YET ANOTHER MODEL OF GAMMA-RAY BURSTS

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ABSTRACT

Sari & Piran have demonstrated that the time structure of gamma-ray bursts (GRBs) must reflect the time structure of their energy release. A model that satisfies this condition uses the electromagnetic emission of energy by the magnetized rotating ring of dense matter left by neutron star coalescence; GRBs are essentially fast, high-field, differentially rotating pulsars. The energy densities are large enough for the power to appear as an outflowing equilibrium pair plasma, which produces the burst by baryon entrainment and subsequent internal shocks. In this paper the magnetic field and the characteristic timescale for its rearrangement—which determines the observed time structure of the burst—are estimated. There may be quasi-periodic oscillations at the rotational frequencies, which are predicted to range up to 5770 Hz (in a local frame). This model is one of a general class of electromagnetic accretion models that includes the Blandford and Lovelace model of active galactic nuclei and that can also be applied to black hole X-ray sources of stellar mass. The apparent efficiency of nonthermal particle acceleration is predicted to be 10%–50%, but higher values are possible if the underlying accretion flow is super-Eddington.

Subject headings: accretion: accretion disks — galaxies: active — galaxies: nuclei — gamma rays: bursts — gamma rays: theory — stars: neutron

1. INTRODUCTION

More than a hundred models of gamma-ray bursts (GRBs), including soft gamma repeaters (SGRs), have been published (Nemiroff 1994); yet there is, at present, no satisfactory and generally accepted model of GRBs. Even early data hinted at cosmological distances (Usov & Chibisov 1975; van den Bergh 1983). At these distances, the inferred energy release of \( \sim 10^{51} \text{ergs} \) in a region constrained by the shortest observed times of variation implies (Cavallo & Rees 1978) the creation of an equilibrium pair and photon plasma with durations in the range \( 10^{-3} \) to \( 10^{-5} \) s and nonthermal spectra.

The expected properties of radiation-pair fireballs (Goodman 1986) do not resemble the observed properties of GRBs. Fireballs are expected to produce a brief pulse (of less than 1 ms, if their energy is released in a region of neutron star dimensions) of blackbody radiation, which is very different from the properties of observed (Fishman et al. 1994) GRBs, which have a complex multi-peaked time structure with durations in the range \( 10^{-2} \) to \( 10^{3} \) s and nonthermal spectra.

For these reasons the hypotheses of cosmological distances and, by implication, of fireball models, were not generally accepted until data from BATSE made the geometrical case for them compelling (Meegan et al. 1992; Piran 1992). This left the problem of reconciling the observed properties of GRBs with the predicted properties of fireballs. Shemi & Piran (1990) had demonstrated that even a small contamination of the fireball with baryonic matter (and its associated nonannihilating electron excess) leads to the conversion of nearly all of the fireball energy to the kinetic energy of a relativistic shell of baryons. Rees & Mészáros (1992) then pointed out that this shell could form a collisionless external shock when it encountered surrounding diffuse matter and suggested that radiation by the shocked matter could resolve the discrepancy between the predicted and observed properties. Relativistic shocks plausibly produce a nonthermal spectrum, optically thin at low frequencies (thus avoiding a Wien slope) and with power-law extensions to higher frequencies. The complex temporal structure of GRBs was attributed to the known complex spatial structure of diffuse matter, such as the interstellar medium. This suggestion was widely accepted (with a sigh of relief) and led to a great deal of work on the properties of fireballs and relativistic shocks.

There are, however, a number of phenomenological problems with these external shock models of GRBs. The characteristic predicted radiation frequency scales roughly inversely with the burst duration (each depends on several parameters raised to a variety of model-dependent powers), in contradiction with the fact that all observed GRBs have spectra with characteristic energies \( \sim 300 \text{keV} \), despite durations ranging from \( \sim 0.01 \) to \( \sim 1000 \) s and possibly longer. There is no obvious anticorrelation of spectral hardness with duration, and in fact, nearly all observations of multi-MeV and GeV photons are from long GRBs (Katz & Canel 1996). In addition, for plausible values of the fireball Lorentz factor \( \sim 10^{3} \) that agree with the observed spectrum, the predicted GRB durations are much shorter than the observed durations of most GRBs. There are two aspects of one problem, that of how to explain the apparent independence of other observed quantities from the burst duration, which has a range of at least 5 orders of magnitude. These arguments suggest that GRB emission is a continuing, quasi-steady process, perhaps reflecting a quasi-steady configuration rather than the consequences of a single catastrophic event.

There are, in addition, more theoretical difficulties with external shock fireball models. The usual suggestion was that they are produced by the coalescence (Goodman 1986; Paczyński 1986; Goodman, Dar, & Nussinov 1987; Eichler et al. 1989) or the formation (Dar et al. 1992) of neutron stars. A fraction of the binding energy, it was argued, escapes as neutrinos and forms \( e^+e^- \) pairs by interactions between two neutrinos. These events release energy over several seconds with a smooth time profile (demonstrated empirically by the duration of neutrino emission from SN
1987A), as a consequence of the diffusion of neutrinos through matter dense enough to be opaque even to them.

Katz & Canel (1996) pointed out that such a model could not explain short GRBs (with durations of less than 2 s), and they proposed collisions between two neutron stars, from which the escape of the neutrinos would be accelerated by the expansion of the debris, as an alternate mechanism. This argument is also applicable to the fine temporal substructure of long GRBs, thus excluding processes dependent on neutrino diffusion from the explanation of almost any GRB.

The most popular version of this model involved the coalescence of two neutron stars, which has also been calculated (Davies, et al. 1994; Janka & Ruffert 1996) to be an inadequate source of neutrinos. This numerical result is easily understood and was predictable: coalescence is nearly adiabatic, because the subsonic velocity of convergence does not make shocks, and known dissipative processes, such as viscosity and internal friction, are weak. As a result, the cold, degenerate neutron star matter stays cold and degenerates and does not radiate many neutrinos.

Finally, Sari & Piran (1997a) demonstrated, given quite general assumptions, that external shocks can only produce smooth, single-peaked GRBs—quite unlike most observed GRBs. This conclusion is very general, because it is essentially geometrical; it does not depend on arguments concerning neutrino diffusion times or other microscopic physical processes. It is applicable to any single release of energy, however thin the relativistic shell and however complex the distribution of surrounding matter. It therefore excludes even the brief (~5 ms) neutrino emission of colliding neutron stars as the explanation of GRBs. They showed that a complex GRB pulse profile can only be the consequence of a complex history of energy release by the engine that powers the GRB. This excludes models in which the GRB energy is released in a single brief event, whether it be the coalescence, birth, or collision of compact objects. It implies that models involving a continuous or intermittent wind (Paczynski 1986) are more likely to be correct that those involving a single fireball (Goodman 1986).

This paper proposes a resolution of these difficulties. As in much previous work, I begin with coalescing neutron stars that release an enormous (~3 x 10^{53} ergs) amount of energy and are estimated (Narayan, Piran, & Shemi 1991; Phinney 1991) to occur at a rate consistent with the rate of GRBs. Such a coalescence results in a rapidly differentially rotating object that contains the mass of the two neutron stars. It is largely supported by internal pressure near its axis, but increasingly supported by angular momentum at larger distances from the axis. The central core may promptly collapse to become a black hole, or it may remain for an extended time as a rapidly differentially rotating neutron star, depending on the uncertain state of matter at densities of up to \( \rho \sim 10^{14} \, \text{g cm}^{-3} \). It is surrounded by a gaseous envelope, which is largely supported by angular momentum and may be qualitatively described as an accretional bulge or a thick accretion disk. For convenience, and because this term is universally used, it is here referred to as a disk, although it is unlikely to be strongly flattened. This configuration was described in the context of GRBs by Narayan, Paczynski, & Piran (1992), Usov (1992, 1994), Thompson (1994), and Mészáros & Rees (1997). It is necessary to explain how it can produce their observed properties.

I propose that this rotating object produces a GRB by the electrodynamic processes that turn rotational energy into particle acceleration in pulsars. The required magnetic energy is only a tiny fraction of the GRB energy and is much less than that required for explaining GRBs as the result of magnetic reconnection. How these processes may work for parameters appropriate to GRBs is described in § 2. In § 3 the origin of the magnetic field is discussed and its magnitude estimated. The duration, rapidity of variation, and temporal complexity of GRBs are discussed in § 4, and possible quasi-periodic oscillations are predicted. The model is compared to magnetic field reconnection in § 5, and radiation by internal shocks is briefly discussed § 6. Section 7 contains a summary and conclusions. GRBs may be considered to be members of a unified class of electro-dynamically accreting objects that includes active galactic nuclei (AGNs) and Galactic black hole X-ray sources such as Cyg X-1; these are discussed in a forthcoming paper (Katz 1997b).

2. ELECTRODYNAMICS

A rotating magnetized object is a source of energy. If it has a magnetic dipole moment that is misaligned with its rotation axis, it radiates a power that is given by the classical magnetic dipole radiation formula. Goldreich & Julian (1969) showed that even an aligned rotor produces a similar power (in a relativistic wind) if the surrounding space is filled with plasma. In the present case the object is differentially rotating, but that makes no essential difference; Blandford (1976) and Lovelace (1976) similarly applied the theory to differentially rotating accretion disks in order to explain AGNs. For a large-scale ordered magnetic field \( B \), a power

\[ P_{\text{rw}} = \frac{B^2 r^6 \Omega^4}{2c^3} \]  

flows outward in a relativistic wind, where \( \Omega \) is the angular velocity, and all parameters are roughly defined mass-weighted means over the disk. Only a fraction \( \sim r^2 \Omega / c \) of the magnetic field lines are open, so the power density on those field lines at the surface is

\[ S_{\text{rw}} = \frac{B^2}{8\pi} c \left( \frac{r^2 \Omega}{c} \right)^3. \]  

Near the surface of a neutron star, or near the last stable circular orbit of a black hole, \( r^2 \Omega / c \sim 0.5 \), so \( S_{\text{rw}} \) is roughly an order of magnitude less than \( c \) times the magnetostatic energy density. Quite apart from uncertainties in the physical parameters, a nonrelativistic treatment is only an approximation. For numerical evaluations, I will take \( (r^2 \Omega / c)^3 = 0.1 \).

The usual estimate of the power of a GRB at cosmological distances is \( P \sim 10^{51} \, \text{ergs s}^{-1} \), although this is known at best to order of magnitude, because their distances must be estimated from uncertain statistical arguments. Taking a radiating surface area of \( 10^{13} \, \text{cm}^2 \) (a neutron star or the inner disk around a black hole of a few \( M_\odot \)) implies \( S \sim 10^{38} \, \text{ergs cm}^{-2} \, \text{s}^{-1} \). If \( S = S_{\text{rw}} \), then \( B \sim 10^{15} \, \text{G} \). This is much larger than known neutron star magnetic fields, although Thompson (1994) discusses the possibility that some neutron stars may be born with such fields (such neutron stars would be nearly unobservable as...
radio pulsars because of their rapid spin-down). In § 3 the generation of magnetic fields of this magnitude as a consequence of neutron star coalescence is further discussed.

In conventional pulsar models, the magnetic field is accompanied (in an inertial frame) by an electric field \( B \partial E / c \) (Goldreich & Julian 1969) that is large enough to pull charge from the stellar surface or to produce \( e^2 \) pairs by a process of vacuum sparking (Sturrock 1971; Ruderman & Sutherland 1975). In the GRB problem, the electric fields are even larger than in pulsars (because of both the assumed larger \( B \) and the large Keplerian \( \Omega \) in the inner disk), so there is no difficulty in providing the pair plasma necessary for the aligned rotor to produce power.

The power density is extremely large; \( S_{\text{tw}} \approx 10^{38} \text{ ergs cm}^{-2} \text{ s}^{-1} \) corresponds to an effective temperature \( k_B T_e \approx 3 \text{ MeV} \). If \( k_B T_e > 20 \text{ keV} \), corresponding to the much lower equilibrium energy density \( E_{\text{eq}} \equiv a T_e^4 > 10^{19} \text{ ergs cm}^{-3} \) and radiated power density \( S_{\text{eq}} \equiv a S_{\text{eq}} T_e^4 > 10^{59} \text{ ergs cm}^{-2} \text{ s}^{-1} \), the energy density is sufficiently large for an opaque equilibrium pair plasma to form. (The temperature is found from the condition that the equilibrium pair density have a Thomson scattering mean free path \( \sim 1 \text{ km} \) but is very insensitive to this assumed length.) The laws of thermodynamics permit this energy and power density to assume the form of a smaller number of more energetic particles, but interactions among them and with the magnetic field will generally lead to rapid equilibration. This is particularly true at higher energy densities, at which the density of interacting particles is higher; the energy density required in GRBs is \( \sim 10^9 \) times the minimum required for an equilibrium pair plasma, so it is safe to assume its formation.

The opaque pair plasma takes the place of the dilute transparent fluid of energetic particles found in radio pulsars, for which the power density is much smaller (\( \sim 10^{27} \text{ ergs cm}^{-2} \text{ s}^{-1} \) for the Crab pulsar, and less for others). Isolated neutron stars with sufficiently strong fields and rapid rotation, such that their power density exceeds \( S_{\text{eq}} \) (such as a newly born neutron star with a field like that of the Crab pulsar but a spin period of less than 3 ms or hypothetical pulsars with much larger field), may resemble low-intensity GRBs of very long duration more than conventional radio pulsars.

The deposition of a super-Eddington power (for example, by viscous dissipation or neutrino transport) within the deep interior of the dense massive disk need not imply a strong mass outflow. Hence pair plasmas with low baryon loading and high Lorentz factors are possible. The Eddington limit only constrains the rate at which radiation can diffuse through matter. A super-Eddington diffusive photon luminosity does not produce an outflowing wind; it does not occur in a hydrostatic configuration. If a power greater than the Eddington limit (but less than the binding energy divided by the hydrostatic relaxation time) is injected into opaque matter, it will only produce a quasi-static structural relaxation (see Katz 1996 for a recent discussion).

The magnetic field inside the disk produces a stress that leads to an outward flow of angular momentum, inward flow of mass, and the release of gravitational energy, as in the usual theory of accretion disks. This process is usually described as dissipation by magnetic viscosity, and its detailed mechanism, as well as the process by which the energy released is thermalized, is a subject of controversy. The relevant component of magnetic stress (in a thin disk) is \( B_r B_{\phi}/4\pi \). Regardless of the detailed mechanism of field production, no other magnetic stress is available to transport angular momentum outward. A Newtonian estimate for the power released by magnetic viscosity is then

\[
P_{\text{visc}} \approx \frac{B_r B_{\phi}}{2} \Omega r^2 h,
\]

where \( h \) is the disk half-thickness. Note that if the viscous disk power is known, the relevant components of the magnetic stress tensor are directly determined; no relation between magnetic stress and fluid pressure (viscosity parameter \( z \)) need be assumed.

Comparing equations (1) and (3) and assuming that the \( B_r^2 \) on the surface of the disk used in equation (1) is comparable to \( B_r B_{\phi} \), leads to

\[
\frac{P_{\text{tw}}}{P_{\text{visc}}} \sim \frac{r}{h} \left( \frac{\Omega r}{c} \right)^3.
\]

The ratio \( r/h \) is variously estimated at 1–10; \( (\Omega r/c)^3 \approx 0.1 \) near a neutron star’s surface or a black hole’s last stable circular orbit, but declines \( \propto r^{-3/2} \) at larger radii. Hence near the surface (or last stable circular orbit), the electrodynamic efficiency

\[
\epsilon_r \equiv \frac{P_{\text{tw}}}{P_{\text{visc}}}
\]

is in the range 0.1–0.5, while it is small in the outer regions of a disk. The efficiency of production of nonthermal power (and ultimately radiation) does not depend on the magnitude of the magnetic field, although it does depend on the field geometry.

The energy available from the accretion of 0.1 \( M_\odot \) is \( \sim 10^{53} \text{ ergs} \), assuming maximally rotating Kerr geometry, so that at least \( \sim 10^{52} \text{ ergs} \) are available for GRB emission. This is sufficient to explain the energetics of GRBs at cosmological distances, as long as the efficiency of emission by the relativistic shocks is not very small. The portion of the accretional energy that does not appear in the relativistic wind is thermalized within the disk and either emitted as neutrinos or swept into the black hole (or onto the rotating neutron star) along with the accreted matter.

3. MAGNETIC FIELDS

Can magnetic fields of \( \sim 10^{15} \text{ G} \) be justified? They exceed by far the known magnetic fields of neutron stars. Assuming equipartition between magnetic and thermal or rotational kinetic energies is not justified; magnetic fields in stars and the interstellar medium are far below these equipartition values. Equipartition with turbulent energy densities is more sensible, for it is roughly correct for the Galactic disk and may not be altogether incorrect for stellar magnetic fields. However, the presence and magnitude of turbulent velocities are uncertain.

When two neutron stars coalesce, they quickly [in a time \( \sim (R^3/GM)^{1/2} \sim 10^{-5} \text{ s} \)] settle into an axisymmetric but differentially rotating configuration. The near adiabaticity of neutron star coalescence implies that thermal convection is weak. Because the flow is axisymmetric, the Cowling theorem (Shu 1992) establishes the fact that it cannot amplify a magnetic field by a dynamo process. It is possible that a weak dynamo may occur because of small deviations from exact axisymmetry, but these cannot be estimated.

Differential rotation generates an azimuthal field from a radial field, requiring only that the field be frozen into the
manner according to the equation
\[ \frac{dB_\phi}{dt} = B_r \frac{d\Omega}{d\ln r} \approx \frac{3}{2} B_r \Omega. \]  
(6)

This is not a dynamo, as the radial field is not regenerated and will eventually decay resistively.

Ropes of (nearly azimuthal) magnetic flux are buoyant and will rise to the surface of the disk on a timescale
\[ t_b \approx h \frac{t}{v_A} \sim \frac{h(4\pi \rho)^{1/2}}{B}, \]  
(7)

where \( v_A = B/(4\pi \rho)^{1/2} \) is the Alfvén speed, and \( \rho \) is a mean density. (Depending on the geometry, this may involve the Parker instability, but the timescale is determined by dimensional considerations alone.) A rough (statistically) steady state will be reached when the rate of field growth (eq. [6]) equals its rate of loss by buoyancy (eq. [7]). Using the likely dominant component \( B_\phi \) for \( B \) in equation (7) permits its magnitude to be estimated. The result is
\[ B_\phi \approx \frac{3}{2} B_r \Omega h(4\pi \rho)^{1/2} \sim 10^{15} B_{13}^{1/2} \Omega_{4}^{1/2} h_6^{1/2} \rho_{13}^{1/4} G, \]  
(8)

where \( B_{13} \equiv B_r/(10^{13} \text{ G}), \Omega_4 \equiv \Omega/(4 \times 10^4 \text{ s}^{-1}), h_6 \equiv h/(10^6 \text{ cm}), \) and \( \rho_{13} \equiv \rho/(10^{13} \text{ g cm}^{-3}). \)

A plausible origin of \( B_\phi \) is the precessional magnetic fields of the neutron stars. Values as large as \( \sim 10^{13} \text{ G} \) are observed for some radio pulsars, justifying the use of \( B \sim 10^{15} \text{ G} \) if the Parker instability or buoyant rise convert \( B_\phi \) to \( B \), with reasonable efficiency. The relative magnitudes of the three components \( B_r, B_\theta, \) and \( B_\phi \) are uncertain, thereby introducing similar uncertainties into estimates such as that of equation (4). Millisecond pulsars are observed with fields as small as \( \sim 10^9 \text{ G} \) (other neutron stars may have even smaller fields, but be unobservable for that reason), implying \( B \sim 10^{13} \text{ G} \).

4. TIMESCALES

4.1. Duration

The magnetic fields determine the rate of accretion and the rate of emission of GRB power. If \( B_r \) is comparable to the \( B_\phi \) estimated from equation (8),
\[ P_{rw} \sim 10^{51} B_{13}^{1/2} \Omega_4^{1/2} h_6^{1/2} \rho_{13}^{1/4} r_6^{-6} \text{ ergs s}^{-1}, \]  
(9)

where \( r_6 \equiv r/(1.5 \times 10^6 \text{ cm}). \) A total GRB energy \( E \sim 10^{51} \text{ ergs} \) implies a duration, assuming efficient conversion of relativistic wind energy, of
\[ t_2 \sim 1.5 B_{13}^{-1} \Omega_4^{-1} h_6^{-1} \rho_{13}^{1/2} r_6^{-6} \text{ s}. \]  
(10)

Short GRBs may be explained by the coalescence of neutron stars with magnetic fields roughly comparable to those observed for most radio pulsars. The shortest GRBs \( (t_2 \sim 10^{-2} \text{ s}) \) may require somewhat larger fields—this is not implausible, because pulsars with such large fields would spin down rapidly and have short observable lives—or somewhat different values of other parameters, or may instead be explained by smaller values of \( E \) and \( P_{rw} \), as is consistent with our poor quantitative understanding of GRB energetics.

GRBs with durations as long as \( \sim 10^4 \text{ s} \) (GRB 940217; Hurley et al. 1994) may be explained by the coalescence of millisecond pulsars with \( B_r \sim 10^9 \text{ G}, \) and even longer durations are possible (e.g., the \( \sim 10^3 \text{ s} \) required to explain the “Gang of Four” apparent repetitions of 1996 October 27–29 as a single event; Meegan et al. 1996; Connaughton et al. 1997). In the present model, long durations pose no intrinsic difficulty and need not be associated with unusually soft spectra, in contrast to external shock models, in which they imply low Lorentz factors and low radiative efficiency. Very long GRBs are faint, on average, in any model in which the total energy of a GRB is limited and are therefore difficult to detect. They may be more frequent than is apparent from intensity-selected (or rate of rise-selected) samples.

4.2. Substructure

In the present model the magnetic field rearranges itself on the timescale \( t_b \), which may be rewritten as
\[ t_b \sim (4\pi \rho)^{1/4} \left( \frac{2h}{3B_r \Omega} \right)^{1/2} \sim 0.01 h_6^{1/2} \rho_{13}^{1/4} B_{13}^{1/2} \Omega_4^{-1/2} \text{ s}. \]  
(11)

Equivalently, \( B_r \) may be eliminated in favor of the duration \( t_d \equiv E/P_{rw} \), yielding
\[ t_b \sim \left( \frac{t_d}{2E} \right)^{1/2} (4\pi \rho)^{1/4} h_6^{3/2} r_6^{1/2} \sim 0.04 \frac{(10^{51} \text{ ergs s}^{-1})^{1/2}}{P_{rw}} \rho_{13}^{1/2} h_6^{3/2} r_6 \Omega_4^{1/2} \text{ s}. \]  
(12)

This timescale characterizes the temporal substructure of GRBs. Their intensity is predicted not to vary much on timescales shorter than \( t_b \). The data appear to be consistent with this prediction.

4.3. Complexity

The complex time structure of GRBs is notorious. It led Stecker & Frost (1973) to point out their qualitative resemblance to the time structure of solar flares. Many different qualitative forms are seen. These include continuous irregular fluctuations; isolated peaks separated by longer intervals with no detected emission; single, simple peaks; peaks with numerous subpeaks; etc. It is beyond the capability of any theoretical model to predict this complex structure; it is a formidable task even to construct suitable measures for its statistical description.

In the present model, the field rearranges itself on the timescale \( t_b \), and there are \( t_d/t_b \sim 10^{2}–10^{3} \) rearrangements (the numerical value depending on poorly known parameters) in a GRB of 10 s duration. Note that \( t_d \) is the length of time during which the GRB radiates strongly, and it may be much less than the measured pulse length if there are peaks of intensity separated by periods of much lower or zero emission. This permits a great variety of complex time structures. Their details depend on unknown details of the disk’s magnetohydrodynamics. For comparison, consider the problem of explaining the solar magnetic cycle. This is incompletely understood, even with the aid of a great body of data. Its observed complexity argues for the plausibility of obtaining the variety of GRB time structure as the result of disk magnetohydrodynamics.

This, however, is only an argument for the plausibility of obtaining the observed GRB temporal phenomenology. It is not a demonstration that the observed heterogeneity and complexity can be obtained. Nevertheless, the situation is better than for external shock models, which, as Sari & Piran (1997a) showed, could not explain the data. Even
before their argument was made, external shock models depended on the hopeful wish that the temporal complexity of GRBs could be attributed to the (also poorly understood) spatial complexity of the interstellar medium.

The failure to demonstrate that the variety of observed pulse forms must occur might be used as an argument against the present model. If it is accepted that all classical GRBs are produced by a single kind of event, this heterogeneity could be used equally as an argument against any model, because we have never understood any physical process producing complex and heterogeneous pulse forms well enough to predict these forms. The alternative hypothesis that many different kinds of events produce GRBs is even more unlikely: apart from the observed homogeneity of GRBs in energy scale, spatial distribution, spectral properties, and even the fact of complex pulse structure, this hypothesis requires the construction of several satisfactory models, when it is difficult enough to find even one! All that can be asked is that a model permit temporal complexity and heterogeneity.

4.4. Quasi-periodic Oscillations

The magnetic field of the accretion disk will not, in general, be axisymmetric. If the deviation from axisymmetry is large, the resulting model resembles magnetic dipole emission (though of a relativistic wind rather than of an electromagnetic wave) more closely than an axisymmetric rotator. Equations (1) and (2) remain valid in either case. The particle flux and the observed radiation may therefore be modulated at the rotation frequency, which varies with \( r \) in the disk. This time dependence will be convolved with the time dependence of acceleration and radiation of the radiating particles and with geometric delays arising from a range of signal path lengths (Ruderman 1975; Katz 1994a), and may thereby be unobservable. However, the acceleration and radiation times may be short (Katz 1994a; Sari, Narayan, & Piran 1996; Sari & Piran 1997b), and the rotational modulation should be looked for in gamma-ray data of sufficiently high temporal resolution. It is also conceivable, although probably unlikely, that sufficient coherent emission occurs for the production of a detectable signal at radio frequencies, as in radio pulsars.

The expected signal would consist of quasi-periodic oscillations (QPOs). A high-frequency cutoff at the maximum rotational frequency of the disk is predicted. For a disk of test particles surrounding a maximally rotating Kerr black hole of 2.8 \( M_\odot \), the frequency (before cosmological redshift) of the last stable circular orbit (Shapiro & Teukolsky 1983) is 5770 Hz. The actual value of the mass and angular momentum parameter can, in principle, be determined from numerical calculations of neutron star coalescence and inferred from observations of their gravitational wave emission. If QPOs with an upper frequency cutoff below this value are observed, cosmological distances will be verified and the redshifts of individual GRBs measured. The maximal rotational frequency of a differentially rotating neutron star depends on its distribution of angular momentum and is significantly smaller. It might be calculable from simulations of neutron star coalescence, but the interpretation of the data would be more ambiguous.

The spectral power density will drop off toward lower frequencies, because both the energy release and the relativistic wind efficiency (eq. [5]) decline with \( r \) (a naive estimate is that the power density varies \( \propto \Omega^2 \)). Disks resulting from neutron star coalescence may have an abrupt outer edge, because viscosity has not had time to extend them to large radii, or because exothermic nuclear reactions in expanded neutronized material expel their lower density regions, so that there may also be an abrupt low-frequency cutoff. The distribution of power will give some indication of the radial structure and distribution of magnetic fields in the disk.

5. Magnetic Reconnection?

It was natural (Stecker & Frost 1973; Katz 1982; Narayan et al. 1992; Katz 1994a) to consider magnetic field reconnection to be the explanation of the complex time structure of GRBs. This is unlikely, as it would require that the entire GRB energy (divided by an efficiency less than 1) pass through the form of magnetic energy. This would require a strong dynamo, but we expect any dynamo to be weak (if present at all), because the flow is axisymmetric, at least to a first approximation.

The mechanism proposed in this paper has an advantage (in comparison with magnetic reconnection) in that the same magnetic flux is a source of emitted power for an indefinite time. Even a small (compared to the accretion energy) magnetic energy can be the conduit through which the entire GRB energy flows (because a GRB is very long, compared to \( r/c \)). In contrast, in a magnetic annihilation model, the magnetic energy must be regenerated rapidly by a dynamo, and all the radiated energy must have at one time assumed the form of magnetostatic energy.

The relativistic wind has a power (eq. [2]) that is the magnetic energy multiplied by a speed \( \propto (\Omega^3/r^2) \). In order to be equally powerful, magnetic annihilation would have to act on a surface covering most of the inner accretion disk throughout the duration of a GRB with a reconnection speed \( v_r \sim c/10 \). This is surely optimistic. In general, \( v_r \) must be less than \( v_A \), which is small in dense matter (§ 3), although it may be relativistic where the wind is generated.

Any process that produces the observed power density produces an equilibrium pair plasma, as discussed in § 2. The electrical conductivity of such a plasma is

\[
\sigma \sim \frac{m_e c^3}{e^2} \sim 10^{23} \text{ s}^{-1},
\]

as may be found by dimensional analysis or by estimating the number of charge carriers and their scatterers (charged particles and photons). The conductivity is approximately independent of temperature for \( k_B T > m_e c^2 \), because the densities of carriers and of scatterers vary each \( \propto T^3 \). The relativistic temperature makes counterstreaming plasma instabilities unlikely, and their high density makes the relative drift velocity of electrons and positrons small.

In a reconnecting current sheet of thickness \( l \), the time \( l/v_f \) required for the flow to regenerate the magnetic flux may be equated with the resistive dissipation time \( \sigma l^2/c^2 \). The result is

\[
l \sim \frac{c^2}{v_f} \sim \frac{c}{v_f} \frac{e^2}{m_e c^2} \sim \frac{c}{v_f} r_e \sim \frac{e^2}{m_e c^2} \sim \frac{e^2}{m_e c^2} h^3 \sim 10^{31} \text{ cm}^{-3}.
\]

where \( r_e \equiv e^2/m_e c^2 = 2.82 \times 10^{-13} \text{ cm} \) is the classical electron radius. The electron density at \( k_B T \sim m_e c^2 \) is \( n_e \sim m_e^2 c^3/h \sim 10^{31} \text{ cm}^{-3} \). The discreteness of the charge carriers limits \( l \) as \( l > n_e^{-1/3} \sim h/|m_e c| \sim 4 \times 10^{-14} \text{ cm} \). This is
consistent with equation (14) only if
\[ \frac{v_f}{c} < \frac{e^2}{hc} \approx \frac{1}{137}. \] (15)

This is a general limit on the speed of magnetic reconnection in a relativistic pair plasma. In the inner region of a disk around a black hole or a neutron star, magnetic reconnection is, at best, an order of magnitude less powerful than the relativistic wind flowing on open field lines. Magnetic flux is probably destroyed by accretion onto the black hole or by advection to infinity in the wind, rather than by reconnection.

6. RADIATION

The Lorentz factor $\gamma$ of a relativistic wind is determined by the energy per baryon in the acceleration region (Shemi & Piran 1990),
\[ \gamma \approx \frac{dT^4}{\rho}, \] (16)
where $T$ is the blackbody temperature, and $\rho$ is the baryonic density. Using equation (2) and $S_{rw} \sim acT^4$, this may be restated as
\[ \gamma \sim \frac{B^2}{8\pi \rho} \left( \frac{r\Omega}{c} \right)^3. \] (17)
Unfortunately, this expression cannot be evaluated numerically, because we know neither the density $\rho$ of injected baryons nor the variation of $B$ with space and time. Instead, values of $\gamma$ can only be estimated by asking what they must be in order to produce the observed properties of GRBs. It is probably fair to make the qualitative statement that values of $\gamma$ are likely to vary spatially over the surface of the disk and temporally on the scale $t_b$.

6.1. Where Do Internal Shocks Occur?

Assume that most of the energy of outflow is contained in a range of Lorentz factors from $\gamma_2$ to $\gamma_1$, where by convention $\gamma_1 > \gamma_2$. If $\gamma_1 \approx \gamma_2$, the fraction of the energy that is available for radiation by internal shocks is small, while if $\gamma_1 \gg \gamma_2$, the kinematics begin to resemble those of external shocks, with the important difference that repeated episodes of flow at $\gamma_2$ permit repeated subshocks of radiation. The result is a very complex relativistic outflow.

Streams from different regions of the disk surface, separated by a distance $\sim r$, will overlap at a distance $\sim yr \sim 10^5 \gamma_3$ cm from the source, where $\gamma_3 \equiv \gamma/10^3$. This is inside the photosphere, the radius of which, for a thin shell of energy $E$, is (Shemi & Piran 1990)
\[ r_{ph} \approx \left( \frac{E}{4\pi \gamma c^2} \right)^{1/2} \sim 5 \times 10^{12} \left( \frac{E_{51}}{\gamma_3} \right)^{1/2} \text{cm}, \] (18)
where $E_{51} \equiv E/(10^{51}$ ergs), and $\kappa \approx 0.34 \text{ cm}^2 \text{g}^{-1}$ is the Thomson opacity. Note that the condition $yr < r_{ph}$ is equivalent to the condition that radiation not escape before it has given its energy to accelerating the baryons, and it must always hold in fireball models. This condition also sets an upper bound on $\gamma$, typically $\sim 10^4$.

If fluid elements with different values of $\gamma$ encounter each other in the optically thick region inside the photosphere, they may produce a shock—the sound speed never exceeds $c/3^{1/2}$, so the values of $\gamma$ need only differ by a factor of $[(3^{1/2} + 1)/(3^{1/2} - 1)]^{1/2} \approx 1.93$ in the limit in which all Lorentz factors are $\gg 1$, in order for the relative velocity to be supersonic—and convert some of their energy to radiation. This radiation, like that of the original fireball, gives its energy to the further acceleration of the baryons. Matter within a thickness $\sim r_{ph}/c^2$ has thus interacted dissipatively and likely shares a single value of $\gamma$ (within a region of opening angle $\sim 1/\gamma$, which is all that a single observer can see), including both fluid from different parts of the disk and fluid accelerated at times differing by less than $\sim r_{ph}/(\gamma c)$. There is therefore a minimum timescale $r_{ph}/(\gamma^2 c) \sim 2 \times 10^{-4} \left( \frac{E_{51}}{\gamma_3} \right)^{1/2} \approx 5 \times 10^5 (E_{51}/\gamma_3)^{1/3} \text{cm}, \] (19)

where $\rho_{-24} \equiv \rho_{\text{d}}/(10^{-24} \text{ g cm}^{-3})$, and $\rho_{\text{d}}$ is the ambient density, at which most of the kinetic energy of the relativistic shell would be dissipated in an external shock. Two pictures may be considered, that of a quasi-steady “Bunsen burner” and that of discrete internal shock interactions.

6.1.1. Bunsen Burners

In a “Bunsen burner,” the wind emerging from the photosphere has not equalized its Lorentz factor across an angular width $1/\gamma$. The relative kinetic energy of the various substreams making up the wind is dissipated in collisionless shocks at $r \sim r_{ph}$, and the observed radiation is produced in this region. Some of the observed radiation may even have diffused outward from smaller radii, because the opacity is almost entirely electron scattering. The characteristic time of variation is
\[ t_v \sim \frac{r_{ph}}{c \gamma^2} \sim 2 \times 10^{-4} \left( \frac{E_{51}^{1/2}}{\gamma_3^{1/2}} \right) \text{s}. \] (20)

This is shorter than any observed timescale of variation in a GRB and also shorter than the expected timescale $t_b$ of variation of the source magnetic configuration. Hence the power released in the internal shock region is expected to be slowly modulated by the changing magnetic source, much as a Bunsen burner produces, at the mouth of a mixing tube, a flame with a short internal combustion and flow time, slowly modulated by a gas valve and an air inlet at the base of the tube. The short timescales (combustion and flow for the Bunsen burner, $t_v$ for the GRB) are not directly observed.

6.1.2. Discrete Internal Shocks

In the discrete internal shock picture, interactions at $r < r_{ph}$ homogenize the Lorentz factor across an angular width $1/\gamma$. The Lorentz factor then varies only on a timescale (in a local observer’s frame) $\sim t_b$, as the source conditions change. Internal shocks occur at a characteristic radius
\[ r_{int} \sim t_b \gamma^2 \sim 10^{15} \left( \frac{t_b}{0.03 \text{s}} \right) \gamma^2 \text{cm}. \] (21)
The inequalities
\[ \gamma \rho \ll r_{\mathrm{ph}} \ll \int \ll r_{\mathrm{ext}} \] (22)
hold over wide and plausible ranges of parameters. Internal shocks in optically thin regions will occur for \( \gamma_{\mathrm{tr}} < \gamma < \gamma_{\mathrm{ext}} \), where the condition \( r_{\mathrm{ph}} \sim \int \) defines
\[ \gamma_{\mathrm{tr}} \equiv \left( \frac{E}{4\pi c^2} \right) \frac{1}{\gamma_1} \left( \frac{1}{\gamma_1 c} \right)^{2/5} \approx 130 E_3^{1/5} \left( \frac{0.03}{t_b} \right)^{2/5}. \] (23)

This sets a lower bound on plausible values of \( \gamma \) and is numerically similar to bounds set by considering \( \gamma \gamma \) pair production (Fenimore, Epstein, & Ho 1993; Woods & Loeb 1995; Piran 1996).

The condition \( r_{\mathrm{int}} \sim \int \) defines
\[ \gamma_{\mathrm{ext}} \equiv \left( \frac{3E}{8\pi \rho_c c^2} \right) \left( \frac{1}{t_5 c} \right)^{3/8} \approx 1900 \left( \frac{E_{51}}{\rho_{-24}} \right)^{1/8} \left( \frac{0.03}{t_b} \right)^{3/8}. \] (24)

From equation (24) it is seen that the observed timescale \( r_{\mathrm{ext}}/(c\gamma^2) \) of variation for external shocks is always \( \leq t_b \), so \( t_b \) is the shortest timescale of observable variation, just like for internal shocks. This is why (see § 1) a GRB produced by neutrinos diffusing through matter of neutron star density cannot have a timescale of variation of less than the diffusion time of several seconds, however high \( \gamma \) may be.

6.2. Efficiency

Cohen et al. (1997) have argued, on the basis of correlations between the spectra and time structure of GRBs, that electron radiation times are typically in the range of tenths of seconds to seconds and frequently determine the observed widths of subpulses. This is not consistent with the “Bunsen burner” model, because there the radiation time must be \( \sim t_e \) (submillisecond) in order to obtain a reasonable efficiency. It is consistent with discrete internal shocks in which each subpulse reflects the interaction of two regions of differing Lorentz factor. In that case, the interaction radius is given by equation (21), where \( t_s \) is taken to be the rise time or peak width of a subpulse; the longer decay time is then the electron radiation time.

6.2.1. Kinematic Bounds

The efficiency with which the kinetic energy of a relativistic outflow may be radiated by internal shocks is subject to a kinematic bound. A collision between a proper mass \( m_1 \) moving with Lorentz factor \( \gamma_1 \) and a proper mass \( m_2 \) moving with a Lorentz factor \( \gamma_2 \) can radiate at most a fraction \( \epsilon_k < 1 \) of their total energy. For parallel motion in the limit \( \gamma_1, \gamma_2 \gg 1 \), use of the relativistic invariant \( E^2 - p^2 c^2 \) leads to the limiting kinematic efficiency
\[ \epsilon_k = 1 - \left[ 1 + \frac{m_1 m_2}{(m_1 + m_2)^2} \left( \frac{\gamma_1 + \gamma_2}{\gamma_1 \gamma_2} \right) - 2 \right]^{1/2}. \] (25)

For the case of equal proper masses \( m_1 = m_2 \), an efficiency \( \epsilon_k > 1/2 \) is obtained for \( \gamma_1/\gamma_2 > 13.93 \), where, by convention, \( \gamma_1 \geq \gamma_2 \). This case is applicable to a flow of fluctuating power but steady baryonic entrainment. For the case of equal energies \( \gamma_1 m_1 = \gamma_2 m_2 \), we find \( \epsilon_k < 1 - 2^{-1/2} \approx 0.293 \) for all \( \gamma_1/\gamma_2 \). This latter case is applicable to a flow of steady power but variable baryonic entrainment. It is not possible to say which limiting case, if either, occurs. For \( \gamma_1/\gamma_2 = 1 + \beta \), with \( \beta \ll 1 \), we can approximate \( \epsilon_k \approx m_1 m_2 \beta^2 /[2(m_1 + m_2)^2] \leq \beta^2/8 \). It may be fair to assume an efficiency of tens of percent, but not of much more than 50%.

6.2.2. The Shocked Region

Because the range—and the mean value—of the Lorentz factor in the flow are poorly known, here I make only crude estimates, following Sari & Piran (1997b), rather than attempt a quantitative analysis of the shock structure. The proper density in the shocked region may be estimated as
\[ n_1 \sim \frac{E}{4\gamma_2^2 c^4 t_b t_d m_p}. \] (26)

Assuming this proper density in both the high- and low-Lorentz factor flows, the results of Katz (1994a) may be adapted to yield the conditions in the shocked matter. The proper energy density
\[ e_2 \sim n_1 \gamma_1^2 m_p c^2, \] (27)

and the magnetic field
\[ B \sim \left( \frac{4\pi c^2 E}{\gamma_2^2 c^4 t_b t_d} \right)^{1/2}, \] (28)

where the magnetic equipartition efficiency \( \zeta \equiv B^2/(8\pi e_2) \) and \( \gamma_{\mathrm{rel}} \equiv \gamma_1/(2\gamma_2) \). The observed synchrotron radiation frequency, assuming electron-proton equipartition (which is a natural consequence of the electrostatic potential required to decelerate the protons also acting on the electrons; Usov & Smolsky 1997), is
\[ \omega_{\mathrm{obs}} \sim \frac{1}{8} \left( \frac{m_p}{m_e} \right)^2 \left( \frac{\gamma_{\mathrm{rel}}}{\gamma_2} \right)^2 \frac{e}{m_e c} \left( \frac{8\pi E_3}{c^2 t_b t_d} \right)^{1/2} \approx 1.4 \times 10^{18} \left[ \frac{\gamma_{\mathrm{rel}}^2}{\gamma_3^2} \left( \frac{1}{\gamma_3} \right) \frac{1}{t_b} \left( \frac{0.03}{t_d} \right)^{1/2} \left( \frac{30 \, \mathrm{s}}{t_b} \right)^{1/2} \right] \mathrm{s}^{-1}, \] (29)

where, now, \( \gamma_3 \equiv \gamma_2/10^3 \). The photon energy \( \hbar \omega_{\mathrm{obs}} \) is approximately 300 KeV, which is typical of GRBs, if the factor in square brackets is \( \approx 330 \); if the other dimensionless factors are of order unity, this is obtained for \( \gamma_{\mathrm{rel}} \approx 0.02\gamma_2 \).

These estimates depend on several uncertain assumptions but show that radiation of the observed frequencies may be produced. The high-energy power-law “tail” of the photon spectrum may be attributed to a high-energy power-law “tail” on the electron distribution function, which is unsurprising, because plasma turbulence must be assumed, in any case, to obtain the necessary collisionless shock.

6.3. External Shocks?

GRBs with only a single pulse may reflect a central energy source that does not have several maxima and minima in output Lorentz factor. An alternative and perhaps more plausible explanation is that in these bursts the Lorentz factor is high enough so that \( r_{\mathrm{int}} > r_{\mathrm{ext}} \). In a case such as this, internal shocks do not occur; essentially all of the interaction is with the external medium, which leads (Sari & Piran 1997a) to a single smooth peak. The interaction radius \( r_{\mathrm{ext}} \) is larger than that obtained with internal shock models, but because of the required larger Lorentz factor, the observed pulse width is typically smaller; \( \gamma < \gamma_{\mathrm{ext}} \).
implies that
\[ t_{\text{obs}} \sim \frac{r_{\text{ext}}}{c^2} < t_b \quad ; \]
just as for internal shocks, the observed time structure reflects the time structure of the energy source.

It is also interesting that, in many GRBs, the intensity drops to zero (to within the accuracy of measurement) between discrete subpulses. If the subpulse structure is the result of internal shocks, it might be thought that the once-shocked matter should later undergo an external shock at larger radii, which would produce a more slowly varying but nonzero intensity between subpulses. The fact that, often, no such flux is observed, even though the radiative efficiency of internal shocks is expected to be less than 50%, may perhaps be attributed to the lower frequencies and lower efficiency of radiation at the lower energy densities and magnetic fields characteristic of external shocks at larger radii.

7. DISCUSSION

The demonstration by Sari & Piran (1997a) that the complex time structure of GRBs must be intrinsic to their energy source has forced the rejection of models involving a fireball produced by a single catastrophic event. Models in which energy is extracted electromagnetically have been considered previously (Usov 1992, 1994; Thompson 1994; Mészáros & Rees 1997), but they suffer from problems resulting from the fact that the known exemplars of this process, the radio pulsars, produce power at a steady rate and have much smaller magnetic fields than are required. Here I have outlined a model that may solve these problems. The magnitude of the required field and its broadband range are explained naturally, as is the irregular and complex time dependence of the emitted power. In accord with much previous work, nonthermal radiation is produced in this model by collisionless shocks internal to the fluctuating and heterogeneous relativistic wind.

The model predicts a minimum timescale of variations and suggests that QPOs of predictable frequency may be observable. Observation of predicted QPOs would be strong evidence in favor of the model and would (given the assumed Kerr black holes) immediately reveal the redshift. Alternatively, if an identification with a galaxy of measurable redshift were made, a known function of the mass and angular momentum of the black hole would be determined, constraining theories of neutron star coalescence.

The development of the GRB model suggests that the latter result from the same process—disk accretion onto a black hole—as that which is believed to power AGNs and Galactic black hole X-ray sources. The parameter regimes are very different, but these diverse objects may all be the consequence of electrodynamic energy extraction from accretion disks. This leads to a unified model that, by emphasizing their similarities, may help to explain all these phenomena. The fact that some AGNs produce the majority of their power in gamma rays of GeV–TeV energy is explained qualitatively.

GRBs have a bimodal distribution of durations (Kouveliotou et al. 1993), and long and short GRBs appear to have different spatial distributions (Katz & Canel 1996) and may represent different populations of objects. The distribution of magnetic fields of neutron stars appears also to be bimodal (although there are strong observational selection effects that are compensated for with difficulty), which may offer a natural explanation. One class of pulsars and rotationally modulated accreting neutron stars has longer periods and \( B \sim 10^{12} \quad \text{G} \); these may be the origin of short GRBs. Lower field \( (B \sim 10^9 \quad \text{G}) \) neutron stars form a second class and are observed as millisecond pulsars and as the presumed low-field neutron stars in unipolar low-mass X-ray binaries (whose rapid rotation may be inferred from their QPOs); these may be the origin of long GRBs. High-field neutron stars are produced with large impulses (as inferred from the space velocities of slow radio pulsars and Her X-1), and low-field neutron stars (at least millisecond pulsars and X-ray sources found in globular clusters) are produced with low impulses (Katz 1975). This implies that if sufficiently accurate coordinates could be obtained (Katz 1997a), short GRBs would be found to be further, on average, from their places of birth (plausibly the disks of their host galaxies) than would long GRBs.

The coalescence of a neutron star and a black hole may produce a different event than the coalescence of two neutron stars, but it is not obvious what the qualitative observable differences would be. Similarly, the “silent” collapse of a rapidly rotating degenerate dwarf may produce a similar configuration to that of neutron star coalescence, though half of the mass, but known degenerate dwarf spins are too slow.

The assumption of magnetic equipartition in the radiation region leads to a field of the same magnitude as that which would follow from the assumption that the outgoing power is carried electromagnetically, as in a radiation field or relativistic magnetohydrodynamic wind (as in pulsars). One consequence of this is that, if synchrotron radiation is efficient, the power of synchro-Compton radiation may approach, but not exceed, that of synchrotron radiation. Another consequence is that it is possible to conceive of GRB models in which there are no shocks and perhaps no baryons, but in which the energy is carried electromagnetically, and particles are accelerated and radiate at great distances from the central object. Such winds must contain some particles, because kinematic constraints prevent pair production by a vacuum electromagnetic wave. Interaction of electromagnetic energy with these particles is analogous to an internal shock with \( \gamma_1 \gg \gamma_2 \). It is unclear how particle acceleration could be arranged at large radii, from which radiation may escape, and avoided closer to the source, where it would produce an equilibrium pair plasma, as discussed in this paper.

After the submission of the original version of this paper, GRB 970228 was discovered to be a source of X-ray and visible emission, continuing over a period of days. This supports the suggestion (§ 4.1) that some GRBs may have such long durations, although the direct evidence so far is limited to the visible and X-ray bands. It is also consistent with the suggestion (§ 6.3) that external shocks may produce radiation at larger radii with lower frequencies and longer durations than the gamma-ray emission. This was predicted and a rough model developed by Katz (1994a, 1994b).

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