Supranova Model for the Delayed Reddened Optical Excesses Detected in Several GRBs

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ABSTRACT

Excess reddened optical emission has been measured in the power-law decay afterglow light curves of several GRBs some tens of days after the GRB events. This emission is claimed to be a signature of a supernova (SN) taking place at the same time as the GRB event, in support of the collapsar scenario. Observations of features in prompt and afterglow X-ray spectra of GRBs require high density, high metallicity material near GRB sources, supporting the supranova (SA) scenario. These conflicting lines of evidence are resolved in likely regimes of parameter space if the supernova remnant (SNR) shell is illuminated and heated by a pulsar wind (PW), and cools following the GRB event. The heating of the SNR shell by the PW is parameterized, and simplified expressions for the synchrotron, synchrotron self-Compton, and thermal radiation fields within the SNR shell and pulsar wind bubble are derived for monoenergetic injection of wind electrons. A cooling SNR shell could produce the excess emission detected from GRB 970228, GRB 980326, and GRB 011121. X-ray synchrotron radiation from the PW electrons provides an important source of ionizing radiation.

Subject headings: gamma-rays: bursts — gamma-rays: theory — radiation processes: nonthermal — pulsars

1. Introduction

Two leading scenarios to explain the origin of gamma-ray bursts are the collapsar and supranova (SA) models (see Mészáros (2002); Dermer (2002) for review). The collapsar model (e.g., Wang & Woosley 2002; Woosley 1993) assumes that GRBs originate from the collapse of a massive star to a black hole. During the collapse process, a nuclear-density,
several Solar-mass accretion disk forms and accretes at the rate of $\sim 0.1-1 \, M_\odot \, s^{-1}$ to drive a baryon-dilute, relativistic outflow through the surrounding stellar envelope. The duration of the accretion episode corresponds to the prompt gamma-ray luminous phase, which is commonly thought to involve internal shocks. A wide variety of collapsar models can be envisaged (Fryer, Woosley, & Hartmann 1999), but their central feature is the one-step collapse of the core of a massive star to a black hole.

Optical observations of GRB afterglows reveal significant evidence for enhanced emission in excess of an extrapolation of the optical power-law decay spectra in a few cases, namely GRB 980326 (Bloom et al. 1999), GRB 970228 (Reichart 1999; Galama et al. 2000), and GRB 011121 (Bloom et al. 2002; Price et al. 2002). The spectrum of the excess is reddened at frequencies above the spectral peak, as might be expected for a thermal emitter. The peak of the excess emission occurs tens of days after the GRB event (in the local frame), and the SN event precedes GRB 011121 by $\sim 3-5 \, (\pm 5)$ days (Bloom et al. 2002), consistent with a one-step collapsar model.

The supranova model (Vietri & Stella 1998) of GRBs involves a two-step collapse process of an evolved massive star to a black hole through the intermediate formation of a neutron star with mass exceeding several Solar masses. The neutron star is initially stabilized against collapse by rotation, but the loss of angular momentum support through magnetic dipole and gravitational radiation leads to collapse of the neutron star to a black hole after some weeks to years. The accretion-induced collapse of a neutron star in a binary system could also form a GRB (Vietri & Stella 1999). A two-step collapse process means that the neutron star is surrounded by a SNR shell of enriched material at distances of $\sim 10^{15}-10^{17}$ cm from the central source. The earlier SN could yield $\sim 0.1-1 \, M_\odot$ of Fe in the surrounding vicinity. The discovery of variable Fe absorption during the prompt emission phase of GRB 990705 (Amati et al. 2000) and X-ray emission features in the afterglow spectra of GRB 991216 (Piro et al. 2000) has been interpreted in terms of the SA model (Lazzati et al. 2001; Böttcher et al. 2002; Ballantyne et al. 2002). Low significance X-ray absorption features observed in GRB 970508 (Piro et al. 1999), GRB 970828 (Yoshida et al. 2001), and GRB 000214 (Antonelli et al. 2000) can also be modeled within the SA scenario, preferably with photoionization reflection spectra (Lazzati et al. 1999). The multiple features detected in GRB 011211 (Reeves et al. 2002) are also suggested to arise from an excited SNR shell at $\sim 10^{15}$ cm from the central source.

A PW and pulsar wind bubble (PWB) consisting of a quasi-uniform, low density, highly magnetized pair-enriched medium within the SNR shell is formed by a highly magnetized neutron star during the period of activity preceding its collapse to a black hole (Königl & Granot 2002). The interaction of the PW with the shell material will fragment and
accelerate the SNR shell, and the PW emission will be a source of ambient radiation that can be Comptonized to gamma-ray energies (Inoue et al. 2002; Guetta & Granot 2002).

Here we propose that the nonthermal leptonic PW radiation provides a source of hard ionizing radiation to heat and photoionize material in the SNR shell, which is described in terms of a partial covering model as a result of the strong PW/SNR interaction. This radiation heats the shell, and the escaping thermalized radiation is seen as the delayed reddened optical excesses in GRB optical afterglow light curves. The source of ionizing nonthermal radiation, which is extinguished when the GRB occurs, would provide a nonthermal precursor to a GRB, though at low flux. Various other features of this system are outlined.

2. Pulsar Wind Physics

A very rapidly rotating, strongly magnetized neutron star is formed in the first step of the SA scenario. For a neutron star with surface polar field of $10^{12} B_{12} \text{G}$, a rotation rate of $10^4 \Omega_4 \text{rad s}^{-1}$, and a radius of $15 r_{15} \text{km}$, the radiated power $L_0 \equiv 7 \times 10^{44} B_{12}^2 r_{15}^6 \Omega_4^4 \text{ergs s}^{-1}$. The quantity $E_{\text{rot}} = \frac{1}{2} j (GM^2/c^2) \Omega \cong 2 \times 10^{53} j_{0.5} M_3^2 \Omega_4 \text{ergs} \equiv 10^{53} E_{\text{rot}}^{53}$ is the initial rotational energy of the neutron star, expressed in terms of the neutron-star mass $M = 3 M_3 \text{M}_\odot$ and the dimensionless angular momentum $j_{0.5} = GM^2/2c$, which remains roughly constant during collapse (Vietri & Stella 1998). The spin-down time $t_{sd} = (E_{\text{rot}}/L_0) \cong 2.8 \times 10^8 j_{0.5} M_3^2 (B_{12}^2 r_{15}^6 \Omega_4^3)^{-1} \text{s}$. The torque equation, $\dot{\Omega} = -K \Omega^n$, where the braking index $n = \ddot{\Omega}/\dot{\Omega}^2$, can be solved to give $\Omega(t) = \Omega_0 (1 + t/t_d)^{-1/(n-1)}$, provided that $K$ remains constant (e.g., Atoyan 1999). (For the Crab, $n = 2.5$.) We assume that the spin-down power is radiated as a PW in the form of particle and Poynting flux, so that the wind luminosity $L_w \cong L_{sd} = (\xi_e + \xi_p + \xi_B) L_{sd} = L_0/(1 + t/t_{sd})^k$, where $k = (n+1)/(n-1)$, and we have divided the wind power into leptonic (“e”), hadronic (“p”), and electromagnetic (“B”) fractions $\xi_i$, $i = e, p, B$.

In the SA model, rotational stability is lost due to spindown, and the neutron star collapses to a black hole to produce a GRB at time $t_{coll}$. Here we assume $t_{coll} \lesssim t_{sd}$, so that for $k > 1$ the wind energy $E_w(t) = \eta (k-1)^{-1} L_0 t_{sd} \left[ 1 - (1 + (t/t_{sd})^{1-k} \right] \rightarrow \eta L_0 t$ when $t \lesssim t_{sd}$, where the parameter $\eta$ roughly accounts for the escape of wind energy from the SNR shell as well as losses of wind energy due to radiation, and is assumed here to be constant with time. The equation for the SNR shell dynamics in a spherical approximation for the SNR shell is $M_{SNR} \dot{R} = 4\pi R^2 p_w(t)$, where $M_{SNR} = m_{SNR} M_\odot$ is the SNR shell mass, the wind pressure $p_w(t) \approx E_w(t)/3V_{pwb}$, the PWB volume $V_{pwb} = 4\pi R^3(t)/3$, and the PWB radius
\( R(t) = \int_0^t dt' \, v(t') \) (Inoue et al. 2002). This equation can be solved to give
\[
R(t) = v_0 t \left(1 + \sqrt{\frac{t}{t_{\text{acc}}}}\right), \quad \text{where} \quad t_{\text{acc}} = \frac{3M_{\text{SNR}}v_0^2}{4\eta L_0} \approx 1.9 \times 10^7 \frac{m_{\text{SNR}} \beta_{-1}^2}{\eta B_{12} r_{15}^3 \Omega_4^2} \text{ s} \quad (1)
\]
\( (\text{Guetta} \& \text{Granot} 2002) \). It follows that \( t_{\text{acc}}/t_{\text{sd}} = E_{\text{SNR}}^{\text{ke}}/\eta E_{\text{rot}} \approx 0.1E_{53}^{\text{SNR}}/\eta E_{53}^{\text{rot}} \), where \( \frac{1}{2}M_{\text{SNR}}v_0^2 \approx 9 \times 10^{51} \text{ ergs} \) is the SNR shell kinetic energy, and \( v_0 = 0.1\beta_{-1}c \) is the SNR shell coasting speed. If \( \eta \approx 1 \), then \( t_{\text{acc}} \ll t_{\text{sd}} \) for standard values, and shell acceleration must be considered. A highly porous shell has \( \eta \ll 1 \). This will occur for strong clumping of the ejecta, which takes place on the Rayleigh-Taylor timescale \( t_{\text{RT}} \approx \sqrt{x/R} \), where the characteristic clumping size scale is \( x \equiv f R \) (Inoue et al. 2002). Hence \( t_{\text{RT}}/t_{\text{acc}} = f^{1/2} (t/t_{\text{acc}})^{3/4} \) in the regime \( t \lesssim t_{\text{acc}} \). This shows that small-scale \( (f \ll R) \) clumping will occur when \( f \ll \sqrt{t/t_{\text{acc}}} \). Only a detailed hydrodynamic simulation can characterize the porosity of the shell due to effects of the pulsar wind, which \( \eta \) parameterizes. Here we assume, in contrast to the picture of Königl \& Granot (2002), that a large fraction of the wind energy escapes the SNR shell and \( \eta \sim 0.1 \), so that shell acceleration can be neglected. This effect also causes the shell to become effectively Thomson thin much earlier than estimated on the basis of a uniform shell approximation, except for high density clouds with small covering factor. A general treatment is needed to treat shell dynamics, including travel-time delays of the pulsar wind and deceleration of the SNR shell by the surrounding medium. When \( t_{\text{coll}} \lesssim t_{\text{acc}}, t_{\text{sd}} \), as assumed here, \( R \approx v_0 t \).

3. Radiation Field in SNR Cavity

The strong PW provides a source of nonthermal leptons, hadrons, and electromagnetic field. Dominant radiation components considered here are leptonic synchrotron and SSC radiation, and thermal emission from the inner surface of the SNR shell which is heated by the nonthermal radiation.

3.1. Wind Synchrotron Radiation

The volume-averaged mean magnetic field \( B_w \) in the SNR cavity powered by the PW is obtained by relating the wind magnetic field energy density \( u_B = B_w^2(t)/8\pi = e_{Bw}u_w(t) = 3e_{Bw}\eta L_0 t/4\pi R^3 \) through the magnetic-field parameter \( e_{Bw} \), also assumed to be constant in time. Thus
\[
B_w(t) = \frac{B}{t}, \quad \text{where} \quad B = \frac{6e_{Bw}L_0 \eta}{v_0^2} \approx 4 \times 10^8 \sqrt{e_{Bw} \eta} \left(\frac{B_{12} r_{15}^3 \Omega_4^2}{\beta_{-1}^{3/2}}\right) \text{ G s} \quad (2)
\]
Let $\gamma_w = 10^5 \gamma_5$ represent the typical Lorentz factors of leptons in the wind zone of the pulsar. This value should remain roughly constant when $t \lesssim t_{sd}$. A quasi-monoenergetic proton/ion wind may also be formed, though we only consider leptonic processes here, and furthermore do not treat nonthermal particle acceleration at the PW/SNR shell boundary shock. When synchrotron losses dominate, the mean Lorentz factor of a distribution of leptons with random pitch angles evolves in response to a randomly oriented magnetic field of mean strength $B$ according to the expression

$$\gamma_w = \left[ 1 - \frac{1}{\gamma} - T_B \left( \frac{1}{t_i} - \frac{1}{t} \right) \right]^{-1},$$

(3)

where $t_i$ is the injection time, $\gamma_i = \gamma_w$ is the injection Lorentz factor, and

$$T_B \gamma_w = \frac{\sigma_T B^2}{6\pi m_e c} \sim 2 \times 10^{13} \epsilon_{Bw} \eta \left( \frac{B^2}{10^4 \gamma_5^4} \right) \gamma_5 \text{ s}.$$  

(4)

Writing the nonthermal electron injection function as

$$\frac{dN_e(\gamma_i; t_i)}{dt_i d\gamma_i} = \frac{\eta \xi_e L_0}{m_e e^2 \gamma_w} \delta(\gamma_i - \gamma_w),$$

(5)

the electron Lorentz factor distribution $dN_e(\gamma; t)/d\gamma$ can be solved to give

$$\frac{\gamma^2 dN_e(\gamma; t)}{d\gamma} = \hat{N}_e T_B \gamma^2 _w \left( \frac{1}{t} + \frac{1}{\gamma} - \frac{1}{\gamma} \right)^{-2}.$$  

(6)

In this expression, $\hat{N}_e \equiv \eta \xi_e L_0 / (m_e e^2 \gamma_w)$, and we introduce the dimensionless quantities $\hat{t} = t/T_B \gamma_w$ and $\hat{\gamma} = \gamma/\gamma_w$. Adiabatic losses, which are significant on time scales $\sim t/3$, are negligible in comparison with synchrotron losses of wind electrons when $\hat{t} \ll 1$, and we restrict ourselves to this regime.

The $\nu L_\nu$ synchrotron radiation flux $\nu L_\nu^{\text{syn}} \approx \frac{1}{2} u_B c \sigma_T \gamma_5^3 N_e(\gamma_s)$, where $\gamma_s = \sqrt{\epsilon/\epsilon_B}$, $\epsilon = h\nu/m_e c^2$, $\epsilon_B = B/B_{cr}$, and the critical magnetic field $B_{cr} = 4.41 \times 10^{13}$ G. Thus

$$\nu L_\nu^{\text{syn}}(\epsilon) = c \sigma_T \frac{B_{cr}^2}{16\pi} \left( \hat{N}_e T_B \gamma_w^2 \right) \left( \frac{u}{t} \right)^{3/2} \sqrt{\epsilon} \left( \frac{1}{t} + \sqrt{\frac{u \gamma_w}{\epsilon t}} - 1 \right)^{-2}, \text{ for } \epsilon \leq \epsilon_{\text{max}},$$

(7)

where $u \equiv B/\gamma_w B_{cr} T_B \approx 4.5 \times 10^{-19} \beta_{-1}^{3/2} / (\sqrt{\eta e_{Bw} B_{12} r_{15}^3 \Omega_4^2} \gamma_5)$, and

$$\epsilon_{\text{max}} \approx \epsilon_B \gamma_w^2 = \frac{B}{B_{cr}} \gamma_w^2 \frac{t}{u} \approx \frac{4.4 \times 10^{-9}}{t} \frac{\beta_{-1}^{3/2} \gamma_5^5}{\sqrt{\eta e_{Bw} B_{12} r_{15}^3 \Omega_4^2}}.$$  

(8)

The synchrotron cooling timescale for wind electrons is $t_{\text{syn}} = \gamma_w T_B t^2$. In the regime $\hat{t} \lesssim \beta_0$ where the wind electrons strongly cool, we can approximate the synchrotron spectrum by

$$\nu L_\nu^{\text{syn}}(\epsilon) \approx \frac{3}{8} \xi_e L_0 \left( \frac{\epsilon}{\epsilon_{\text{max}}} \right)^{1/2} H(\epsilon; \epsilon_0, \epsilon_{\text{max}}),$$

(9)
where $H(x; a, b) = 1$ for $a \leq x < b$, and $H = 0$ otherwise. The value of $\hat{t}$ at $t_{sd}$ is
\[
\frac{t_{sd}}{T_B \gamma_w} \approx 1.4 \times 10^{-5} \frac{j_{0.5} M_z^2 \beta_{-1}^3}{e_{Bw}^3 \eta B_{12}^{18} \Omega_{14}^{11} \gamma_5}.
\]
(10)

Strong cooling holds when $\hat{t} < t_{sd}/T_B \gamma_w$, noting that time in physical units can be found from equation (4).

### 3.2. Wind Synchrotron Self-Compton Radiation

The relative importance of synchrotron self-Compton (SSC) to synchrotron cooling is given by the ratio $\rho = u_{syn}/u_B$ of the synchrotron radiation energy density to $u_B = B^2 / 8\pi t^2$ (see eq.[2]). The PW synchrotron radiation energy density $u_{syn} \simeq \kappa (\nu L_{\nu}^{max})/4\pi R^2 c$, where $\nu L_{\nu}^{max} \simeq \xi_{e} L_0$ and the parameter $\kappa \geq 1$ is a reflection factor such that $\kappa = 1$ corresponds to perfect absorption or direct escape (small covering factor) of the PW synchrotron radiation, and $\kappa \gg 1$ corresponds to a highly reflecting SNR shell with large covering factor. One obtains $\rho = 0.033 \kappa \xi_{e}\beta_{-1}/e_{Bw}\eta$. For large porosity with $\kappa \approx 1$ and $\eta \approx 0.1$, we see that the SSC contribution is small compared to the synchrotron flux when $\xi_{e}\beta_{-1}/e_{Bw} \lesssim 1$. We restrict ourselves to this regime where equation (3) is valid.

In the $\delta$-function approximation for the energy gained by a photon upon being scattered by relativistic electrons (Dermer, Sturner, and Schlickeiser 1997), with scattering restricted to the Thomson regime, the $\nu L_{\nu}$ SSC spectrum from the cooling wind electrons is given by
\[
\nu L_{\nu}^{SSC} \simeq \frac{\sigma_T \kappa \xi_{e} L_0}{8 \pi v_0^2 T_B} \frac{\dot{N}_\epsilon}{t^2} \sqrt{\frac{\epsilon}{\epsilon_{max}}} \left( \frac{1}{a_0 (x_0 + x)} + a_0^{-2} \ln \left( \frac{x}{a_0 + x} \right) \right)|_{x_0}^{x_1},
\]
(11)
where $a_0 = \hat{t}^{-1} - 1$, $x_0 = \gamma_w \sqrt{\epsilon_0/\epsilon}$, and $x_1 = \gamma_w \sqrt{\max(\epsilon, \epsilon^{-1}, \epsilon_{max})/\epsilon}$, where we have used approximation (9) for the synchrotron spectrum.

### 3.3. Thermal Emission from the Interior of the Shell

The SNR shell is pictured as having been shredded by the PW and therefore highly porous, though permeated with small scale ($\sim 10^{12}$-10$^{14}$ cm) density irregularities consisting largely of metals (Dermer & Mitman 1999; Böttcher et al. 2002). Note that $0.1 m_{56, -1} M_\odot$ of $^{56}$Ni provides a radioactive decay power of $\sim 2 \times 10^{42} m_{56, -1}$ ergs s$^{-1}$ before decaying on an $\approx 120$ day timescale. The partial covering factor $p_c$ ($\leq 1$) corresponds to the covering fraction by optically-thick, dense shell inhomogeneities. Recognizing the severe limitations
in the following expression, we determine the effective temperature $T_{\text{eff}}$ in the interior of the shell through the relation

$$4\pi R^2\sigma_{\text{SB}}T_{\text{eff}}^4p_c \approx A_p\frac{\xi_e L_0}{\kappa} + 2 \times 10^{42}m_{56,-1}\exp(-t/120 \, \text{d}) \, \text{ergs s}^{-1},$$

(12)

where $A(\leq 1)$ is an absorption coefficient, dependent in general on the evolving radiation spectrum, though $A$ is here assumed constant. The wind heating is more important than heating from the decay of radioactive Ni when $(A/0.1)(p_c/0.1)(\xi_e/1/3) \gtrsim m_{56,-1}/B_{12}^2r_{15}^6\Omega_4^4$. In this regime,

$$T_{\text{eff}} = \left(\frac{A\xi_e L_0}{4\pi R^2\sigma_{\text{SB}}}\right)^{1/4} \approx \frac{825A^{1/4}}{\sqrt{t/t_{\text{sd}}}} \left(\frac{\xi_e}{1/3}\right)^{1/4} \frac{B_{12}^{3/2}r_{15}^{9/2}\Omega_4^{5/3}}{\sqrt{J_{0.5}\beta_{-1}M_3}} \, \text{K},$$

(13)

letting $\kappa \approx 1$ because $p_c \sim 0.1$. For standard parameters, the shell is heated to $\sim 10^4$ K temperatures when $t/t_{\text{sd}} \ll 1$. When $B_{12} \gtrsim 4$, this occurs at $t/t_{\text{sd}} \lesssim 1$. The PW heat source is extinguished once the GRB occurs, and the thermal emission from the SNR shell will decay on the dynamical timescale $t_{\text{dyn}} = R(t)/c \cong 324\beta_{-1}(j_{0.5}M_e^2/B_{12}^2r_{15}^6\Omega_4^4)(t/t_{\text{sd}})$ d, noting that the cooling timescale of a hot shell is generally shorter than $t_{\text{dyn}}$.

4. Results and Discussion

Fig. 1 shows the total $\nu L_\nu$ flux composed of synchrotron, SSC, and thermal components, calculated according to the above simple analysis. Standard parameters are $j_{0.5} = B_{12} = r_{15} = \Omega_4 = \beta_{-1} = \kappa = M_3 = m_{52,-1} = \gamma_5 = 1$, $\xi_e = e_{BW} = 1/3$, and $p_c = \eta = A = 0.1$. The heavy solid and dashed curves correspond to standard parameters with $t/t_{\text{sd}} = 0.1$ and 0.01 as labeled, with $t_{\text{dyn}} = 32.4$ and 3.2 days and $T_{\text{eff}} = 1700$ and 5420 K, respectively ($t_{\text{sd}} = 3240$ d). This system is in accord with the standard SA model and might explain the delayed excess emission in GRB 011121, contrary to the conclusions of Bloom et al. (2002). The set of three dotted curves employ standard parameters, except that $B_{12} = 4$, and $t/t_{\text{sd}} = 1, 0.1, and 0.01$ as labeled ($t_{\text{sd}} = 202$ d), with $t_{\text{dyn}} = 20.2, 2.0, and 0.2$ days (= 17 ks), and $T_{\text{eff}} = 3760, 11900, and 37630$ K, respectively.

The intense PW X-ray synchrotron emission provides a source of ionizing photons throughout the SNR shell that would decay on the timescale $t_{\text{dyn}}$ (the line flux would decay on $\approx 2t_{\text{dyn}}$). The $B_{12} = 4, t/t_{\text{sd}} = 0.01$ case might correspond to a GRB 011211-type system. The PW radiation field provides an additional source of ionizing photons in the SA model not considered by, e.g., Kumar & Narayan (2002), and would be energetically important for this case if $\gamma_5 \approx 0.1$. A varying ionization flux impinging on an expanding SNR shell will produce a characteristic kinematic signature of shell illumination. Our results differ from
the conclusions of Guetta & Granot (2002), because we assume that the SNR shell highly porous, as in the scenario of Inoue et al. (2002).

The thin solid curve is the result for standard parameters except that $B_{12} = 2$, $\gamma_w = 10^6$, and $t = 0.1 t_{sd} = 81$ d. For this case, $t_{dyn} = 8.1$ d and $T_{eff} = 4360$ K. It is apparent that a sizeable parameter regime could make quasi-thermal emission between $\sim 10^{14}$-$10^{15}$ Hz at the level of $\sim 10^{42}$-$10^{43}$ ergs s$^{-1}$ some tens of days after the GRB event.

Thus, even if the GRB takes place months after the SN, a delayed reddened excess from the cooling shell may be seen, for plausible parameters, some tens of days after the GRB. The PW heating is in addition to the heating due to the Ni→Co→Fe chain, and may relax the need to have $m_{56,-1} \gg 1$, as required to explain the unusually bright light curve of SN 1998bw (Iwamoto et al. 1998). The thermal emission from the SNR shell heated by the PW is proposed to be the origin of the reddened optical excess emission detected in GRBs 970228, 980326, and 011121.

This model predicts decaying nonthermal ionizing hard X-ray emission, which could be detected with INTEGRAL or Swift from nearby $z \lesssim 0.1$-0.3 GRBs when the PW synchrotron flux is at the level $\gtrsim 10^{45}$ ergs s$^{-1}$. A nonthermal optical emission signature would accompany the reddened excess, with a relative intensity that depends sensitively on $\gamma_w$. A hard X-ray sky survey instrument such as EXIST could monitor for the PW nonthermal emission preceding aligned and off-axis GRBs. The X-ray synchrotron radiation provides a source of external radiation to enhance photomeson production by energetic hadrons over the standard GRB model (Dermer 2002; Atoyan & Dermer 2001). High-energy neutrino and GLAST gamma-ray observations can distinguish the collapsar model and different versions of the SA model (A. M. Atoyan & C. D. Dermer, 2002, in preparation).

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Fig. 1.— Total model spectral energy distribution composed of synchrotron, SSC, and shell thermal emission from a PW and wind-heated SN shell. Standard parameters, given in the text, are used, except where noted in the legends.