_m}, is a constant. An interesting case is that of GRB 970508, where the ratio of maximum $F_{\rm X}$ to maximum F_γ is $\sim\frac{1}{2}$, while the ratio of $F_{\rm O}$ to F_γ is $\sim 1/10$. Considering the drastic simplifications involved in the simplest model (homogeneous medium, constant equipartition fields, etc.), this order-of-magnitude agreement is perhaps encouraging. However, $F_{\nu_m} \sim$ constant is clearly an approximation that need not always hold for less simple models. In Table 1, the top four rows of the last two columns show that for isotropic but inhomogeneous outflows one can expect F_{ν_m} to either increase or decrease with ν_m as the later decreases, while the previous columns of the top four lines in the same table show that either of these behaviors may also arise as a result of an anisotropic outflow. Thus, the ratio of F_{ν_m} at γ-rays and lower wavelengths could be either larger or smaller than 1, simply on this basis. In a simple bimodal anisotropic model, one can simply have more or less energy in the large angle slower outflow seen later than in the early and harder axial outflow.

A related question is the observability of radio fluxes of order mJy around 10^{10} Hz, as reported for GRB 970508 (Frail et al. 1997). Radio fluxes of this magnitude arise naturally in a simple isotropic homogeneous fireball model after a week or so, since the self-absorption frequencies (overestimated by 10^2 by Mészáros & Rees 1997a) are in the range of 10^{11} Hz initially and drop to 10^{10} Hz in about a week and the flux level is well below the brightness temperature limits for incoherent synchrotron radiation. What is more interesting is the relatively large value of the R flux ($\sim$mJy) relative to the O flux (~ 50 μJy). In this afterglow, therefore, F_{ν_m} first decreases with ν_m between γ and O energies but then increases with decreasing ν_m between the O and the R bands. The simplest explanation may be in terms of a jet-wind two-component model. The decrease between γ and O could be due to expansion of a jet into an inhomogeneous medium (e.g., Table 1, fourth row), and the increase between O and R could be due to a surrounding low-Γ wind at larger angles, which shocks at later times, as suggested by Wijers et al. (1997). The slow wind need not have an angular dependence; a growth and decay of the flux is approximated also by behaviors represented in Table 2 in the last two columns for expansion into either a homogeneous or an inhomogeneous external medium. A wind with $\Gamma \sim 3$–10 can match the delayed emergence (~ 1 week) of the radio from the wind blast wave. The tables give only illustrative values for selected spectral and density exponents, which can easily be changed to fit a particular observed rate of growth and decay. Actually, the radio flux of GRB 970508 at 10^{10} Hz (Frail et al. 1997) at first increased and then appears to flatten, except for decaying oscillations, which could be due to scintillation (Goodman 1997; Waxman et al. 1997) followed by a slow decline.

5.5. *Burstless Afterglows and Afterglowless Bursts*

An interesting consequence of anisotropic models (Mészáros & Rees 1997b; Rhoads 1997) is that there could be a large fraction of detectable afterglows for which no γ-ray event is detected. If the observer lies off-axis to the jet, then the detected "afterglow" can be approximated by the isotropic model calculated for an $E(\theta_{\rm obs})$, $\Gamma(\theta_{\rm obs})$ corresponding to the offset angle $\theta_{\rm obs}$ of the observer to the jet axis. As Γ drops after the deceleration shock, the causal angle includes an increasing amount of the solid angle toward the jet as well as toward the equator, and depending on the values of j and k, the observed flux would generally decrease in time. However, for some choices of j, k an increase in the light curve might be possible, depending on the normalization. It could be that gravitational energy is converted more efficiently into kinetic energy of expansion at large angles, where the opacity is larger (cf. Paczyński 1997). After $\Gamma(\theta_{\rm obs})$ has dropped to the point where the central jet portion $\theta \lesssim \theta_o$ is detectable, the late stages of the jet emission would become visible at a later stage when it is bright only at wavelengths longer than γ-rays. If the jet contained substantially more energy than the off-axis regions so that it dominates the flux even after expanding for a longer time than the initially observed off-axis region, one would expect an additional increase or flattening of the light curve at this point. Details would be further complicated by contributions from the equator and the back side of an opposite jet, if $\theta_{\rm obs} \gtrsim \pi/4$. The statistics of afterglows not detected in γ-rays can be calculated from equations (11) and (14).

The converse question is why some bursts (e.g., GRB 970111) have been detected in γ-rays but not in X or O, even though it was in the field of view of *BeppoSAX*, which would have been expected to detect it if the X- to γ-ray ratio had been comparable to GRB 970228 (a weak X-ray afterglow may have been detected; Costa 1997). One reason may be that the γ-ray emission is due to internal shocks (which leave essentially no afterglows; Mészáros & Rees 1997a) and the environment has a very low density, in which case the external shock can occur at much larger radii and over a much longer timescale than in usual afterglows, and the X-ray intensity is below the threshold for triggering. This may be the case for GRBs arising from compact binaries that are ejected from the host galaxy into an external environment that is much less dense that the ISM assumed for usual models. Another possibility for an unusually low density environment, made up only of very high energy but extremely low density electrons, is if the GRB goes off inside a pulsar cavity inflated by one of the neutron stars in the precursor binary. Such cavities can be as large as fractions of a parsecond or more, giving rise to a deceleration shock months after the GRB with a consequently much lower brightness that could avoid triggering and detection.

The lack of an afterglow in some bursts may also be due to occurrence in an unusually high density environment (e.g., a star-forming region or the inner kiloparseconds of a late-type spiral, where failed supernova or hypernova progenitors may reside; e.g., Paczyński 1997). This could lead to a more rapid onset of the deceleration leading to the X-ray phase, and it would also imply an increased neutral gas column density and optical depth in front of the source. A special case is that of GRB 970828, where X-rays have been observed but no optical radiation down to faint levels (Groot et al. 1997). The presence of a significant column density of absorbing material has been inferred from the low-energy turnover of the X-ray spectrum (Murakami et al. 1997), and the corresponding dust absorption may, in fact, be sufficient to cause the absence of optical emission (Wijers & Paczyński 1997, private communication). The difference between the low-density and high-density environment cases could be tested if future observations of afterglows reveal a correlation with the degree of galaxy clustering or with individual galaxies.

In conclusion, the absence of detected afterglows in many bursts is not surprising, while there may be detected afterglows also in some cases where a corresponding $\gamma\gamma$-ray burst has not been detected; and when afterglows are detected, a wide diversity of behaviors may be the rule, rather than the exception.

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