A GAMMA-RAY BURST MODEL WITH SMALL BARYON CONTAMINATION

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ABSTRACT

We present a scenario ("supranova") for the formation of gamma-ray bursts (GRBs) occurring when a supramassive neutron star (SMNS) loses so much angular momentum that centrifugal support against self-gravity becomes impossible, and the star implodes to a black hole. This may be the most baryon-clean environment proposed so far, because the SN explosion in which the SMNS forms sweeps the medium surrounding the remnant, and the quickly spinning remnant loses energy through magnetic dipole radiation at a rate exceeding its Eddington luminosity by some 4 orders of magnitude. The implosion is adiabatic, because neutrinos have short mean free paths, and silent, given the prompt collapse of the polar caps. However, a mass of $M_t \approx 0.1\ M_\odot$ in the equatorial belt can easily reach centrifugal equilibrium. The mechanism of energy extraction is via the conversion of the Poynting flux (due to the large-scale magnetic field locked into the minitorus) into a magnetized relativistic wind. Occasionally, this model will produce quickly decaying or nondetectable afterglows.

Subject headings: black hole physics — gamma rays: bursts — instabilities — relativity — stars: neutron

1. INTRODUCTION

A successful model for the progenitors of gamma-ray bursts (GRBs) must make contact with the fireball model (Rees & Mészáros 1992), which so successfully predicted the discovery of afterglows in the radio (Paczyński & Rhoads 1993), optical (Mészáros & Rees 1997a; Katz 1994), and X-ray bands (Vietri 1997), and explained radio flares (Goodman 1997). While some features of afterglows remain at present unexplained, there can be little doubt that the overall energetics, timescales, power-law behavior, and the appearance of longer wavelength radiation at later times are all nicely accounted for.

The crucial point of the fireball model is that it channels the large explosion energy ($>10^{53}$ ergs, barring beaming) in a locale with small baryon contamination ($\approx 10^{-4}\ M_\odot$; Rees & Mészáros 1992). From this point of view, hypernova explosions (Woosley 1993; Paczyński 1998) are problematic: massive stars are well known to have powerful winds that may carry away upward of $10\ M_\odot$, thus baryon-contaminating the environment in which the GRB is eventually to go off. Obviously, a very thin beam may be a way out of this conundrum, but this begins to look unlikely, given the currently available observational evidence (Grindlay 1998). In addition, scenarios involving accretion-induced collapse of white dwarfs (Usov 1992; Blackman, Yi, & Field 1996; Yi & Blackman 1997) are troublesome, both because of the large quantities of baryons lying about and because of the long timescales required to accrete sizeable amounts of mass.

Given the considerable specific angular momenta involved, one plausibly baryon-clean model is the merger of a NS-NS or NS–black hole binary (Narayan, Paczyński, & Piran 1992), which is expected to take place outside the environment where the binary formed (Bloom, Sigurdsson, & Pols 1998). However, at least some bursts (though not all; see Wijers & Galama 1998) seem to arise within dense environments, suggesting star-forming regions; such is the case of GRB 970111 (Feroci et al. 1998), GRB 970828 (Murakami et al. 1997; Groot et al. 1998a), and GRB 980326 (Groot et al. 1998b). Since it is possible that

GRBs arise from several distinct sources (Mészáros, Rees, & Wijers 1998b), we propose here a different model, in which the initial explosion is triggered by the gravitational implosion to a black hole of a supramassive neutron star; we call this model a "supranova."

In the following discussion, we first discuss some properties of equilibrium supramassive neutron stars, then the implosion to a back hole and the extraction of the energy that leads to the GRB. Reasons why this model remains baryon-clean are discussed in § 3, and in the last section we discuss some observational consequences.

2. ON SUPRAMASSIVE NEUTRON STARS AND THEIR IMPLOSION

All known equations of state for matter at nuclear densities allow the construction of equilibrium models of supramassive neutron stars (SMNSs), i.e., rotating sequences with constant baryon number but different angular momenta; no member of this sequence may be nonrotating, because they have masses at infinity larger than the largest mass for static models (for recent work on this topic, see Cook, Shapiro, & Teukolsky 1994; Salgado et al. 1994). Equilibrium sequences for models with baryon numbers exceeding a critical $\mathscr{B}_c$, above which no static equilibria are possible, start out exactly at break-up angular speed; depending on the equation of state, these models have masses in the range $M \approx 2$–$3.5\ M_\odot$, equatorial radii $R_{eq} \approx 11$–18 km, and angular velocities $\omega \approx 8000$–$12,000\ \mathrm{s}^{-1}$. Although fast-rotating, the total angular momentum of these models is not large; measured in units of the critical angular momentum $J_c = GM^2/c$ of a black hole of the same mass at infinity, they have $j \approx 0.6$–0.78. An important feature of these models is that as they lose angular momentum through most of the models' parameter space, they speed up; they contract so as to reduce their moment of inertia, and increase their angular velocity. For this reason, angular momentum loss through, for instance, magnetic dipole radiation actually speeds up. As the losses mount, the models reach a point, at about half of the initial angular momentum, where they become secularly unstable to axisymmetric perturbations. The evolution of these models beyond this point has not been investigated as yet, but the eventual fate of the star is sealed; since it cannot connect to a static, stable configuration, it will implode to a black hole.

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The above results hold for every equation of state, so that we regard them as rather robust. Currently, there seems to be no obvious theoretical reason why these stars should not form in the collapse of massive stars. A recent claim (Andersson 1998) that a new type of r-modes leads to such a copious emission of gravitational waves that any fast-spinning neutron star (NS) will slow down within a year, would seem to imply that SMNSs, which speed up as they lose angular momentum, should collapse to a black hole faster than expected on the basis of pure magnetic dipole radiation (eq. [1], derived below). However, these computations only concern idealized NS models, and neglect proper, general relativistic treatment of viscosity, convection, and nonlinear effects, all of which may slow down or even damp this instability; they will be neglected henceforth.

An approximate estimate of the time it takes for this configuration to collapse is given by assuming that the major source of angular momentum loss is through the usual magnetic dipole radiation. The star's rotational energy is $E_{rot} = J\omega/2$, and since ω is roughly constant as the star loses angular momentum passing through different stages of equilibrium (Salgado et al. 1994; Cook et al. 1994), $\dot{E}_{rot} \approx \omega\dot{J}/2$; equating this to the classical Pacini (1967) formula, one gets $\dot{J}$ and the spin-down time to halve J, as required by the onset of the instability,

$$t_{sd} \equiv \frac{J}{\dot{J}} = 10 \text{ yr } \frac{j}{0.6}\left(\frac{M}{3\ M_\odot}\right)^2\left(\frac{15 \text{ km}}{R_{eq}}\right)^6$$
$$\times \left(\frac{10^4 \text{ s}^{-1}}{\omega}\right)^4\left(\frac{10^{12} \text{ G}}{B}\right)^2 \ . \quad (1)$$

Since this is much longer than the viscous timescale inside the NS, on which the configuration evolves after it has become secularly unstable toward the final collapse, t_{sd} is a good estimate of the timescale between the first explosion (a normal supernova where a SMNS is formed) and the implosion to a black hole.

As the star loses angular momentum, the centrifugal force weakens until support against self-gravity becomes impossible, and the star implodes. The implosion is likely to be strongly nonhomologous, the central densest regions forming a black hole first, with the outermost regions collapsing later. To a first approximation, the implosion will be "silent," with most matter being swallowed down the hole promptly, since the total angular momentum of the collapsing NS is subcritical ($j \approx 0.6$–0.78). In particular, it is likely that all mass lying close to the spin axis ends up in the hole, a point already made by Mészáros & Rees (1997b) and Paczyński (1998). Furthermore, the neutrinos' mean free path inside the NS is expected to be small, so the collapse will most likely be adiabatic. Further baryon contamination must be small; for instance, it is difficult to think of any outwardly propagating shock capable of ejecting baryons, because of the lack of a hard surface or energetic phenomena capable of releasing considerable counterpressures. In addition, for reasonable magnetic fields ($B < 10^{15}$ G), any magnetic-related phenomena using all the energy in the magnetic field could push at most $\approx 10^{-5}\ M_\odot$ in centrifugal equilibrium, all most likely in the equatorial plane.

It is likely that this implosion is accompanied by the release of a copious amount of energy in gravitational waves; in fact, the known axisymmetric instability of the equilibrium sequence is just secular (Friedman, Ipser, & Sorkin 1988). It is possible that this leads to a nearby, distinct equilibrium model that becomes dynamically unstable as it loses angular momentum; the most likely source of instability is to a bar formation mode ($m = 2$), which leads to a release of gravitational waves (Chandrasekhar 1970).

Not all mass will be immediately accreted, however; the outermost layers, in fact, have centrifugal accelerations only $\approx 10\%$ below the local gravitational attraction. Thus, even if the pressure gradient were suddenly removed, these layers would only contract a bit (barring large energy losses due to shocks caused by the converging flow), and halt in centrifugal equilibrium at a radius $\approx 10\%$ smaller than the initial one. We cannot compute M_l at this stage, so we give only a suggestive argument showing that, possibly, $M_l \approx 0.1\ M_\odot$. Stellar dynamical simulations of self-gravitating galaxies in purely rotational equilibrium show that these objects are unstable to bar-forming $m = 2$ modes, a well-known result (Hohl 1971; Ostriker & Peebles 1973; Binney & Tremaine 1987). One may see these as energy-conserving instabilities, where stars may acquire some energy because the gravitational potential in which they move is not constant in time. In such systems, typically $\approx 1\%$ of the mass (originally, all stars on high angular momentum orbits) reaches orbits outside the main body of the resulting bar, becoming detached (but not unbound) from the galaxy. We may expect the total fraction of mass lost in hydrodynamic systems to be comparable to this, since here too parcels of fluid on high angular momentum orbits interact with the rest of the system, largely through the gravitational potential; in fact, according to the models of Salgado et al. (1994), the contribution of the pressure gradient to the support of the outer layers of the star is much smaller than the centrifugal one. Moreover, neutron stars differ from stellar dynamical systems because of losses by gravitational waves (GWs), and possibly neutrinos, which weaken gravitational binding of the outermost layers. It is sufficient that $\approx 10\%$ of the whole star mass at infinity ends up as GWs for the equatorial belt to become rotationally supported; this corresponds to 0.3 $M_\odot$ radiated away, well within current estimates of GW emission efficiencies in nonaxisymmetric collapses to black hole (Smarr 1979). Neutrinos not trapped inside the collapsing SMNS will also carry away binding mass, or may provide pressure for pushing the equatorial belt into centrifugal equilibrium. We thus expect that the SMNS might typically shed a few percent of its total mass to centrifugal equilibrium. This implies $M_l \gtrsim 0.1\ M_\odot$.

The ensuing configuration is of a largish black hole ($M \approx$ 2–3 $M_\odot$), spinning at subcritical speed ($j \approx 0.6$), surrounded by a thin equatorial belt of matter of mass M_l, which could not be immediately accreted because its angular momentum, supplemented by the factors discussed above, makes it rest in centrifugal equilibrium. The binding energy of the leftover mass M_l is

$$E_b = 3 \times 10^{52} \text{ ergs } \frac{M_{SMNS}}{2\ M_\odot}\ \frac{M_l}{0.1\ M_\odot}\ \frac{20 \text{ km}}{R_{eq}} \ . \quad (2)$$

This binding energy must be radiated as the leftover accretes onto the black hole, and it is this energy release that powers the GRB proper, not the gravitational collapse per se. Even the most powerful bursts (Kulkarni et al. 1998) only require $M_l \approx 0.1\theta^2\ M_\odot$, where $\theta^2/4$ is the beaming fraction.

The leftover material will be threaded by the same large-scale magnetic field lines that threaded it inside the NS; the field will be amplified by shearing induced by differential rotation inside the accreting torus. Assuming that the angular

speed difference between the front and back edges nearly equals ω, a field of initial strength $B \approx 10^{12}$ G will be amplified to 10^{15} G (i.e., well below equipartition with differential rotation, an upper limit that it is not necessary to attain) in ≈ 10 differential rotations; the amplification timescale is $t_a \approx 0.01$ s, comparable to the bursts' rise times. At this point, the torus will radiate at the rate $\dot{E} \approx -B^2\omega^4 r^6/6c^3 = 10^{50}$ ergs s^{-1}, enough to power a burst. The details of the conversion of Poynting flux into a magnetized relativistic wind have been discussed in both different and similar contexts by Usov (1992, 1994), Mészáros & Rees (1997b), and Katz (1997). Alternatively, one may invoke the Blandford-Znajek (1977) mechanism to extract the much larger hole's rotational energy (Mészáros et al. 1998b).

The energy release rate computed above of course exceeds the Eddington luminosity of the minitorus by more than 10 orders of magnitude. The energy deposition occurs outside the minitorus, but it is of course quite likely that some mass is blown off the torus, for instance by photons that hit it and heat it. This baryon outflow will certainly contaminate the zones of the relativistic wind that are most distant from the spin axis, because we do not expect any efficient angular momentum exchange in the process. Still, this implies that the outflow may be nonisotropic, and moderately beamed. The situation that we envision is similar to that described by Mészáros & Rees (1997b), with a smoothly tapering beam.

3. A CLEAN ENVIRONMENT

The volume immediately surrounding the SMNS will be vacated of baryons by two effects. First, the SN explosion accompanying the SMNS's formation will sweep away ISM baryons. We can gauge the importance of this effect by considering the fact that most of the mass around the Crab pulsar, the only plerionic remnant for which relevant data are available, is well known to be located in the filaments, which amount to no more than $\sim 5 M_\odot$ (Fesen, Shull, & Hurford 1997), within about 3′. At a canonical distance of 2 kpc, this corresponds to a present average baryon number density $n \approx 10$ cm^{-3}. Assuming uniform expansion in the period from t_{sd} to the Crab's current age implies that at t_{sd}, the average baryon density was $n \approx 10^7$ cm^{-3}. Within a distance $D = 10^{15}$ cm, more than enough to guarantee millisecond variability (which requires $D \approx 10^{12}$ cm; Rees & Mészàros 1994), this corresponds to a total baryon mass of $M \approx 3 \times 10^{-5} M_\odot$, below the $10^{-4} M_\odot$ upper limit required to get a significantly relativistic expansion. We regard the above estimate as an upper limit because the Crab is well known to have an unusually low expansion velocity (a full factor of 3 below other Type II SNRs; Fesen, Shull, & Hurford 1997), and because it considers the mass in the filaments as if it were uniformly distributed over the whole volume rather than concentrated in the observed thin shell.

The second effect is due to the large energy release by the magnetic dipole rotation after the explosion,

$$\dot{E} = -\frac{B^2 R^6 \omega^4 \sin^2\alpha}{6c^3}$$

$$= -3 \times 10^{43} \text{ ergs s}^{-1} \left(\frac{B}{10^{12} \text{ G}}\right)^2$$

$$\times \left(\frac{R}{15 \text{ km}}\right)^6 \left(\frac{\omega}{10^4 \text{ s}^{-1}}\right)^4 \sin^2\alpha; \qquad (3)$$

since a large fraction of this energy will be converted to photons, the ensuing luminosity in fact exceeds the Eddington luminosity for the $\approx 3 M_\odot$ NS by about 4 orders of magnitude. This is of course identical to the well-known plerion model for SN remnants (Ostriker & Gunn 1971; Reynolds & Chevalier 1984), except that the total energy released ($\approx 10^{53}$ ergs) is much larger; for this reason, we cannot directly apply these models to our case, except to note that its effect will surely be to lower the naive baryon mass estimate obtained above ($\approx 3 \times 10^{-5} M_\odot$).

4. OBSERVATIONAL IMPLICATIONS

We now discuss observational signatures of the occurrence of this model. It shares with all models for which GRBs are generated from young, massive stars dying possibly exotic deaths, all predictions that associate GRBs with star-forming regions. In particular, GRBs' afterglows should be localized within star-forming galaxies, and their redshift distributions ought to closely match the history of star formation in the universe (Madau et al. 1996; Lilly et al. 1996). Furthermore, the optical afterglow may at times be absorbed by dust present in the star-forming regions, and the X-ray afterglow may show absorption by large equivalent column depths by neutral elements. In fact, in the t_{sd} yr elapsed since the first explosion, the SN shock front will have reached a distance of $R_s = 10^{18}$ cm $(v_{exp}/3 \times 10^9$ cm s$^{-1})(t_{sd}/10$ yr). The total column depth of baryons, at distance R_s, is $N_b = 10^{21}$ cm$^{-2}(M_{ej}/10 M_\odot)(R_s/10^{18}$ cm)$^{-2}$, where M_{ej} is the ejecta mass; this column depth may perhaps be observable (Murakami et al. 1997). Thus, we may sometimes expect to see the ionization edges discussed by Mészáros & Rees (1998a, 1998b) arising from dense, stationary outlying material.

We argued above that in its maxi-Crab phase, the SMNS will vacate of baryons a sizeable cavity around itself. Outside this cavity, however, there will be an outwardly increasing density gradient, a well-known property of Sedov solutions. In fact, matter behind the SN shock is subsonic with respect to the outwardly propagating shock, so that it must be in approximate pressure equilibrium with previously shocked material. However, since the shock is decelerating, just-shocked material will be at a lower temperature than material shocked long before, and in pressure equilibrium this implies higher densities. The exact Sedov solutions yield for the innermost regions $n_{ISM} \propto r^d$, where $d = 3/(\gamma - 1) = 9/2$ (Landau & Lifshitz 1979). If the time between the SN explosion and the burst is rather short, it may happen that the collision between different shells, which gives rise to the burst proper (Paczyński & Xu 1994), takes place in a region that the maxi-Crab effect has not yet managed to empty of baryons. In this case, the afterglow will propagate into the Sedov solution, i.e., inside an outwardly increasing density gradient. In the afterglow, the time t_{nr} of the transition to the nonrelativistic expansion regime is reasonably sensitive to the density gradient; Mészáros et al. (1998a) find $t_{nr} \approx t_\gamma \Gamma^q$, where t_γ is the duration the burst in the γ-ray band, $\Gamma \approx 100$ is the initial Lorenz factor of the ejecta, and $q = (8 + 2d)/(3 + d)$, for adiabatic expansion (note that we define as $-d$ what Mészáros et al. 1998a call $+d$). For expansion in a wind, $d = -2$, and thus $q = 4$, while for expansion into an adiabatic Sedov solution, $d = 9/2$, $q = 2.2$, so that the transition to the nonrelativistic regime is anticipated by up to 4 orders of magnitude in time. For $t_\gamma = 1$ s, $q = 2.2$ and $t_{nr} \approx 7$ hr. After t_{nr}, the luminosity steepens considerably, to $t^{-2.4}$ or steeper (Wijers, Rees, & Mészáros 1997). This

provides a natural explanation for those cases in which either no afterglow is seen (in the hard X-ray GRB 970111, Feroci et al. 1998; in the optical GRB 970828, Groot et al. 1998a), or it is seen to decrease more steeply than the other afterglows (GRB 980326; Groot et al. 1998b).

The formation rate of SMNSs in an L_* galaxy, $1/t_{SMNS}$, is of course expected to be of the order of, but less than, the rate of formation for pulsars, in our Galaxy currently about $1/t_p$, where $t_p \approx 100$ yr. If the formation rate is around a factor of 100 lower than this, as suggested by rough considerations of the range of mass and angular momentum over which SMNSs form, there is ample room to account for all observed GRBs, including a moderate degree of beaming. It may prove interesting to search for such objects. The fraction of the time between each birth through which they are active in the maxi-Crab phase is t_{sd}/t_{SMNS}, so that there will be one L_* galaxy every t_{SMNS}/t_{sd} containing an active one. Given the luminosity density of the universe ($j_0 = 1.7 \times 10^8 \, L_\odot \, \mathrm{Mpc}^{-3}$ in the V band) and the typical luminosity of an L_* galaxy like our own ($L_* = 10^{10} \, L_\odot$ in the V band; Binney & Tremaine 1987), we find that the closest such active maxi-Crab should lie at a distance $R = (3L_* t_{SMNS}/4\pi j_0 t_{sd})^{1/3} = 52$ Mpc, where we used equation (1) and $t_{SMNS} = 10^5$ yr, which is necessary if SMNSs must account for GRBs. At this distance, the flux of equation (3)

corresponds to $\approx 10^{-10}$ ergs s^{-1} cm^{-2}, pulsed with frequency $\nu = \omega/2\pi \approx 1000$–2000 Hz. If the Crab pulsar is anything to go by, we expect an X-ray flux of $\approx 10^{-11}$ ergs s^{-1} cm^{-2} (only a fraction of which will be pulsed) from a pointlike source well offset from the center of the galaxy. A detection of the pulsations from such a source is beyond the capabilities of *Rossi XTE*, but well within reach of very large area X-ray telescopes such as *XMM* and, especially, the proposed *Constellation-X*.

In summary, we presented a supranova model for gamma-ray bursts, i.e., a model with a two-step collapse to a black hole not induced by accretion. The advantages of this model lie in its ability to keep the environment baryon-free, both because of the SN explosion that goes with the first collapse, which cleans up the surroundings, and because of the clean, silent collapse of the NS once centrifugal support weakens critically. Second, we discuss a rotating magnetic dipole radiating much energy through its large rotation rate, which injects into the NS's surroundings a luminosity exceeding Eddington's by 4 orders of magnitude, and which speeds up as a consequence of energy losses. This also contributes to keeping the NS's environment baryon-clean. A rough estimate suggests that the total environment pollution by baryons may be in the range of $\approx 10^{-4} \, M_\odot$.

REFERENCES

Andersson, N. 1998, ApJ, 502, 708
Binney, J., & Tremaine, S. 1987, Galactic Dynamics (Princeton: Princeton Univ. Press)
Blackman, E. G., Yi, I., & Field, G. B. 1996, ApJ, 473, L79
Blandford, R. D., & Znajek, R. L. 1977, MNRAS, 179, 433
Bloom, J., Sigurdsson, S., & Pols, O. 1998, MNRAS, in press (astro-ph/9805222)
Chandrasekhar, S. 1970, Phys. Rev. Lett., 24, 611
Cook, G. B., Shapiro, S. L., & Teukolsky, S. A. 1994, ApJ, 424, 823
Feroci, M., et al. 1998, A&A, 332, L29
Fesen, R. A., Shull, J. M., & Hurford, A. H. 1997, AJ, 113, 354
Friedman, J. L., Ipser, J. R., & Sorkin, R. D. 1988, ApJ, 325, 722
Goodman, J. J. 1997, NewA, 2, 449
Grindlay, J. E. 1998, ApJ, in press (astro-ph/9808242)
Groot, P. J., et al. 1998a, ApJ, 493, L27
———. 1998b, ApJ, 502, L123
Hohl, F. 1971, ApJ, 168, 343
Katz, J. I. 1994, ApJ, 432, L107
———. 1997, ApJ, 490, 633
Kulkarni, S. R., et al. 1998, Nature, 393, 35
Landau, L. D., & Lifshitz, E. M. 1979, Fluid Dynamics (Oxford: Pergamon)
Lilly, S. J., et al. 1996, ApJ, 460, L1
Madau, P., et al. 1996, MNRAS, 283, 1388
Mészáros, P., & Rees, M. J. 1997a, ApJ, 476, 232
———. 1997b, ApJ, 482, L29
———. 1998a, ApJ, 502, L105

Mészáros, P., & Rees, M. J. 1998b, MNRAS, in press (astro-ph/9806183)
Mészáros, P. J., Rees, M. J., & Wijers, R. A. M. J. 1998a, ApJ, 499, 301
———. 1998b, NewA, in press (astro-ph/9808106)
Murakami, T., et al. 1997, in Gamma Ray Bursts, ed. C. Meegan, R. Preece, & T. Koshut (New York: AIP), 435
Narayan, R., Paczyński, B., & Piran, T. 1992, ApJ, 395, L83
Ostriker, J. P., & Gunn, J. E. 1971, ApJ, 164, L95
Ostriker, J. P., & Peebles, J. P. E. 1973, ApJ, 186, 467
Pacini, F. 1967, Nature, 216, 567
Paczyński, B. 1998, ApJ, 482, L29
Paczyński, B., & Rhoads, J. E. 1993, ApJ, 395, L83
Paczyński, B., & Xu, G. 1994, ApJ, 424, 708
Rees, M. J., & Mészáros, P. 1992, MNRAS, 258, 41P
———. 1994, ApJ, 430, L93
Reynolds, S. P., & Chevalier, R. A. 1984, ApJ, 278, 630
Salgado, M., Bonazzola, S., Gourgoulhon, E., & Haensel, P. 1994, A&A, 291, 155
Smarr, L. L. 1979, Sources of Gravitational Radiation (Cambridge: Cambridge Univ. Press)
Usov, V. V. 1992, Nature, 357, 472
———. 1994, MNRAS, 267, 1035
Vietri, M. 1997, ApJ, 478, L9
Wijers, R. A. M. J., & Galama, T. 1998, ApJ, in press (astro-ph/9805341)
Wijers, R. A. M. J., Rees, M. J., & Mészáros, P. 1997, MNRAS, 288, 51P
Woosley, S. E. 1993, ApJ, 405, 273
Yi, I., & Blackman, E. G. 1997, ApJ, 482, 383