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TWO POPULATIONS AND MODELS OF GRBs

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TWO POPULATIONS AND MODELS OF GAMMA-RAY BURSTS

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ABSTRACT

Gamma-ray burst statistics are best explained by a source population at cosmological distances, while spectroscopy and intensity histories of some individual bursts imply an origin on Galactic neutron stars. To resolve this inconsistency I suggest the presence of two populations, one at cosmological distances and the other Galactic. I build on ideas of Shemi & Piran (1990) and of Rees & Mészáros (1992) involving the interaction of fireball debris with surrounding clouds to explain the observed intensity histories in bursts at cosmological distances. The distances to the Galactic population are undetermined because they are too few to affect the statistics of intensity and direction; I explain them as resulting from magnetic reconnection in neutron star magnetospheres. An appendix describes the late evolution of the debris as a relativistic blast wave.

Subject headings: gamma rays; bursts — MHD — stars: neutron

1. INTRODUCTION

Attempts to explain all the observed gamma-ray bursts (GRBs) with a single population of sources have become progressively more difficult. On one hand, their distribution on the sky has been observed, with steadily improving precision (Atteia et al. 1987; Megan et al. 1992), to be isotropic, an observation which is naturally explained (Usov & Chibisov 1975; Goodman 1986; Paczyński 1986; Dar et al. 1992; Dermer 1992; Fenimore et al. 1992; Mao & Paczyński 1992a; Mészáros & Rees 1992; Paczyński 1992a; Piran 1992) if they are at cosmological distances. On the other hand, a number of GRBs have been reported to show (Higdon & Lingenfelter 1990) spectral features at a few tens of keV and at about 400 keV, which are readily interpretable as cyclotron lines and the two-photon positron annihilation line from the surface of magnetized neutron stars at Galactic distances, but which are inexplicable at cosmological distances. If their validity is accepted, the data appear irreconcilable.

The problem is further complicated by the fact that straightforward models of radiation transport in GRBs at cosmological distances (Goodman 1986; Paczyński 1986) predict very brief bursts of radiation with thermalized spectra, in contradiction to observation, while attempts (Brainerd 1992; Harding & Leventhal 1992; Katz 1992; Li & Dermer 1992; Li & Liang 1992; Mao & Paczyński 1992b; Wasserman 1992) to explain the spatial anisotropy and log *N* versus log *S* or *V/V*_{max} distributions of GRBs in Galactic models require the assumption of a spherically symmetric halo of ~100 kpc radius. Finally, the soft gamma repeaters (SGRs) introduce additional confusion. The fact of their repetition and the identification of one of them (1979 March 5; Cline 1980) with a supernova remnant in the LMC point strongly to a Galactic population, while the presence of spectral features and an 8 second periodicity (1979 March 5) indicate origin on a magnetic neutron star. However, it is unclear that SGRs should be

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(NASA-CR-194762) TWO POPULATIONS AND MODELS OF GAMMA RAY BURSTS (Washington Univ.) 19 p

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considered GRBs at all because their properties, including their spectra, are very different, and arguments made for SGRs may be irrelevant to the problems of GRBs.

As a first step toward resolving the apparent inconsistencies, I consider the obvious possibility that there are two distinct populations of GRBs. A cosmological population C accounts for most GRBs, and explains the statistics of isotropy and the $\log N$ versus $\log S$ or V/V_{\max} distributions. A population G includes those GRBs with spectral lines (a minority), and originate on neutron stars at Galactic distances. An individual GRB cannot be assigned to a population unless it shows spectral lines, but the majority (probably overwhelming) of those without spectral lines must be members of population C. The SGR may be members of population G. I do not consider more exotic possibilities, such as GRBs arising within the Oort cloud (Ruderman 1975), because they all seem far-fetched even though I myself have discussed one of them (Katz 1993). This suggestion of a cosmological population (to explain the statistics) and a Galactic population (to explain the spectral lines) is not related to the suggestion of Lingenfelter & Higdon (1992) (disputed by Paczyński 1992b) that the statistics may be explained by two Galactic populations.

In the absence of information about the intensity statistics and angular distribution of population G alone, it is not possible to discriminate between a disk and a halo origin. It will be important to obtain this information, which may be reducible from archival data.

In § 2, I discuss a possible mechanism for GRBs at cosmological distances, building on recent suggestions by others. In § 3, I more briefly discuss magnetic reconnection models of GRBs at Galactic distances. The 1979 March 5 event in the LMC poses an acute problem of gamma-gamma pair production (Carrigan & Katz 1992) which must be faced, whether or not GRBs of population G have comparable distances and luminosities, and even if there is no population G. § 4 contains a summary discussion. Unfortunately, unambiguous observational tests of the ideas discussed here will not be easy.

2. POPULATION C: GAMMA-RAY BURSTERS AT COSMOLOGICAL DISTANCES

The well-known failure of straightforward fireball models to explain the spectral and temporal properties of GRB led Shemi & Piran (1990) to consider neutrino-produced fireballs loaded with small (but not zero) amounts of ordinary matter; they found that the fireball could couple nearly all of its energy to the matter and (with the right values of the parameters) could accelerate it to relativistic velocity. Rees & Mészáros (1992) and Mészáros & Rees (1993) then pointed out that the interaction of this relativistic debris with surrounding matter, such as that suggested by Narayan, Paczyński, & Piran (1992), might be characterized by times consistent with the range of GRB rise times and durations of (10^{-3} – 10^3 s).

I build on these ideas. The environments of GRBs at cosmological distances are open to much speculation (for example, are they low-density Galactic halos or dense nuclei of galaxies?), but the strong clumpiness of interstellar matter is a consequence of immutable atomic physics (cooling rates), and isolated discrete clouds are likely under a very wide range of conditions. The rarity of GRBs makes it possible to assume favorable conditions, if these lie in the range of plausibility; there is no great difficulty if a considerably larger number of fireballs occurring in less favorable circumstances do not produce observable GRBs.

GRBs at cosmological distances require the radiation of $\sim 10^{51}$ ergs in observable gamma-rays. The complex chain of processes which lead to gamma-ray emission must be moderately efficient when the parameters do have favorable values, because the energy radiated as neutrinos by neutron star collapse, formation, or coalescence is unlikely to exceed $0.3 M_{\odot}$.

$c^2 \approx 5 \times 10^{53}$ ergs, and may be considerably less. Thus, while we are entitled to assume favorable circumstances to explain the rare observable GRBs, when these circumstances occur the resulting processes must be reasonably efficient. The fraction of neutrino energy converted to an electromagnetic fireball is small. Efficient conversion requires neutrino-neutrino collisions at angles in excess of 90° , but the neutrinos are generally expanding outward in a neutrino fireball, with velocity vectors which are tending toward radial outflow. Optimal head-on collisions are particularly rare. The conversion of electromagnetic energy to particle kinetic energy also has an efficiency less than 1. The final conversion to observable gamma-rays is the hardest part of the problem; it is easy to see how this could fail entirely.

Relativistic invariants alone limit the amount of kinetic energy available for radiation by fireball debris. If a relativistic debris cloud with speed $\beta_F c$, Lorentz factor γ_F , and proper mass per unit area σ sweeps up a proper mass per unit area $\alpha\sigma$, then the efficiency of radiation of the debris kinetic energy can be as large as

$$\epsilon = \frac{\alpha}{\alpha + \gamma_F(1 - \beta_F)} \quad (1)$$

Values of $\epsilon > \frac{1}{2}$ are obtained for $\alpha > 1/[\gamma_F(1 + \beta_F)] \approx 1/(2\gamma_F)$. Collisions with a very broad range of clouds of circum-fireball matter are consistent with efficient conversion of kinetic energy to radiation. This is fortunate, because efficient production of GRBs requires that common circum-fireball environments produce observable GRBs in most directions; it is not possible to insist on fortuitous geometries or on special values of the parameters.

Most of the kinetic energy will become available when the debris has swept up only a very little matter. If the energy of the explosion is Y then the proper mass of debris is $Y/\gamma_F c^2$, and half the kinetic energy will become available when the swept-up proper mass is $Y/2\gamma_F^2 c^2$. In a uniform medium of density ρ this will occur at an interaction radius

$$r_I = \left(\frac{3Y}{8\pi\gamma_F^2 c^2 \rho} \right)^{1/3} \sim 2 \times 10^{15} \text{ cm}, \quad (2)$$

where the numerical estimate assumed $Y = 10^{52}$ ergs, $\rho = 10^{-24} \text{ g cm}^{-3}$ and $\gamma_F = 10^4$.

The hardest part of the problem is turning the kinetic energy of the relativistic debris into the observed gamma-rays. The collision length of relativistic protons in ordinary matter is about 50 g cm^{-2} , or about 100 Mpc at typical interstellar densities. Clearly, some collective process is necessary, and it must couple the proton and ion energy into that of electrons, which radiate more readily. Even relativistic electrons do not radiate rapidly under interstellar conditions; the radiation length of a 10^{13} eV electron (corresponding to equipartition with a $\gamma_F \sim 10^4$ proton) for Compton scattering on a 3 K blackbody radiation field is $\sim 10^{23}$ cm, excessive by many orders of magnitude.

In order to obtain short pulses of radiation at distances of order those given by equation (2) it is necessary that a coherent relativistically expanding front of radiating particles be directed nearly toward the observer. It is not sufficient that individual particles be observed only when directed toward the observer, a condition met by most relativistic radiation processes. Therefore, ambient magnetic fields must not deflect particles significantly from their initial spherical expansion. This condition will be satisfied if the magnetic energy E_{mag} in the interaction sphere of radius r_I is very much less than the debris energy, so the debris can sweep away the ambient magnetic field without significant deflection. If equipartition is assumed between the ambient magnetic field B_{ISM} and an ambient turbulent velocity field v_{ISM} , then $E_{\text{mag}}/Y \sim v_{\text{ISM}}^2/\gamma_F^2 c^2 \ll 1/\gamma_F^2$, so

proton

relativistic

that the ambient field may safely be ignored.

2.1. Shock Structure

When the debris shell collides with a cloud of ambient matter the resulting flow may be complex. If the shell and the cloud each initially had uniform density and velocity and negligible (on a relativistic scale) temperature, the geometry is slab-symmetric, and all bulk velocities are normal to the planes of symmetry, then the resulting shock structure is shown in Figure 1. There are two shocks S1 and S2 and, in general, a contact discontinuity CD separating shocked fireball debris from shocked cloud. The equations relating the conditions in the four regions are cumbersome, except in the special symmetric case in which debris and cloud initially had the same composition and proper density. In this case, which I assume, there is no contact discontinuity and conditions in regions 2 and 3 are identical, as are those in regions 1 and 4.

FIG. 1.—Flow geometry in frame of contact discontinuity CD. S1 and S2 are shocks and numbers denote regions of fluid.

The relativistic shock conditions (Landau & Lifshitz 1959) may be used to determine physical conditions. The thermodynamic variables are the proper internal energy density e , the proper pressure p , and the proper enthalpy density $w = e + p$. In the unshocked cloud $e_1 = n_1 m_a c^2$, where n_1 is the proper atomic number density and m_a is the proper mass per atom, and $p_1 = 0$. In the shocked cloud $p_2 = e_2/3 \gg e_1$ in the extreme-relativistic (ER) limit. I shall refer to the frame of the unshocked interstellar material as the local observer's frame; transformation to our frame requires application of the cosmological redshift. Then, to lowest nontrivial order in $e_1/e_2 \ll 1$, the velocities of the fluids with respect to the shock front S1 are

$$\frac{v_1}{c} = \left[\frac{(p_2 - p_1)(e_2 + p_1)}{(e_2 - e_1)(e_1 + p_2)} \right]^{1/2} \approx 1 - \frac{e_1}{e_2} \quad (3)$$

and¹

$$\frac{v_2}{c} = \left[\frac{(p_2 - p_1)(e_1 + p_2)}{(e_2 - e_1)(e_2 + p_1)} \right]^{1/2} \approx \frac{1}{3} \left(1 + 2 \frac{e_1}{e_2} \right). \quad (4)$$

¹ Note the assertion in the first edition of Landau and Lifshitz (1959) that in the ER limit $v_2 \rightarrow c/3^{1/2}$ is a typographical error; the correct limit $v_2 \rightarrow c/3$ is given in later editions.

The velocity discontinuity v_{12} between fluids 1 and 2, measured in the frame of either, is obtained from the relativistic expression for the subtraction of velocities

$$\frac{v_{12}}{c} = \frac{v_1/c - v_2/c}{1 - v_1 v_2/c^2} \approx 1 - 2 \frac{e_1}{e_2}; \quad (5)$$

this velocity is also the velocity v_{2L} of shocked fluid 2 in the local observer's frame. The velocity v_1 is also the speed of the shock S1 in that frame.

Fluids 2 and 3 have the same velocity and, given our assumptions that $n_1 = n_4$ and $e_1 = e_4$, the same values of the thermodynamic variables. Then the velocity of fluid 3 in the frame of shock S2 is $-v_2$, and the speed of fluid 4 in that same frame is $-v_1$. The expressions for combinations of relativistic velocities may then be used to obtain the following results in the local observer's frame:

$$\frac{v_3}{c} = \frac{v_{12}}{c} \approx 1 - 2 \frac{e_1}{e_2}, \quad (6)$$

observer's

$$\frac{v_{S2}}{c} \approx 1 - 4 \frac{e_1}{e_2}, \quad (7)$$

$$\frac{v_4}{c} \approx 1 - 2 \left(\frac{e_1}{e_2} \right)^2. \quad (8)$$

It is now possible to calculate e_2 from the debris Lorentz factor γ_F , defined in the local observer's frame, using equation (8):

$$\gamma_F \equiv \frac{1}{[1 - (v_4/c)^2]^{1/2}} \approx \frac{e_2}{2e_1}, \quad (9)$$

$$e_2 \approx 2n_1 m_a c^2 \gamma_F. \quad (10)$$

The Lorentz factor γ_{2L} of the shocked material in the local observer's frame is obtained from $v_{2L} = v_{12}$, equations (5) and (10):

$$\gamma_{2L} \approx \gamma_F^{1/2}. \quad (11)$$

The detailed mechanics of the shock are obscure, but must be collisionless in order to form a shock at all. The shocked matter need not be in thermodynamic equilibrium. Heating of the shocked material by plasma instabilities is the source of dissipation; the distribution functions of particle energies will not be (relativistic) Maxwellians, but are more likely to be power laws. The distribution of energy between electrons and ions is uncertain. In a highly relativistic shock, as we expect here, electrons and ions are kinematically very similar (identical in the ER limit), so I will assume that the distribution of particle energies is independent of species; roughly half the post-shock energy resides in electrons. Any neutral matter does not interact with the shock, so that the density n_1 refers only to the ionized component. At distances $\sim r_I$ (eq. [2]) the interstellar material will largely have been ionized by the flash of radiation associated with the fireball or by collision with debris.

Using the shock jump conditions for the proper enthalpy w , $w_1 = n_1 m_a c^2$, and the ER limit $w_2 \approx 4e_2/3$ yields the proper atomic density

$$n_2 \approx 2 \left(\frac{n_1 e_2}{m_a c^2} \right)^{1/2}. \quad (12)$$

Define γ_2 by the relation

$$\gamma_2 \equiv \frac{e_2}{n_2 m_a c^2}; \quad (13)$$

then the mean energy (in the frame of fluid 2) per particle is $\gamma_2 \mu m_p c^2$, where μm_p is the mean proper mass per particle. For pure ionized hydrogen $\mu = 0.5$, while for the usual cosmic abundances (fully ionized) $\mu = 0.62$. The mean Lorentz factor of an electron (in the frame of fluid 2) is

$$\gamma_{2e} = \gamma_2 \frac{\mu m_p}{m_e}. \quad (14)$$

The proper density n_2 (eq. [12]) may be rewritten, using equation (13), as

$$n_2 \approx 4n_1 \gamma_2. \quad (15)$$

reproducing the result $n_2 = 4n_1$ for a strong but nonrelativistic shock ($\gamma_2 \rightarrow 1$) in a gas with an adiabatic exponent of 5/3. The Lorentz factor γ_2 is found from its definition (eq. [13]), and equations (9) and (15):

$$\gamma_2 \approx \frac{1}{2} \left(\frac{e_2}{e_1} \right)^{1/2} \approx \left(\frac{\gamma_F}{2} \right)^{1/2}. \quad (16)$$

In the local observer's frame most of the particles are narrowly collimated in the direction of the motion of fluid 2, and the

typical Lorentz factor is larger than those given in equations (14) and (16) by a factor $\sim \gamma_{2L}$ (eq. [11]). The angular width of collimation depends on the angular distribution of the particle momenta in the proper frame of fluid 2, which is unknown. If this is isotropic, then the locally observed angular width (for electrons as well as ions)

$$\theta_0 \sim \gamma_{2L}^{-1} \approx \gamma_F^{-1/2} \quad (17)$$

exponent $-1/2$

Note that this is a much broader angular distribution than the locally observed radiation pattern from a single particle, whose Lorentz factor is $\sim \gamma_F$ (ions) or $\sim \gamma_F m_p/m_e$ (electrons).

2.2. Time Dependence

The geometry of radiation from an advancing spherical shock is shown in Figure 2. The distant (but cosmologically local) observer may first see a flash of radiation from the fireball, itself, whose arrival time is taken as $t = 0$. If the fireball is a consequence of the merger of binary neutron stars (Eichler et al. 1989), the initial pulse includes bursts of neutrino and gravitational radiation, as well as electromagnetic radiation. Their emission is essentially simultaneous, although their arrival may be affected by dispersion arising from plasma refraction, neutrino rest mass (if any), etc. The initial electromagnetic flash is expected to have a thermal spectrum and to be extremely brief ($\ll 10^{-4}$ s) because of the small size of the fireball (Goodman 1986; Carrigan & Katz 1992); if, as assumed here, the fireball energy is largely converted (Shemi & Piran 1990) to kinetic energy of debris this initial flash may be unobservably faint. However, if it is observed the time interval between it and the rest of the GRB is an important constraint on the emission geometry.

FIG. 2.—Emission geometry

Radiation emitted from a point (r, θ) on the expanding spherical shell arrives at the observer at a time

$$t \approx \frac{r(1 - \cos \theta)}{c} + r \left(\frac{1}{v} - \frac{1}{c} \right) \approx \frac{r(1 - \cos \theta)}{c} + \frac{r}{u}, \quad (18)$$

where v is the shell's expansion velocity and the parameter (dimensionally but not physically a velocity) $u \equiv vc/(c - v)$. If the angular distribution of radiated intensity, measured in the local observer's frame, is $f(\theta')$, where θ' is the angle from the normal to the radiating surface, then the energy dE radiated by a patch of area dA is

$$dE = f(\theta') dA. \quad (19)$$

Radiation directed toward the observer has $\theta' = \theta$. Using $dA = 2\pi r^2 \sin \theta d\theta$ and $dt = r \sin \theta d\theta/c$ yields the observed power

$$P(t, r) = \frac{dE}{dt} = 2\pi c r f(\theta). \quad (20)$$

The function $f(\theta)$ is proportional to a convolution of the angular distribution of the radiating particles and their radiation pattern; as previously discussed, the latter is expected to be narrower than the former. A plausible guess is then

$$f(\theta) \propto \frac{1}{\theta^2 + \theta_0^2}, \quad (21)$$

where the angular width $\theta_0 \sim \gamma_F^{-1/2}$ is essentially the same as that of the momentum distribution (eq. [17]). The pulse shape is then obtained from equations (20) and (21), using equation

exponent $-1/2$

$$P(t, r) \propto \begin{cases} \frac{r}{[(t - r/u)2c/r + \theta_0^2]} & t \geq \frac{r}{u} \\ 0, & t < \frac{r}{u} \end{cases} \quad (22)$$

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where the approximation $\cos \theta \approx 1 - \theta^2/2$ has been used. The pulse shape $P(t, r)$ has been plotted in Figure 3, where the dimensionless parameter $\tau \equiv 2ct/r\theta_0^2$ has been defined.

FIG. 3.—Radiation pulse emitted at radius r , plotted for $\theta_0^2 = 20c/u$. Dashed lines denote rise and fall for a cloud sharply circumscribed between $\theta_1 = (\tau_1\theta_0^2 - 2c/u)^{1/2}$ and $\theta_2 = (\tau_2\theta_0^2 - 2c/u)^{1/2}$. Note that abrupt rise at $\tau = 2c/(\theta_0^2 u)$ is obtained for a uniformly filled medium, independent of the angular distribution of radiation, and does not require a cloud bounded in an angle. Units of the ordinate are arbitrary.

The pulse form of equation (22) should be regarded only as an envelope, for the actual pulse shape will be modulated by the spatial distribution of matter which the debris shell sweeps up. One striking feature of equation (22) is its abrupt rise and pronounced time-skewness (Weisskopf et al. 1978), consistent with the observed rapid rises of some GRBs (Bhat et al. 1992; Fishman 1993). The characteristic width of this function is

$$\Delta t \sim \frac{r\theta_0^2}{2c} \sim \frac{r}{\gamma_F c}. \quad (23)$$

Use of $r \sim r_1$ (eq. [2]) and $\gamma_F \sim 10^4$ yields $\Delta t \sim 10$ s, the right order of magnitude for the duration of GRBs. Much longer or shorter Δt may be possible for plausibly different values of the parameters, particularly the cloud density, which is uncertain even to order of magnitude.

The debris shell and shock propagate into a very heterogeneous medium. The effects of structure in θ are shown by the dashed lines in Figure 3, which assume a cloud uniform in the range $\theta_1 \leq \theta \leq \theta_2$, with abrupt boundaries. A more realistic gradual density profile or shape would produce a gradual rise and decay; the abrupt rise remains if (and only if) the cloud includes the line $\theta = 0$. Thus this abrupt rise is expected for some, but perhaps not all, GRBs, in accord with observations.

A complete intensity profile of a GRB requires the integration of equation (22) over r :

$$\mathcal{P}(t) = \int_0^{ut} P(t, r)g(r)dr, \quad (24)$$

where the weighting function $g(r)$ includes both the fact that the energy available for radiation falls off for $r > r_1$ and the effects of clumpiness of the ambient matter as a function of r . The observable region is a narrow half-cone of apical angle $\sim \theta_0$: unless the scale of spatial structure is $< r_1\theta_0 \sim 2 \times 10^{13}$ cm, clumpiness is more likely to be apparent as a function of r than of θ , justifying the use here of equation (22) which ignored any dependence of density on θ .

In the absence of spatial heterogeneity $g(r)$ may be taken to impose a cutoff at $r \approx r_1$, so that

$$\mathcal{P}(t) \approx \int_0^{\min(r_1, ut)} P(t, r)dr \propto \int_0^{\min(r_1, ut)} \frac{r^2 dr}{2ct + r[-(2c/u) + \theta_0^2]}. \quad (25)$$

The integral is elementary, but cumbersome, and of limited quantitative interest because of the artificiality of the assumed uniform density; the rise time is r_1/u . Two possibilities should

1. For a shock S1 propagating through a homogeneous medium $v/c \approx 1 - e_1/e_2 \approx 1 - 1/(2\gamma_F)$ and $u \approx 2\gamma_F c$. Then, for r_I given by equation (2) and $\gamma_F = 10^4$, the rise time is several seconds, given by equation (23) and is inconsistent with very rapid rise times. This corresponds to a long GRB pulse envelope.

2. On the other hand, the fireball debris propagates through a vacuum or very low density intercloud medium with a speed $v/c = (1 - 1/\gamma_F^2)^{1/2}$, so that $u \approx 2\gamma_F^2 c$; impacts upon small discrete clouds scattered within a region of size $\sim r_I$ will only introduce a timewidth $\sim 10^{-3}$ s, or less. This might not measurably broaden the abrupt rise given by equation (22) and in Figure 3.

Integration over r introduces two broadenings, one $O(r_I/2\gamma_F^2 c)$ associated with the entire emission region of size $\sim r_I$, where the low intercloud density is appropriate, and another $O(r_c/2\gamma_F c)$ associated with individual clouds of size $r_c \ll r_I$. These broadenings may each be much less than the envelope width (eq. [23]), permitting an observed signal resembling that of Figure 3. Several sub-pulses may be observed if the debris shell collides with several isolated clouds in a medium dense enough to slow the intercloud shock, so that interpulse times are $O(r_I/2\gamma_F c)$.

The actual situation is much more complicated than can be discussed here. For example, debris shell and clouds need not have the same proper densities, and each is likely to be spatially heterogeneous. Shocks propagating through heterogeneous media vary their strength, and produce associated continuous rarefactions and compressions. It is plausible that the complex structure of observed GRBs could be explained by the interaction of relativistic debris shells with clumpy media, but more quantitative results would require numerical relativistic hydrodynamic calculations.

2.3. Radiation

The hardest part of the problem is turning electron kinetic energy into the observed radiation. Even though the typical Lorentz factor of an electron in the local observer's frame is $\sim \gamma_{2L} \gamma_{2e} \sim \mu \gamma_F m_p/m_e$ (from eqs. [11], [14], and [16]), their rate of synchrotron radiation and Compton scattering in plausible interstellar magnetic and radiation fields is low. I therefore make the radical suggestion that the collisionless shock produces approximate equipartition between the magnetic energy density and the particle energy density. The proper magnetic field B_2 and energy density in fluid 2 are then, using equation (10),

$$\frac{B_2^2}{8\pi} = \zeta e_2 \approx 2\zeta n_1 \gamma_F m_p c^2, \quad (26)$$

where $\zeta \leq \frac{1}{2}$ is a phenomenological parameter describing the approach to equipartition. The synchrotron energy loss time for an electron with Lorentz factor given by equation (14), assuming no correlation between the direction of the electron momentum and that of the magnetic field, is then, taking pure hydrogen composition,

$$t_{r2} \approx \zeta^{-1} \left(\frac{n_1}{1 \text{ cm}^{-3}} \right)^{-1} \gamma_F^{-3/2} 1.1 \times 10^7 \text{ s}. \quad (27)$$

The radiating volume is moving toward the observer with a bulk Lorentz factor γ_{2L} (eq. [11]), so that application of a Lorentz transform yields the local observer's measured radiation time

$$t_{\text{obs}} = \gamma_{2L} \left(\frac{t_{r2} - \frac{v_{2L}^2 t_{r2}}{c^2}}{c} \right) \\ = \frac{t_{r2}}{\gamma_{2L}} \approx \zeta^{-1} \left(\frac{n_1}{1 \text{ cm}^{-3}} \right)^{-1} \gamma_F^{-2} 1.1 \times 10^7 \text{ s}. \quad (28)$$

insert comma after (14)

$$\left(t_{r2} - \frac{v_{2L}^2 t_{r2}}{c^2} \right)$$

For a plausible interstellar cloud density $n_1 > 1 \text{ cm}^{-3}$ and $\gamma_F \approx 10^5$, t_{obs} may be a millisecond or less, as required by rapid time-structure in some GRBs (Bhat et al. 1992). This justifies the assumption, made implicitly in the discussion of GRB rise times and pulse lengths, that shock-accelerated electrons radiate instantaneously; properly, the pulse profiles predicted by equations (24) and (25) should be convolved with a broadening function which includes the radiation time, and which has a width t_{obs} in the local observer's frame.

The characteristic frequency of synchrotron radiation, measured in the frame of fluid 2, is obtained from standard expressions using equations (14) and (26). For pure hydrogen the result is

$$\nu_2 \sim \zeta^{1/2} \left(\frac{n_1}{1 \text{ cm}^{-3}} \right)^{1/2} \gamma_F^{3/2} 3 \times 10^{11} \text{ s}^{-1}, \quad (29)$$

while Lorentz transformation to the local observer's frame, using equation (11), yields

$$\nu_{\text{obs}} \sim \zeta^{1/2} \left(\frac{n_1}{1 \text{ cm}^{-3}} \right)^{1/2} \gamma_F^2 3 \times 10^{11} \text{ s}^{-1}. \quad (30)$$

Mev photons may be observed for $\zeta = \frac{1}{2}$ if $n_1 \sim 1 \text{ cm}^{-3}$ and $\gamma_F \sim 4 \times 10^4$, for example. Note that if $\gamma_F < 5 \times 10^5 \zeta^{-1/3} (n_1 / 1 \text{ cm}^{-3})^{-1/3}$ then the photon energies are below the pair production threshold (this condition is Lorentz-invariant, but is most easily evaluated using eq. [29] in the frame of the shocked fluid, in which the photon distribution is isotropic), and there can be no pair-production catastrophe.

It has also been observed (Fishman 1993) that many GRBs, or subpulses within them, show a progressive spectral softening with time. This is qualitatively explained using Figure 2 and equation (18). If the radiation field is isotropic in the frame of fluid 2 (as will be the case if the particle distribution and magnetic field directions are isotropic) and has a characteristic photon energy, then in the local observer's frame the spectral hardness above this characteristic spectral peak will be a decreasing function of θ , because of the angular dependence of the Doppler shift. Higher frequency photons are preferentially observed from smaller values of θ , which arrive earlier in the burst or subpulse, while lower frequency photons are observed over a wider range of θ and hence over a longer time. A quantitative prediction for the spectral evolution with time could be made by numerical integration of the synchrotron emission function, but would depend on (uncertain) assumptions made regarding the energy and angular distributions of the radiating electrons. Spectral softening within subpulses is also produced by the more rapid radiation of the most energetic electrons. The decline of γ_F as more interstellar matter is swept up leads to a spectral softening throughout a GRB, from one subpulse to the next.

When shock-accelerated relativistic electrons enter the interstellar medium they may coherently radiate synchrotron radiation, in analogy to the coherent electromagnetic pulse (EMP) produced by terrestrial high-altitude nuclear explosions (Karzas & Latter 1962, 1965). Although radiation rates by individual electrons in interstellar fields are low, the coherent emission of many electrons may be significant. The observed synchrotron frequency of this natural EMP may extend up to the visible or ultraviolet bands, but the intensity may be reduced at higher frequencies at which the electrons do not radiate coherently. The cutoff frequency depends on the electron density as well as on their Lorentz factor, and is difficult to estimate. The problem is particularly complex because the debris energy density far exceeds that of the ambient field, which is strongly affected by the entry of the relativistic particles (even aside from the plasma processes in the collisionless shock), in contrast to the terrestrial case. At lower frequencies an initial impulsive pulse is strongly affected by intergalactic

plasma dispersion, with an electron column density $\sim 10^{23}(\Omega_p/0.1)$ electrons cm^{-2} , where Ω_p is the fraction of the cosmological closure density present in ionized intergalactic gas. The resulting pulse may be very long (10^9 Hz radiation is delayed by a few minutes for $\Omega_p = 0.1$), and therefore difficult to detect. If it were observed it would have a characteristic dependence of frequency on time [$\nu \propto (t - t_0)^{1/2}$, a "cosmic whistler", although in a terrestrial whistler $\nu \propto (t - t_0)^{2/3}$; Stix 1962] which would directly determine the mean intergalactic electron density. Unfortunately, the limiting amplitude of the radiated EMP field is comparable to the ambient interstellar field in the source region (the EMP field screens the electron current from the ambient field, and limits the radiated intensity), leading to an upper bound $\sim B_{\text{ISM}}^2 cr^2/2 \sim 10^{30}$ ergs s^{-1} on the radiated power which is very small.

250
exponent $-1/2$
exponent -2

3. POPULATION G: GALACTIC GAMMA-RAY BURSTERS

GRBs which show spectral features, typically around a few tens of keV and at 400 keV, have long been identified with Galactic magnetic neutron stars and are inexplicable at cosmological distances. Their distances cannot be determined from available data, and could be less than 100 pc, ~100 kpc, or anything in between. The familiar arguments concerning the mechanisms of GRBs at Galactic distances center on two issues: the source of energy, and the physical conditions in the emitting region. The problems are harder, the greater the assumed distances. The observation of the 1979 March 5 event at a likely distance of 55 kpc (Cline 1980) forces the consideration of distances of that order, and of correspondingly high luminosities, even though it is unclear whether it (a SGR) was a member of the population G of GRBs, or represented a distinct third class of objects.

The central problem of distant and luminous gamma-ray sources is gamma-gamma pair production (Cavallo & Rees 1978; Schmidt 1978; Katz 1982; Epstein 1985; Carrigan & Katz 1992). This process does not permit the escape of a large luminosity of MeV gamma-rays from a small region unless they are collimated, and thus excludes many models of GRB or SGR at Galactic halo or cosmological distances. It is well known that this problem is avoided in a collimated relativistic outflow of radiating matter, a consideration which led to the popularity of fireball models, in which an opaque cloud of radiation and pair gas adiabatically expands and cools until its particles' velocity vectors are collimated outward. However, fireballs are conspicuously incapable of producing low redshift (400 keV) annihilation lines, line features at tens of keV, or the observed long and complex time structure. The interaction of fireballs with their environment may solve the temporal problem, as discussed in § 2, but offers no hope of solving the spectral problem. The case for magnetic neutron stars for GRBs or SGRs with spectral lines remains as strong as the data.

features

It is usually assumed that the radiating region of a GRB or SGR in a neutron star model is dominated by pair plasma, with $n_+ \gg n_i$, where n_+ and n_i are the positron and ion densities, respectively. This assumption is made, in analogy to fireballs, even for nonfireball models, perhaps because the threatened gamma-gamma pair production catastrophe seems a likely source of dense pair plasma, and because the observed annihilation line requires a source of positrons. However, the assumption of pair dominance may not be justified in non-fireball models, such as are required to explain population G GRBs. If sufficient gamma-ray collimation is present to avoid a gamma-gamma pair production catastrophe, then the production of pairs may be negligibly small. When the observation of annihilation radiation provides empirical evidence for the production of some positrons, it should be remembered that an observably narrow annihilation line requires temperatures less than 50 keV, and may be produced by a comparatively small number of positrons precipitated onto the cool neutron star surface; a hot pair plasma does not produce a recognizable annihilation line.

If a pair plasma is not an expanding fireball, it must be trapped on magnetic field lines (Carrigan & Katz 1992). Gravitation is unimportant for pairs, so they fill a magnetosphere (presumably of a magnetic neutron star). However, they quickly (in a free-flight time) precipitate onto the stellar surface, where they annihilate, because they more rapidly radiate their transverse momentum by the cyclotron process; even if the radiation density is not sufficient to maintain most leptons in excited Landau (magnetic) states (a condition satisfied under only the most extreme conditions), their interaction with the radiation field destroys their transverse adiabatic

invariant. In this magnetosphere-filling geometry the emergent radiation, by whatever process, is not collimated, and gamma-gamma pair production imposes its usual limits on the emergent flux of Mev gamma-rays.

It may be more satisfactory to consider an electron-ion plasma with only a small admixture of positrons (sufficient to produce the observed annihilation line), so that $n_+ \ll n_-$. An electric field may accelerate with the electrons, which radiate by bremsstrahlung or by the cyclotron process after elastic scattering on the ions raises them to excited Landau states. Because most of the leptons are negative, they form a broadly collimated beam; if they are relativistic the resulting radiation is similarly collimated and there is no gamma-gamma pair production catastrophe or limit (other than a Planck function at an effective temperature characterizing the electron distribution function) on the emergent intensity.

Planck

In contrast, an electric field acting on a pair gas heats it but imparts no net momentum to the leptons; two counterstreaming beams of gamma-rays readily produce pairs, rapidly achieving equilibrium with them and limiting the emergent intensity. A minority admixture of positrons in an electron-ion plasma produces only a proportionately small counter current of gamma-rays to those produced by the electrons. This countercurrent removes an equal current of electron-produced gamma-rays by pair production, but the remaining electron-produced gamma-rays form a collimated beam and escape comparatively freely, suffering little or no (depending on the degree of collimation) gamma-gamma pair production. This is a consequence of the net momentum imparted to the lepton-photon system by the electric field in analogy to the momentum imparted to a sector of a fireball by adiabatic expansion.

It is possible to make simple rough estimates of the parameters of an electrically heated ion-electron sheet plasma, which might be the source region of a GRB of population G, following Katz (1993). Consider a sheet of thickness L , be composed of positive ions of charge Z and density n_i ($n_e = Zn_i$), have with transverse optical depth τ and temperature T , and radiating power per unit area P . Define the dimensionless power per unit area $p \equiv Ph^3/(m_e^4 c^6)$, dimensionless thickness $l \equiv Lm_e c^2/e^2$, and temperature $t \equiv k_B T/(m_e c^2)$, where these quantities have been scaled to values characteristic of a relativistic electron (or pair) gas. The characteristic radiant intensity $m_e^4 c^6/h^3 = 4.3 \times 10^{35}$ ergs/(cm² s) and length $e^2/m_e c^2 = 2.8 \times 10^{-13}$ cm (the classical electron radius). The optical depth is

radiating

$$\tau \sim n_e \sigma_0 L, \quad (31)$$

where the characteristic cross section $\sigma_0 \equiv e^4/(m_e^2 c^4)$. At non-relativistic energies the appropriate cross section is the Thomson cross section $8\pi\sigma_0/3$, but at semi-relativistic energies of interest σ_0 may be a fair approximation. The observation of a nonthermal spectrum implies that τ cannot much exceed unity, but $\tau \sim 1$ and $\tau \ll 1$ are each possible.

The large field of a magnetized neutron star has a number of effects. It enters the argument of the effective Coulomb logarithm in collisional processes, typically reducing it to $\ln \Lambda \approx \ln [k_B T m_e c / (\hbar e B)]$ (Katz 1982). Both bulk motion and current flow are restricted to be parallel to the field lines, justifying the assumption of thin sheet geometry and making the field distribution nearly force-free ($\mathbf{J} \times \mathbf{B} = 0$). Perhaps most important, it means that any electron energy resulting from motion perpendicular to the field is immediately radiated. Even in conditions characteristic of the 1979 March 5 event at 55 kpc the radiation density is far below Planckian (for $t \sim 1$), so that electrons may be assumed to be in their ground magnetic state until collisionally excited, and then to radiate as if in vacuum. The radiation rate per unit area may therefore be estimated using standard expressions for elastic scattering:

$$P \sim n_e n_i k_B T \left(\frac{k_B T}{m_e} \right)^{1/2} \sigma_0 Z^2 \left(\frac{m_e c^2}{k_B T} \right)^2 L \sim \frac{n_e^2 e^4 LZ \ln \Lambda}{m_e c t^{1/2}}.$$

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This expression may be inverted, using the definitions of p , l , t , and $\alpha \equiv e^2/(\hbar c)$, to give

$$n_e \sim \left(\frac{pt^{1/2}}{l\alpha^3 Z \ln \Lambda} \right)^{1/2} \frac{m_e^3 c^3}{\hbar^3}; \quad (32)$$

the characteristic density $m_e^3 c^3/\hbar^3 = 1.8 \times 10^{31} \text{ cm}^{-3}$. An alternative expression for n_e is obtained from equation (31):

$$n_e \sim \frac{\tau}{l\alpha^3} \frac{m_e^3 c^3}{\hbar^3}. \quad (34)$$

Equating the expressions (33) and (34) yields a result for the thickness:

$$l \sim \frac{\tau^2 Z \ln \Lambda}{pt^{1/2}\alpha^3}. \quad (35)$$

Very roughly, $p \sim 10^{-4}$ for the 1979 March 5 event, if in the LMC, and $p \sim 10^{-7}$ for a typical observed GRB if at 100 kpc distance, corresponding to $L \sim 1.4\tau^2 \text{ cm}$ and $L \sim 1.4 \times 10^3 \tau^2 \text{ cm}$, respectively, where $Z = 26$, $t = 0.1$, and $\Lambda = 10$ were taken. The densities are correspondingly high.

If the radiating sheet is driven by magnetic reconnection, as is plausible, then the radiated power may be related to the electrical work done:

$$P \sim \sigma_{e1} E^2 L, \quad (36)$$

where E is the electric field (properly, its component parallel to \mathbf{B}) and a nonrelativistic expression (Spitzer 1962) is used for the (cgs) electrical conductivity:

$$\sigma_{e1} \approx 2 \left(\frac{2}{\pi} \right)^{3/2} \frac{(k_B T)^{3/2}}{m_e^{1/2} e^2 \xi Z \ln \Lambda} \approx \frac{\tau^{3/2}}{\xi Z \ln \Lambda} \frac{m_e c^3}{e^2}, \quad (37)$$

where the parameter $\xi \geq 1$ is a correction factor which allows for the possibility of anomalous (plasma instability) resistivity when current densities and electron drift velocities are large and for the decrease in conductivity when electron velocities approach c . The characteristic conductivity $m_e c^3/e^2 = 1.1 \times 10^{23} \text{ s}^{-1}$. Equations (36) and (37) may be combined with the definitions of the various dimensionless parameters to give the electric field, current density, and electron drift velocity:

$$E \sim \left(\frac{p\alpha^3 \xi Z \ln \Lambda}{\tau t^{3/2}} n_e m_e c^2 \right)^{1/2}; \quad (38)$$

$$j = \sigma_{e1} E \sim \frac{pt\alpha^3}{\tau \xi^{1/2} \ln \Lambda} \frac{m_e^3 c^7}{e^5}; \quad (39)$$

$$v_{df} \frac{j}{n_e e} \sim \frac{ct^{1/2}}{\xi^{1/2}} \sim \left(\frac{k_B T}{n_e} \right)^{1/2} \frac{1}{\xi^{1/2}}. \quad (40)$$

The characteristic current density $m_e^3 c^7/e^5 = 6.5 \times 10^{38} \text{ esu}/(\text{cm}^2 \text{ s})$. The drift velocity is thus comparable to the electron thermal velocity unless $\xi \gg 1$; ion-acoustic instability is likely unless $\xi > m_i/(Zm_e)$.

GRB mechanisms based on magnetic reconnection, as in solar flares, have been qualitatively discussed for many years (Ruderman 1975). They may explain GRB energetics and phenomenology (Katz 1982). Magnetic reconnection by sheet currents provides a natural explanation of the electrically heated sheets discussed in this section. If this model is assumed, then Maxwell's equations provide an additional relation among j , L , and the magnetic field B , permitting further constraints to be placed on the parameters. Because j is parallel to B , the direction of B rotates across a sheet current without changing its magnitude B . If the total angle of rotation across a uniform current sheet is π , then

$$L = \frac{Bc}{4j}. \quad (41)$$

Defining the usual characteristic magnetic field $B_c \equiv m_e^2 c^3 / (e\hbar) = 4.4 \times 10^{13}$ gauss and the dimensionless parameter $b \equiv B / (4B_c)$, equation (41) may be rewritten

$$l \sim \frac{\tau b \xi^{1/2} Z \ln \Lambda}{p t \alpha^2} \quad (42)$$

Equating this expression to equation (35) yields

$$\tau \sim \frac{b \alpha \xi^{1/2}}{t^{1/2}} \quad (43)$$

For plausible $t \sim 1$ and $b \sim 0.02$, $\tau \sim 10^{-4} \xi^{1/2}$; either the emission region is very optically thin or the resistivity is dominated by plasma wave scattering, and is far in excess of its independent particle value. Either or both of these possibilities is acceptable.

The electrical field (eq. [38]) may be evaluated, using equation (34) for n_e and equation (35) for l , and defining the characteristic electric field $E_c \equiv m_e^2 c^3 / (e\hbar) \equiv B_c = 4.4 \times 10^{13}$ cgs:

$$E \sim \left(\frac{p \xi Z \ln \Lambda}{l t^{3/2}} \frac{m_e^4 c^5}{\hbar^3} \right)^{1/2} \sim \frac{p \alpha^2 \xi^{1/2}}{\tau t^{1/2}} E_c \quad (44)$$

If the energy release is driven by magnetic reconnection it is proper to use equation (43) for τ , yielding

$$E \sim \frac{p \alpha}{b} E_c \quad (45)$$

At Galactic halo distances the brightest GRBs may have E sufficient to produce vacuum breakdown into a pair gas (Smith & Epstein 1993), but in less intense or closer GRBs this will not occur, and the resistively heated ion-electron plasma discussed here may be sufficient. On a microscopic level, of course, the power is determined by the magnetic field strength and configuration, and by the mechanisms of plasma resistivity which drive reconnection.

The observed nonthermal gamma-ray spectrum of GRBs requires the presence of a nonthermal distribution of "runaway" electrons, consistent with the large v_d , and collective processes discussed above. It may be relevant that values of $t \sim 1$, consistent with the radiation of the bulk of the power of GRBs at \sim Mev energies, correspond to a maximum in the conductivity and therefore, under conditions of a fixed potential drop, to a maximum in the power dissipation.

If $Z \gg 1$, σ_{ei} increases approximately linearly with n_+ in the range $Z n_i < n_+ < Z^2 n_i$ because of the increasing density of charge carriers without a corresponding decrease in their scattering length. This is in contrast to the usual near-independence of density of σ_{ei} . Pair production thus may provide a natural thermostat at $t \sim 1$, reducing the power dissipated in regimes at the pair production threshold by increasing the conductivity under conditions of constant current. It is evident, of course, that observable cyclotron and annihilation lines require much lower values of t , and are plausibly produced by energy and positrons precipitated on the dense surface layers of the neutron star, cooled by blackbody radiation.

4. DISCUSSION

The fundamental problem of GRB phenomenology is the apparent inconsistency between their spatial distribution, which points strongly toward a cosmological origin, and the spectral features observed in some GRBs, apparently inconsistent with such an origin. In this paper, I have tried to reconcile these apparently contradictory data by assuming two disjoint populations of GRBs. Because the argument for a cosmological population C is statistical, while the argument for a Galactic population G is based on the observation of spectral lines from only a minority of GRBs (perhaps from only a minority of that population), it is difficult to assign an individ-

this

electric

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ual GRB to either population, unless it shows spectral lines or is identified with another astronomical object (the SGR of 1979 March 5 is the only good extant example of such an identification).

Fortunately, the models discussed in this paper predict another potential distinguishing characteristic. GRBs in population C produce their radiation by the synchrotron process of relativistic electrons. If the magnetic field is ordered the radiation will be linearly polarized. GRBs in population G produce their radiation by the cyclotron process of semi-relativistic electrons, Coulomb scattered into excited Landau states. This radiation is elliptically polarized, with a substantial circular component. In contrast, radiation produced by annihilation in a pair gas, such as the initial burst from a fireball or the trapped pair gas discussed by Carrigan & Katz (1992), is unpolarized.

The predicted characteristic frequency of gamma-ray emission in fireball debris impact models (eq. [30]) is fairly sensitive to the initial fireball Lorentz factor γ_F . The value of γ_F depends on physical conditions within the fireball (Shemi & Piran 1990), and a wide range of γ_F would be expected, with a corresponding range in spectra. In particular, smaller values of γ_F would lead to X-ray, ultraviolet, or visible bursts without gamma-ray emission; these should be searched for. Bursts of lower frequency radiation should have longer durations and smoother time histories (eq. [28]), and perhaps also lower efficiencies, as accelerated electrons undergo adiabatic expansion (coupling their kinetic energy to ion motion) before they radiate.

Relativistic synchrotron models, such as those discussed here, generally involve a power-law distribution of radiating electrons up to a cutoff Lorentz factor γ_{2e} (eq. [14]), and therefore a power-law synchrotron spectrum up to a cutoff frequency ν_2 (eq. [29]). The power-law indices are uncertain, but using the usual estimate for the spectral index $p = (\eta + 2)/(\eta - 1)$ of relativistic particles accelerated by a shock of compression ratio η and the relation $s = (p - 1)/2$ for the synchrotron spectral index yields $s = 3/[2(\eta - 1)] = 3(\gamma_a - 1)/4 = 1/4$ for a strong shock in a relativistic fluid for which $\eta = 7$ and the adiabatic exponent $\gamma_a = 4/3$. Because $p = 3/2$ most of the particle energy resides in the most energetic particles, in contrast to the usual case with $p > 2$. The radiation spectrum $F_\nu \propto \nu^{-s}$ implies a visible power $\sim 10^{-4}$ of that in soft gamma-rays; a bright GRB of 10^{-5} ergs/(cm² s) extrapolates to $F_\nu \sim 5 \times 10^{-23}(\nu/10^9 \text{ Hz})^{-1/4}$ ergs/(cm² s Hz). However self-absorption limits the brightness temperature at radio frequencies. When the debris shell has a radius r_s the effective radiating area is $\sim r_s^2/\gamma_F$ because of the collimation, and the brightness temperature $T_b < \gamma_F m_p c^2/k_B$. The resulting power spectral density at a distance D is $< 2\pi\nu^2 m_p r_s^2/D^2 \sim 3 \times 10^{-29}(\nu/10^{10} \text{ Hz})^2 (r_s/2 \times 10^{15} \text{ cm})^2$ ergs/(cm² s Hz). As the debris shell expands this may become significant. Lower energy electrons radiate more slowly in a given field, with $t_r \propto \nu^{-1/2}$ (in any frame), so that an observed t_r of 1 msec at 300 keV corresponds to a minimum pulse length of ~ 0.3 s in visible light and of $\sim 10^2$ s at 1 GHz. Radiation from expanded shells with $r_s \gg r_i$, where γ_F is too small to produce gamma-rays and B is small, will have much longer pulse lengths.

Relativistic electrons of lesser energy will also be produced after the debris have been slowed by interaction with surrounding gas (and gamma-ray emitting electrons are no longer accelerated), so that visible and radio-frequency pulses may follow the high energy pulse by seconds, minutes, or even hours, and be characterized by smaller γ_F and correspondingly longer durations.

During the lengthy process of review of this paper, several groups (Kouveliotou et al. 1993; Lamb, Graziani, & Smith 1993; and Mao, Narayan, & Piran 1993) reported that the durations of GRB are bimodally distributed. This may be

explained, using equation (28), as a consequence of the well-known bimodal (or multi-modal) distribution of interstellar gas densities in ionized and neutral regions; the observed duration of a GRB is determined by γ_F (which may be distributed over a wide range) and the gas density in a small region (see eq. [2]) surrounding the fireball.

The natural next step in the development of population C GRB models described here would be to perform numerical one-dimensional special-relativistic hydrodynamic calculations of relativistic debris clouds propagating into model distributions of clumpy interstellar media. The predicted pulse shapes and spectra may then be calculated using equation (24) convolved with the radiation timescale of electrons as a function of their energy.

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APPENDIX

As relativistic debris sweeps up surrounding diffuse matter it forms a relativistic blast wave. The debris energy degrades. Its radiation shifts to lower frequencies and increases in duration, and may be observable at frequencies from X-rays down to radio. This appendix collects a few simple results concerning relativistic blast waves.

First, assume that a negligible fraction of the blast wave's energy is radiated, as in the Taylor-Sedov nonrelativistic blast wave. As long as the blast wave remains relativistic its radius $r \approx ct$, where r and t are its radius and age as measured by a local observer. The dependence of other quantities on r and t is more interesting. Conservation of energy gives

$$\gamma_F(r) \approx \frac{3}{4\pi} \frac{Y}{r^3 \rho c^2}. \quad (46)$$

Define a characteristic radius

$$r_0 \equiv \left(\frac{3Y}{4\pi \rho c^2 \gamma_0} \right)^{1/3} \sim r_1 \gamma_0^{1/3}, \quad (47)$$

where $\gamma_0 \equiv \gamma_F(0)$ is the original debris Lorentz factor, so that

$$\gamma_F(r) \sim \gamma_0 \left(\frac{r}{r_0} \right)^{-3}. \quad (48)$$

The relativistic particles, after being thermalized in a shock, are roughly isotropic in the frame of that shock and have an angular dispersion $\theta_0 \sim \gamma_F^{-1/2}$ (eq. [17]) as measured in the local observer's frame. As a result, they fill a shell of thickness

$$\Delta r \sim \left(\frac{2}{\gamma_F} \right) r, \quad (49)$$

where the chief contribution to Δr is their angular dispersion (their dispersion in speed contributes only a thickness $O[r/\gamma_F^2]$). The energy density measured in the local observer's frame is then

$$\mathcal{E}(r) \approx \frac{Y}{4\pi r^2 \Delta r} \sim \frac{3Y^2}{32\pi^2 r^6 \rho c^2}, \quad (50)$$

while the proper energy density is less by two factors of $\gamma_F^{1/2}$:

$$e(r) \sim \frac{Y}{8\pi r^3} \quad (51)$$

(using eq. [46], this is equivalent to eq. [10], up to factors $O[1]$). Then, as in equation (26), the magnetic field in the comoving frame is

$$B_2(r) \sim \left(\frac{\zeta Y}{r^3} \right)^{1/2}. \quad (52)$$

In this frame the electric fields are small and the usual synchrotron radiation expressions may be used (in the local observer's frame there are large electric fields which reduce the electrons' magnetic acceleration and permit them to remain collimated despite the presence of transverse magnetic fields).

Using the previous results for electron energies, the characteristic synchrotron radiation frequency, measured in the comoving frame, is

$$\nu_2(r) \sim \nu_2(r_0) \left(\frac{r}{r_0} \right)^{-9/2}, \quad \text{frak nu} \quad (53)$$

where $v_2(r_0)$ is given in equation (29). In the local observer's frame this frequency is

$$v_{\text{obs}}(r) \sim v_{\text{obs}}(r_0) \left(\frac{r}{r_0} \right)^{-6}, \quad (54)$$

where $v_{\text{obs}}(r_0)$ is given in equation (30). The electron energy-loss time, measured in the comoving frame, is

$$t_{\text{loss}}(r) \sim t_{r2}(r_0) \left(\frac{r}{r_0} \right)^{9/2}, \quad (55)$$

where t_{r2} is given in equation (27). In the local observer's frame this time is

$$t_{\text{obs}}(r) \sim t_{\text{obs}}(r_0) \left(\frac{r}{r_0} \right)^6, \quad (56)$$

where $t_{\text{obs}}(r_0)$ is given in equation (28).

The radiated power, measured in the comoving frame, is

$$P_{\text{cm}}(r) \sim \frac{2\pi}{9} \frac{\zeta e^4 Y \gamma_0 n_1}{m_e^2 c^3} \left(\frac{m_p}{m_e} \right)^2 \left(\frac{r}{r_0} \right)^{-3}, \quad (57)$$

while the radiated power measured in the local observer's frame is

$$P_{\text{obs}}(r) \sim \frac{2\pi}{9} \frac{\zeta e^4 Y \gamma_0^2 n_1}{m_e^2 c^3} \left(\frac{m_p}{m_e} \right)^2 \left(\frac{r}{r_0} \right)^{-6}. \quad (58)$$

The efficiency of radiation over a range $\sim r$ around radius r is

$$\epsilon(r) \sim \frac{P_{\text{obs}}(r)}{Y} \frac{r}{c\gamma_F^{1/2}(r)} \sim \frac{2\pi}{9} \frac{\zeta e^4 \gamma_0^{3/2} n_1 r_0}{m_e^2 c^4} \left(\frac{r}{r_0} \right)^{-7/2}; \quad (59)$$

the factor $\gamma_F^{1/2}(r)$ results from the Doppler shortening of the observed duration of the radiation.

Defining a characteristic time

$$t_c \equiv \frac{r}{c\gamma_F^{1/2}(r)} \sim \left(\frac{r}{r_0} \right)^{5/2} t_0, \quad (60)$$

where $t_0 \equiv r_0/(c\gamma_0^{1/2})$, yields the scaling laws:

$$\gamma_F(t_c) \sim \gamma_0 \left(\frac{t_c}{t_0} \right)^{-6/5}, \quad (61)$$

$$\mathcal{E}(t_c) \sim \mathcal{E}(t_0) \left(\frac{t_c}{t_0} \right)^{-12/5}, \quad (62)$$

$$e(t_c) \sim e(t_0) \left(\frac{t_c}{t_0} \right)^{-6/5}, \quad (63)$$

$$B_2(t_c) \sim B_2(t_0) \left(\frac{t_c}{t_0} \right)^{-3/5}, \quad (64)$$

$$v_2(t_c) \sim v_2(t_0) \left(\frac{t_c}{t_0} \right)^{-9/5}, \quad (65)$$

$$v_{\text{obs}}(t_c) \sim v_{\text{obs}}(t_0) \left(\frac{t_c}{t_0} \right)^{-12/5}, \quad (66)$$

$$P_{\text{cm}}(t_c) \sim P_{\text{cm}}(t_0) \left(\frac{t_c}{t_0} \right)^{-6/5}, \quad (67)$$

$$P_{\text{obs}}(t_c) \sim P_{\text{obs}}(t_0) \left(\frac{t_c}{t_0} \right)^{-12/5}, \quad (68)$$

$$\epsilon(t_c) \sim \epsilon(t_0) \left(\frac{t_c}{t_0} \right)^{-7/5}, \quad (69)$$

Numerical evaluation shows that $\epsilon(r_0) \gg 1$ for plausible parameters. Such a high nominal efficiency of radiation justifies the assumption of efficient gamma-ray production, but is inconsistent with the assumption of energy conservation made to obtain the blast wave results of equations (46)–(69). The opposite limit, in which the collision between fireball debris and interstellar matter is completely inelastic (analogous to the late-time "snowplow" limit of nonrelativistic blast wave theory) is easily treated.

In a narrow angular segment $d\Omega$ a debris proper mass $M d\Omega/(4\pi) = Y d\Omega/(4\pi\gamma_0 c^2)$ sweeps up an interstellar proper mass $\rho r^2 dr d\Omega$ in traveling a thickness dr . In the local observer's frame the interstellar matter is at rest, and the momentum of the debris is $M c(\gamma^2 - 1)^{1/2} d\Omega/(4\pi)$. After the collision the total proper mass is $[M/(4\pi) + \rho r^2 dr] d\Omega$. The Lorentz factor $\gamma + d\gamma$ of the new center-of-momentum frame is found from the expression of a relativistically moving object of the total proper mass:

for

$$\frac{M d\Omega}{4\pi} c(\gamma^2 - 1)^{1/2} = \left(\frac{M d\Omega}{4\pi} + \rho r^2 dr d\Omega \right) c[(\gamma + d\gamma)^2 - 1]^{1/2}. \quad (70)$$

If the radiation is emitted isotropically in the center-of-momentum frame its emission does not change the velocity of this frame. Elementary calculation yields the differential equation

$$\frac{d\gamma}{dr} = - \frac{4\pi r^2 \rho}{M} \frac{\gamma^2 - 1}{\gamma}. \quad (71)$$

Integration yields the result

$$\begin{aligned} \gamma^2 - 1 &= (\gamma_0^2 - 1) \exp [-8\pi r^3 \rho / (3M)] \\ &= (\gamma_0^2 - 1) \exp (-2r^3 / r_0^3). \end{aligned} \quad (72)$$

The relativistic debris are rapidly degraded by radiation at $r \approx r_0$. Once γ decreases significantly the efficiency of radiation declines and the non-radiating blast wave results (46)–(69) become applicable, with Y , γ_0 , and r_0 assuming the values appropriate to the debris at the transition between the efficiently radiating and weakly radiating regimes. In practice, while gamma-ray emission may occur in the strongly radiating regime, emission of lower energy photons will be predominantly from the weakly radiating regimes.

The exponential dependence of γ on r^3 (eq. [72]) implies that the transition between the two regimes occurs for $r \sim O(r_0)$, even though the degradation of Y and γ may be large. This degradation may be estimated by setting $\epsilon(r) = 1$ and adopting $r \approx r_0$. The surviving fireball energy Y' and Lorentz factor γ' may be approximated

$$\begin{aligned} Y' &\sim Y \epsilon_0^{-2/3}, \\ \gamma_0' &\sim \gamma_0 \epsilon_0^{-2/3}, \end{aligned} \quad (73)$$

where ϵ_0 is the nominal radiative efficiency obtained from equation (59) using the original fireball Y and γ_0 . The surviving Y' and γ_0' may then be used to describe the subsequent evolution and radiation of the blast wave.

REFERENCES

- Beia, J.-L., et al. 1987, *ApJS*, 64, 305
 Berger, P. N., et al. 1992, *Nature*, 359, 217
 Blandford, J. J. 1992, *Nature*, 355, 522
 Briggs, B. J., & Katz, J. I. 1992, *ApJ*, 399, 100
 Calzo, G., & Rees, M. J. 1978, *MNRAS*, 183, 359
 Chevalier, T. L. 1980, *Comm. Astrophys.*, 9, 13
 Chevalier, T. L., Kozlovsky, B. Z., Nussinov, S., & Ramaty, R. 1992, *ApJ*, 388, 164
 Dermer, C. D. 1992, *Phys. Rev. Lett.*, 68, 1799
 Dermer, D., Livio, M., Piran, T., & Schramm, D. S. 1989, *Nature*, 340, 126
 Epstein, R. I. 1985, *ApJ*, 297, 555
 Epstein, R. I., Ho, C., Klebasadel, R. W., & Laros, J. 1992, *Nature*, 357, 140
 Fishman, G. J. 1993, in *Compton Gamma-Ray Observatory*, Symp., ed. M. Friedlander, N. Gehrels, & D. J. Macomb (New York: AIP), 669
 Fishman, G. J. 1986, *ApJ*, 308, L47
 Fishman, G. J., & Leventhal, M. 1992, *Nature*, 357, 388
 Higdon, J. C., & Lingenfelter, R. E. 1990, *ARA&A*, 28, 401
 Lingenfelter, R. E., & Latter, R. 1962, *J. Geophys. Res.*, 67, 4635
 Lingenfelter, R. E., & Latter, R. 1965, *Phys. Rev.*, B137, 1369
 Lingenfelter, R. E., & Latter, R. 1982, *ApJ*, 260, 371
 —. 1992, *Ap&SS*, 197, 163
 —. 1993, in *Compton Gamma-Ray Observatory*, Symp. ed. M. Friedlander, N. Gehrels, & D. J. Macomb (New York: AIP), 1090
 Kouvelioutou, C., et al. 1993, *preprint* *ApJ*, 413, L101
 Kouvelioutou, C., Graziani, C., & Smith, I. A. 1993, *ApJ*, 413, L11
 Landau, L. D., & Lifshitz, E. M. 1959, *Fluid Mechanics* (Reading: Addison-Wesley)
 Liang, E. P., & Dermer, C. D. 1992, *Nature*, 359, 514
 Liang, E. P., & Liang, E. P. 1992, *ApJ*, 400, L59
 Lingenfelter, R. E., & Higdon, J. C. 1992, *Nature*, 356, 132
 Liang, E. P., Narayan, R., & Piran, T. 1993, *preprint* *ApJ*, in press
 Liang, E. P., & Paczyński, B. 1992a, *ApJ*, 388, L45
 —. 1992b, *ApJ*, 389, L13
 Liang, E. P., Fishman, G. J., Wilson, R. B., Paciesas, W. S., Pendleton, G. N., Brackbill, J. M., Brock, M. N., & Kouvelioutou, C. 1992, *Nature*, 355, 143
 Mészáros, P., & Rees, M. J. 1992, *ApJ*, 397, 570
 —. 1993, *ApJ*, 405, 278
 Narayan, R., Paczyński, B., & Piran, T. 1992, *ApJ*, 395, L83
 Paczyński, B. 1986, *ApJ*, 308, L43
 —. 1992a, *Nature*, 355, 521
 —. 1992b, *Acta Astron.*, 42, 1
 Piran, T. 1992, *ApJ*, 389, L45
 Rees, M. J., & Mészáros, P. 1992, *MNRAS*, 258, 41P
 Serman, M. A. 1975, *Ann. N.Y. Acad. Sci.*, 7th Texas Symp. on Relativistic Astrophysics, ed. P. G. Bergmann, E. J. Fenyves, & L. Motz, 262, 164
 Schmidt, W. K. H. 1978, 271, 525
 Smith, I. A., & Piran, T. 1990, *ApJ*, 365, L55
 Smith, I. A., & Epstein, R. I. 1993, *ApJ*, 410, 315
 Spitzer, L. 1962, *Physics of Fully Ionized Gases* (Interscience: New York)
 Tatarski, V. L. 1962, *Theory of Plasma Waves* (McGraw-Hill: New York)
 Tatarski, V. L., & Chibisov, G. 1975, *Soviet Astron.-AJ*, 19, 115

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OF POOR QUALITY

Wasserman, I. 1992, ApJ, 394, 565
 Weisskopf, M. C., Sutherland, P. G., Katz, J. I., & Canizares, C. R. 1978, ApJ,
 223, L17

Note added in proof:

The analysis of the flow shown in Figure 1 may readily be extended to the case $n_1 \neq n_4$. Then $e_2 = e_3 = 2\gamma_F(e_1 e_4)^{1/2} = 2\gamma_F m_a c^2 (n_1 n_4)^{1/2}$. The Lorentz factor describing the speed of the shocked matter in the local observer's frame $\gamma_{12} = (\gamma_F/2)^{1/2} (n_4/n_1)^{1/4}$. The internal energy of the shocked interstellar matter is described by the Lorentz factor (eq. [13]) $\gamma_2 = (\gamma_F/2)^{1/2} (n_4/n_1)^{1/4}$. The shock compression $n_2/n_1 = 4\gamma_{12}$. The equipartition magnetic field $B_2 \propto e_2^{1/2} \propto (n_1 n_4)^{1/4}$, so that the co-moving synchrotron frequency $\nu_2 \propto \gamma_2^2 B_2 \propto (n_4/n_1)^{3/4}$, and the observed frequency $\nu_{obs} \propto n_4/n_1$. Unshocked relativistic shells are very thin ($\Delta r/r \sim \gamma_F^{-2}$) so that $n_4 \gg n_1$ is expected; this substantially reduces the γ_F required to obtain the observed high energy gamma-rays.

The passage of the shock S2 through the debris shell leads to a reflected rarefaction which reduces both n_4 and γ_F . In the local observer's frame $v_{S2}/c \approx 1 - (2/\gamma_F)(n_1/n_4)^{1/2}$ (eq. [7]). If $n_4/n_1 \sim O(\gamma_F^2)$ then $1 - v_{S2}/c \sim O(\gamma_F^{-2})$ and S2 will not pass entirely through the shell until S1 has passed through a thickness $O(r)$ of interstellar cloud. In contrast, for $n_4 \sim n_1$ S2 traverses the entire debris shell and is reflected from its back surface as a rarefaction after S1 has traveled only $O(r/\gamma_F)$. Qualitatively similar remarks apply to debris shells with continuous distributions of γ_F .

T. Piran has independently obtained similar results.