Early GeV afterglows from gamma-ray bursts in pulsar wind bubbles

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ABSTRACT

Gamma-ray bursts may occur within pulsar wind bubbles (PWBs) under a number of scenarios, such as the supranova-like models in which the progenitor pulsar drives a powerful wind shocking against the ambient medium before it comes to death and produces a fireball. We here study the early afterglow emission from GRBs expanding into such a PWB environment. Different from the usual cold GRB external medium, the PWBs consist of a hot electron-positron ($e^+e^-$) medium with typical ‘thermal’ Lorentz factor of the order of $\gamma_w$, the Lorentz factor of the pulsar particle wind. After GRB blast waves shock these hot $e^+e^-$ pairs, they will emit synchrotron radiation peaking at GeV bands. It is shown that GeV photons suffer negligible absorption by the soft photons radiation field in PWBs. Thus, strong GeV emissions in the early afterglow phases are expected, providing a plausible explanation for the long-duration GeV emission from GRB940217 detected by EGRET. Future GLAST may have the potential to test this GRB-PWB interaction model.

Key words: gamma rays: bursts—pulsars: general

1 INTRODUCTION

Although the fireball shock scenario of GRBs and their afterglows has been well demonstrated by multiwavelength observations of afterglows (e.g. Wijers, Rees & Mészáros 1997; Piran 1999; van Paradijs, Kouveliotou & Wijers 2000), the nature of the central engine is not yet well known. There is observational evidence for the association of long duration bursts
with star-forming regions (Bloom et al. 2002), or possibly with supernovae (SNe; Galama et al. 1998). One class of candidates of the central engine involves that the supernova leaves behind a rotationally-supported supermassive or massive neutron star, which then shrinks by shedding angular momentum via a pulsar-type wind and finally collapses to a black hole surrounded by a disk (Supranova model; Vietri & stellar 1998) or transits to a strange star (Wang et al. 2000), producing the GRB fireball. The pulsar-wind bubble (PWB) forms when the relativistic wind from the pulsar shocks against its ambient matter and creates a pulsar nebula like the famous Crab nebula. Though different physical processes are involved for the death of the neutron star in these two scenarios, the PWB and the SN ejecta behaviors as well as the rotation energy of the progenitor pulsar are nearly identical; hence the two scenarios have rather similar implications to the evolution of GRB afterglows. This class of models in which SN explosion preceded the GRB event has recently gained support from the detections of strong Fe emission features in the X-ray afterglow spectra of some GRBs (e.g. Piro et al. 1999), particularly that of GRB991216 (Piro et al. 2000), as well as an absorption feature in the prompt emission of GRB990705 (Amati et al. 2000). These detections implies that a large amount ($\gtrsim 0.1 M_\odot$) of pure iron is located in the vicinity ($r \lesssim 10^{16}$ cm) of the GRB source (e.g. Vietri et al. 2001). Such a surprising large iron mass is most naturally produced in a SN explosion and the inferred distance of the SN ejecta implies the SN event precedes the GRB by months to years or shorter (see, e.g., Böttcher, Fryer & Dermer 2002).

Effects of wind nebulae resulted from the central ‘pulsar’ on the GRB afterglows have recently been examined by Inoue et al. (2002) and Königl & Granot (2002; hereafter KG) from different points of view. Assuming the SN ejecta may be fragmented by the plerion radiation, Inoue et al. (2002) study the inverse Compton effect of this radiation field on GRB external shock decelerated by the outlying baryonic material. Different from Inoue et al. (2002), in the model of KG, the afterglow shock wave propagates in a bubble composed mainly of shocked electrons and positrons that emanate from the central pulsar. KG argue that the pulsar wind nebula environment, bounded by the SNR shell, can account for the

* Strange quark matter is conjectured to be more stable than hadronic matter (Witten 1984). The existence of strange matter is allowable within uncertainties inherent in a strong-interaction calculation (Farhi & Jaffe 1984). Strange stars, composed of this kind of quark matter, may exist and could be born from a massive neutron star as it spins down (Wang et al. 2000)
† In the literature, Fe emission features are usually interpreted in the supranova model (Vietri & Stellar 1998). We note however that Fe line production is irrelevant to whether the progenitor pulsar collapses to a black hole or converts to a strange star.

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high electron and magnetic energy fractions ($\epsilon_e$ and $\epsilon_B$, respectively) inferred in a number of afterglow sources, as the PWBs are expected to have a significant $e^+e^-$ component and to be highly magnetized. Following the previous treatment of the PWB structure that was developed for the Crab and other nebulae (e.g. Rees & Gunn 1974; Kennel & Coroniti 1984, hereafter KC84), and assuming the magnetic pressure remains in approximate equipartition with the particle pressure throughout the shocked-wind bubble, they developed the structure model of PWB. On the other hand, a one-zone simplified version of the shock-wind model of PWB was developed by Chevalier (2000) and is shown to be well consistent with the X-ray luminosity of nebulae. Guided by these works, we attempt to give a simplified analytic description of the PWB that may exist in the supranove-like models (section 2), and study the early afterglow emission from blast wave propagating in the hot $e^+e^-$ bubble (section 3). The characteristic synchrotron frequencies of early afterglows are much higher than that of the afterglows in a typical interstellar medium (ISM). In section 4, we give a brief summary.

2 PULSAR WIND BUBBLES

Before collapsing, the supermassive or massive neutron star loses its energy through a pulsar-type wind, whose luminosity can be estimated from the magnetic dipole formula $L_w = B^2 R_{eq}^6 \Omega^4/6c^3 = 3 \times 10^{44} \text{ergs}^{-1} B_{12}^2 R_{eq,13}^6 \Omega_4^4$ and the spin-down time, in which half of the neutron star rotation energy $\Delta E_{rot} \sim 10^{53} \text{erg}$ need to be lost, is

$$t_{sd} = \Delta E_{rot}/L_w = 10 \text{ yr} \Delta E_{rot,53} B_{12}^{-2} R_{eq,13}^{-6} \Omega_4^{-4},$$

where $B = 10^{12} B_{12} \text{ Gauss}$, $R_{eq} = 13 R_{eq,13} \text{Km}$ and $\Omega = 10^4 \Omega_4 \text{s}^{-1}$ are, respectively, the dipole magnetic field, the typical equatorial radius and the angular velocity of the neutron star. Due to the unknown variations of these parameter values, we below consider $t_{sd}$ as a free parameter and parameterize the wind luminosity as $L_w = \Delta E_{rot}/t_{sd} = 3 \times 10^{44} \text{ergs}^{-1} \Delta E_{rot,53} t_{sd,1}^{-1}$, where $t_{sd} \equiv 10 \text{ yr} t_{sd,1}$.

We assume that the PWB is elongated along the rotation axis of the pulsar with an outer radius $R_b$ equal to that of the SN ejecta in this direction and pulsar wind is shocked at radius $R_s$ (Rees & Gunn 1974; KC84; KG). The pressure force of the expanding PWB is expected to accelerate the SN ejecta into a thin shell (Reynolds & Chevalier 1984) moving with a velocity $v_{ej} = 0.05-0.1 c$ as a significant amount of $\Delta E_{rot}$ is deposited into the bubble. Approximately, the outer radius of the nebula at time $t_{sd}$ equals to $R_b = v_b t_{sd} = 10^{18} \text{cm} v_{b,0.1} t_{sd,1}$, where $v_b = v_{ej} = 0.1 c v_{b,0.1}$ is the expansion velocity of the outer edge of the PWB. Fe line emission
requires that the pre-ejected SN ejecta should be no farther than \( r \sim 10^{16}\) cm. To reconcile this short Fe line-emitting distance with the radial distance of the PWB, KG suggest that the PWB become elongated in the polar direction because of anisotropic mass outflow from the GRB progenitor star. We note that the anisotropic pulsar wind power could also give rise to the elongated configuration. In fact, some pulsar wind nebulae have been found to show such anisotropic structure (e.g. Gaensler et al. 2001).

The pulsar rotation energy is assumed to go into a highly relativistic wind of the Lorentz factor \( \gamma_w \). Spectral modelling of the Crab nebula has given a current value of \( \sim 3 \times 10^6 \) for \( \gamma_w \). However, people do not know the initial value of \( \gamma_w \) for pulsars in their early years. A much lower initial value \( \gamma_w \leq 10^4 \) for Crab pulsar is inferred by Atoyan (1999) when interpreting the radio spectrum of the Crab nebula. In light of these factors, we regard that the value of \( \gamma_w \) is quite uncertain and may span a range of two orders. The wind luminosity consists of electromagnetic and particle contributions with their flux ratio denoted by the magnetization parameter \( \sigma \). KC84 note that \( \sigma \) must be smaller than 0.1 in order to have a high efficiency of the wind power into synchrotron radiation. We assume that the pulsar wind flux is dominated by \( e^+e^- \) pairs, so \( L_w \approx 4\pi n_1 \gamma_w^2 R_s m_e c^3 \), where \( n_1 \) is the proper pair density in the wind just before the shock front.

As the young SNR expands at a velocity \( v_b \ll c \), the wind energy excluding the part that goes into radiative and adiabatic loss will accumulate within a volume confined by SNR. The characteristic radius \( R_s \) of the shock front is determined by the balance between the ram pressure of the wind \( L_w/4\pi c R_s^2 \) and the total magnetic and particle pressure \( P \) in the bubble. Since the sound speed in the shocked fluid \( c/\sqrt{\gamma} \) is much larger than that of the SN ejecta, no pressure gradient in the nebular bubble can be maintained for dynamically important time, and the bubble is virtually isobaric (Reynolds & Chevalier 1984). At the shock front \( R_s \), the wind \( e^+e^- \) particles may acquire a power-law energy distribution of the form \( N(\gamma) \propto \gamma^{-p} \) for \( \gamma \geq \gamma_m \), where \( n = \int N(\gamma) d\gamma \) is the number density of the shocked relativistic particles, \( \gamma \) is the Lorentz factor of the particles and \( \gamma_m \) is the minimum value of \( \gamma \). Usually \( p > 2 \) is assumed to keep the total energy of the shocked particles finite and \( p = 2.2 \) is inferred for Crab nebula from its X-ray spectrum (e.g. Pravdo & Serlemitsos 1981). \( p \approx 2.2 \) seems also consistent with the simulation result of the first-order Fermi particle acceleration (Kirk et al. 2000). From the shock jump conditions, the internal pressure immediately behind the shock front is \( p_b = n_1 m_e c^2 \gamma_w^2 \) and the energy density is \( e_b = 3 n_1 m_e c^2 \gamma_w^2 \). As in Chevalier (2000), we assume that the energy density in the downstream region is divided between a
fraction $\epsilon_e$ into particles and a fraction $\epsilon_B = 1 - \epsilon_e$ into the magnetic field, similar to the definition in the blast-wave model of GRB afterglows (Sari et al. 1998). The plausibility for incorporating the equipartition assumption have been accounted for in Chevalier (2000) and KG.

With these assumptions, the magnetic field in the bubble is

$$B^b = \left(\frac{6\epsilon_BL_w}{R_s^2c}\right)^{1/2} = 0.17 \text{Gauss}$$

where $\alpha \equiv R_s/R_b$ and $\epsilon_B = 0.5\epsilon_{B,0.5}$. The minimum Lorentz factor of the shocked $e^+e^-$ is $\gamma_m = (p-2)/(p-1)\epsilon_e\gamma_w = 8 \times 10^3\epsilon_{e,0.5}\gamma_{w,5}$, where $\epsilon_e = 0.5\epsilon_{e,0.5}$ and $\gamma_w = 10^5\gamma_{w,5}$. The synchrotron lifetime of the shocked $e^+e^-$ with the average Lorentz factor $\bar{\gamma}_e = \epsilon_e\gamma_w$ is estimated to be

$$\tau_{\text{syn}} = \frac{24\pi m_e c}{\sigma_T \bar{\gamma}_e (B^b)^2} = 2 \times 10^6 \text{sec}$$

where $\sigma_T$ is the Thompson cross section. Thus, for $t_{sd} \lesssim 10\text{yr}$ and $\gamma_w$ ranging from $10^4$ to $10^6$, the synchrotron lifetime of the $e^+e^-$ in the bubble is much shorter than the age of the bubble and the PWB is in the strongly cooling regime for parameters of interest. In this case, the radial width of the bubble $\Delta_b$ is much smaller than the radius of the bubble $R_b$, which means $\alpha \sim 1$, and we can express the bubble volume as $V = 4\pi R_b^2\Delta_b$. Now let’s estimate the radial width $\Delta_b$ at the time ($t_{sd}$) when GRB-PWB interaction begins. Equating the ram pressure with the internal pressure of the bubble, we have

$$\frac{1}{3} \left[ \frac{(B^b)^2}{8\pi} + \epsilon_p \right] = \frac{L_w}{4\pi c R_s^2},$$

where $\epsilon_p$ is the internal energy density of the particles. The particle energy density can be regarded as the sum of the contributions from both the $e^+e^-$ that have suffered strongly cooling and those that have not suffered any cooling yet, i.e.

$$\epsilon_p = \left( \int_{0}^{t_{sd}-\tau_{\text{syn}}} \epsilon_e L_w \frac{\gamma^b_e}{\gamma_{e}} dt + \int_{t_{sd}-\tau_{\text{syn}}}^{t_{sd}} \epsilon_e L_w dt \right) / 4\pi R_b^2 \Delta_b,$$

where $\gamma^b_e$ is the Lorentz factor of cooled $e^+e^-$ in the bubble at time $t_{sd}$: $\gamma^b_e = 6\pi m_e c / (\sigma_T (B^b)^2 (t_{sd} - t))$. Finally we get

$$\Delta_b = \frac{6\pi m_e c}{\sigma_T (B^b)^2 \epsilon_B \gamma_{w}} \ln \frac{t_{sd}}{\tau_{\text{syn}}} + \tau_{\text{syn}} \bar{\gamma}_e \frac{c}{3} \simeq 3 \times 10^{16} \text{cm} \gamma_{w,5}^{-1} \epsilon_{e,0.5}^{-1} \epsilon_{B,0.5}^{-1} t_{sd,1}^3 \Delta E_{\text{rot,53}}^{-1} \epsilon_{b,0.1}^2,$$

consistent with our previous assumption $\Delta_b \ll R_b$.

The innermost part of the bubble is composed of newly injected pairs from the pulsar that have suffered negligible cooling and the minimum Lorentz factor is $\gamma^m_m$. For a nonmagnetic ultrarelativistic shock, the downstream fluid just behind the shock front moves with
a velocity $v_d \approx \frac{1}{3}c$ relative to the shock front. The bulk of the shocked $e^+e^-$ may move at a speed between $v_d$ and the bubble expanding speed $v_b$. As a rough estimate, the width of these 'uncooled' $e^+e^-$ region is $\Delta_1 \approx \bar{v} \tau_{\text{syn}} \approx 10^{16}\text{cm}\gamma_{w,5}^{-1}\epsilon_{e,0.5}^{-1}\epsilon_{B,0.5}^{-1}t_{sd,1}^2\Delta E_{\text{rot},53}^{-1}v_{b,0.1}^2$, where we have taken $\bar{v}$ being $0.2c$.

3 GRB-PWB INTERACTION AND EARLY GEV AFTERGLOWS

3.1 Early GeV afterglows and GRB940217

At the time $t_{sd}$ after the SN explosion, the pulsar collapses to a black hole or converts to a strange star and the GRB goes off, sending a fireball and a relativistic blast wave into the PWB. Different from the usual cold proton-electron preshock medium in the standard model of afterglows, the PWB is hot, composed of randomly moving, relativistic $e^+e^-$ pairs, and highly magnetized. The initially freely-expanding fireball shell will be decelerated by the swept-up hot $e^+e^-$ pairs. We can calculate the extent to which the shell is decelerated by PWB gas. The total enthalpy, representing the effective inertia that decelerates the shell, is $U \sim 4\pi R_b^2 \Delta_b (e_b + p_b) \approx 1.2 \times 10^{51}\text{erg}\gamma_{w,5}^{-1}\epsilon_{e,0.5}^{-1}\epsilon_{B,0.5}^{-1}t_{sd,1}^2\Delta E_{\text{rot},53}$. If we temporarily neglect the radiative loss of the GRB shell (i.e. the shell energy is conserved), its Lorentz factor will be reduced from the initial value $\Gamma_0$ to $\Gamma_f = (E_0^s/U)^{1/2} \approx 10(E_{0,53}^s)^{1/2}\gamma_{w,5}^{1/2}\epsilon_{e,0.5}^{1/2}\epsilon_{B,0.5}^{1/2}t_{sd,1}^{-1}\Delta E_{\text{rot},53}^{-1/2}$ when the shell collides with the outer supernova remnant, where $E_0^s = 10^{53}\text{erg}E_{0,53}^s$ is the initial isotropic kinetic energy of the GRB shell. The deceleration of the GRB shell starts when the energy of the swept-up medium equals to $E_0^s/\Gamma_0^2$, lagging the GRB prompt emission by a time $\Delta t_{lag} = R_b/2\Gamma_0^2c = 180s t_{sd,1}v_{b,0.1}\Gamma_0^{-2} \Gamma_{0,300}$, where $\Gamma_0 = 300\Gamma_{0,300}$ is the initial Lorentz factor. As the forward shock expanding into the hot PWB, the average Lorentz factor per GRB-shocked $e^+e^-$ particle is $\bar{\gamma}_e^s = \epsilon_e\gamma_w \Gamma = 1.5 \times 10^7\epsilon_{e,0.5}\gamma_{w,5}\Gamma_{300}$, if the post GRB-shocked medium has the same equipartition factors $\epsilon_e, \epsilon_B$ as those in the pulsar wind shocks. According to the shock jump conditions, the minimum Lorentz factor of the GRB-shocked $e^+e^-$ is

$$\gamma_m^s = \frac{p-2}{p-1}\epsilon_e\gamma_w \Gamma = 2.5 \times 10^6\epsilon_{e,0.5}\gamma_{w,5}\Gamma_{300}$$

for $p = 2.2$ and the magnetic field in the post GRB-shocked medium is

$$B^s = \Gamma B^b = 51\text{Gauss}\Gamma_{300}\epsilon_{B,0.5}^{3/2}t_{sd,1}^{-3/2}\Delta E_{\text{rot},53}^{1/2}v_{b,0.1}^{-1}.$$
loss time. Taking the shock-acceleration time as the gyroperiod and assuming synchrotron cooling, we get

$$\gamma_M^s \simeq \left( \frac{3\pi q_e}{\sigma_T B^s} \right)^{1/2} = 1.9 \times 10^7 \Gamma_{300}^{-1/2} \epsilon_{0.5}^{-1/4} \epsilon_{0.5}^{3/4} \Delta E_{\text{rot},53}^{-1/4} \Gamma_{0.1}^{1/2}, \tag{9}$$

where $q_e$ is the electron charge. The observed characteristic synchrotron emission photon energy for $e^+e^-$ with the Lorentz factor $\gamma_M^s$ is

$$h\nu^s = \Gamma (\gamma_M^s)^2 \frac{q_e B^s}{2\pi m_e c} = 1.1 \text{GeV} \Gamma_{300}^4 \epsilon_{e,5}^2 \epsilon_{0.5} \Delta E_{\text{rot},53}^{-1/2} \Delta E_{\text{rot},53}^{1/2} \tag{10}$$

and the maximum synchrotron emission photon energy is

$$h\nu^s_M = \frac{3q_e^2}{m_e c \sigma_T} \Gamma = 45 \text{GeV} \Gamma_{300}, \tag{11}$$

with the spectrum being $F_\nu \propto \nu^{-p/2}$ between them.\n
When the GRB blast wave swept-up matter cools, an initial GeV burst occurs at the deceleration radius. The deceleration of the GRB shell starts when the energy of the swept-up medium equals to $E_0^s/\Gamma_0^2$, and the corresponding radial distance that the shell travelled in the PWB is $x_{\text{dec}} = (E_0^s/\Gamma_0^2 \epsilon R_0^2 (e_b + p_b)) = 3 \times 10^{13} \text{cm} \Gamma_{0.300}^{-2} \epsilon_{0.5}^{-2} \pi R_{13}^3$. Please note that, as the shell slows down in the PWB, the radial distance $x$ that the shell travelled is much shorter compared to its radius. In this case, the angular spreading time scale\footnote{Note that for a range of parameter values, the assumed $x_{\text{dec}}$ may become larger than $\gamma_M^s$, then one may ask what will happen. To answer this, one must know the detailed microphysics of how the shock is formed. If there is a process which thermalises particles faster than the gyroperiod, then this process could lead to higher energy ‘thermalised’ particles rather than the ‘nonthermal’ maximum produced by the standard Fermi acceleration picture (J. G. Kirk 2001, private communication). In this case, the Fermi picture would not be relevant. If there is no such a process, then the Lorentz factor of the GRB-shocked $e^+e^-$ pairs implied by the jump conditions cannot be reached. The energy must be radiated away within the shock structure itself (by synchrotron radiation) and the jump conditions must be modified to account for the energy ‘leak’ (J. G. Kirk 2001, private communication). Thus, strong high-energy emissions are also expected.} of the shell emission $\Delta t_{\text{ang}} = R_b/2\Gamma^2 c$ (Katz 1994a; Fenimore et al. 1996) dominates over the radial time scale $\Delta t_{\text{rad}} = x/2\Gamma^2 c$, Thus an interesting situation arise: though the time that the shell spends on moving through the PWB is very short, the shock generated energy radiates over a much longer observed time which is determined by the shell angular spreading time. Therefore the shell angular spreading time is relevant to the flux and light curve in the following calculation. So, the effective deceleration time should therefore be $t_{\text{dec}} = R_b/2\Gamma^2 c = 180 \sec \epsilon R_{13} / \epsilon_{0.1} t_{13} \Gamma_{0.300}^{-2}$. As the shocked PWB medium is in the fast cooling regime, the power emitted is simply
that given to the shocked $e^+e^-$, that is $\epsilon_e$ times the power generated by the shock, i.e. $\int F_\nu d\nu = \epsilon_e \frac{dE_s}{dt} \frac{1}{4\pi d_L^2}$, where $d_L$ is the luminosity distance of the GRB source. Integrating the fast cooling spectrum with $\nu_m^s \gg \nu_e^s$ over frequency (Sari et al. 1998), we have

$$\int F_\nu d\nu = \nu_m^s F_m^s \left[ 2 + \frac{2}{p-2} \left( 1 - \left( \frac{\gamma_m^s}{\gamma_m^s} \right)^{2-p} \right) \right] = \epsilon_e \frac{dE_s}{dt} \frac{1}{4\pi d_L^2},$$

(12)

where $E^s$ is the isotropic kinetic energy of the decelerated GRB shell. Strictly, to calculate the radiated power at the deceleration radius, we need to calculate the total number of swept-up relativistic electrons and their radiative power as in Dermer et al. 1999. But as a rough estimate, we assume that the bulk of the shell kinetic energy is dissipated into the shock energy at the deceleration radius (Rees & Mészáros 1992; Dermer et al. 1999) and get the fluence at $\nu_m^s \sim 1$GeV band (for $p = 2.2$): $t_{\text{dec}}, \nu F_\nu(1\text{GeV}) = 8 \times 10^{-6}\text{ergcm}^{-2}\epsilon_e, 0.5 E_{53}^s d_{28}^2$. Clearly, this intensity of GeV emission is well above the detectability of EGRET detector. This synchrotron GeV emission will fade with time and to a certain time it will be too weak to be detected, leading to an extended GeV afterglow emission.

Below we will study the fading behavior of GeV afterglow. According to Eq. (12), the observed afterglow flux at $\sim 1$GeV is

$$\nu F_\nu(1\text{GeV}) = \nu_m^s F_m^s \left( \frac{\nu}{\nu_m^s} \right)^{-p/2+1} \propto (\nu_m^s)^{p/2-1} \frac{dE_s}{dt} t_{\text{dec}}.$$  

(13)

The duration of emissions from the shell scales with $\Gamma$ as $t \sim R_b/2\Gamma^2 c \propto \Gamma^{-2}$. As the shocked gas is partially radiative, the shock energy $E^s$ decreases as $E^s \propto \Gamma^\epsilon$ in the limit that $\Gamma_0 \gg \Gamma \gg 1$ (Böttcher & Dermer 2000), where $\epsilon$ is the fraction of the shock-generated thermal energy that is radiated. For a fast-cooling shock, $\epsilon = \epsilon_e$. Considering $\nu_m^s \propto \Gamma^4$, from Eq.(13) we find the GeV spectral power flux decays with time as

$$\nu F_\nu(1\text{GeV}) \simeq 4.4 \times 10^{-8}\text{ergcm}^{-2}\text{s}^{-1} \left( \frac{t}{t_{\text{dec}}} \right)^{-p-1-\epsilon/2} \epsilon_e, 0.5 E_{53}^s d_{28}^2 t_{\text{dec}}^{-1} v_{b,0.1}^{-1} \Gamma_{0.300}^2.$$  

(14)

For comparison with EGRET observation and future GLAST observation, we now calculate the energy fluence the detector can collect during a certain time duration $t$ (see Fig.4 of Zhang & Meszaros 2001), i.e. the product of $\nu F_\nu$ Gev flux and the observing time.

$$t\nu F_\nu(1\text{GeV}) \simeq 8 \times 10^{-6}\text{ergcm}^{-2} \epsilon_e, 0.5 E_{53}^s d_{28}^2 \left( \frac{t}{t_{\text{dec}}} \right)^{-p-2-\epsilon/2}.$$  

(15)

The fluence threshold for EGRET is roughly $\sim 2.1 \times 10^{-6}\text{ergcm}^{-2}$ for short integration time regime (Zhang & Meszaros 2001). According to Eq. (15), the decreasing of the fluence from the initial value at time $t_{\text{dec}}$ to the sensitivity threshold of EGRET implies that the observed GeV emission lasts
for \( p = 2.2 \). We think that this mechanism provides a plausible explanation for the famous long-duration (\( \sim 5400 \) sec) high-energy photons, with energies ranging from 36 Mev to 18 GeV, from GRB940217 detected by EGRET (Hurley et al. 1994). The future GLAST detector is much more sensitive than EGRET and its fluence threshold is roughly \( \sim 1.2 \times 10^{-9}t^{1/2}\text{ergcm}^{-2} \) for a long integration time regime, so a more extended detected GeV emission time is expected with \( t_{\text{dur}} \sim 1.8 \times 10^6 \text{sec} \). The power-law fading behavior of GeV emissions in our scenario distinguishes itself from some other models, such as afterglow inverse-Compton (e.g. Katz 1994b) and hadron process models (e.g. Katz 1994b).

### 3.2 Synchrotron self-Compton and nebula-induced external Compton components at GeV band

We need to consider the synchrotron self-Compton (SSC) and nebula-induced external Compton (EC) components at GeV band in order to compare it with synchrotron component obtained in the above subsection. The GRB-shocked \( e^+e^- \) gas will inverse-Compton scatter synchrotron photons emitted by themselves to higher energies. Because the shock is fast-cooling, the ratio of the SSC to synchrotron luminosity is approximately given by (Sari & Esin 2001)

\[
L_{\text{SSC}} = \begin{cases} 
\frac{\epsilon_e}{\epsilon_B} L_{\text{syn}} & \text{if } \epsilon_e \ll \epsilon_B \\
\left( \frac{\epsilon_e}{\epsilon_B} \right)^{1/2} L_{\text{syn}} & \text{if } \epsilon_e \gg \epsilon_B
\end{cases}
\]

For fast-cooling shocks, \( L_{\text{syn}} \sim 4\pi d_L^2 \nu_m^s F_m^s \) and \( L_{\text{SSC}} \sim 4\pi d_L^2 \nu_m^{\text{SSC}} F_m^{\text{SSC}} \), respectively. Note that \( \frac{\nu_{m}}{m_c} \gg m_e c^2 \), the SSC energy emission peaks at \( \nu_{KN} = \Gamma \gamma^2 m_e c^2 \sim 3\text{TeV} \frac{1}{2} d_0^{-1/4} \gamma^{3/4} \nu_{b,0.1}^{1/2} \Delta E_{\text{rot},53}^{1/4} \), above which the Klein-Nishina effect suppresses Compton scattering. Therefore, at \( \sim 1\text{GeV} \) band, the intensity of the SSC fluence component is generally below that of the synchrotron one so long as \( \epsilon_e \) is not much larger than \( \epsilon_B \), though at TeV band, it dominates over the latter.

Another important process that might contribute to the high-energy component is the Comptonization of external nebular soft photons by relativistic \( e^+e^- \) in the expanding GRB shell. The luminosity of the EC emission can be obtained analogously to the SSC case, by using \( \nu_m^E F_m^E \sim \tau_e \Gamma^2 \gamma^s m c \nu_m^s F_m^b \) where \( \tau_e \) is the optical depth of the shocked material. The
number density of $e^+e^-$ in the bubble is $n_e^b = 3\gamma_w n_1 = 3L_w/4\pi\gamma_w R_b^2 m_e c^3 = 0.03\gamma^{-1}_{w,5} t^{-3}_{sd,1} \Delta E_{\text{rot}} v_{b,0.1}^{-2}$, so $\tau_e \sim 1/3\sigma_T n_e^b \Delta x \sim 10^{-10} \gamma_w^{-1} t^{-3}_{sd,1} \Delta E_{\text{rot}} v_{b,0.1}^{-2} \Delta x_{16}$, where $\Delta x$ is the radial distance the shell travelled in the bubble. The luminosity of the PWB $\nu_m^b F_{\nu_m}^b$ can be obtained from Eq. (9) in Chevalier (2000): $\nu_m^b F_{\nu_m}^b \sim 10^{-14}\text{ergcm}^{-2}\text{s}^{-1}\epsilon_{0.5} \gamma_{w,5}^{0.2} t^{-1}_{sd,1} d_L^{-2}$ for $p = 2.2$. We finally get the EC luminosity at time $t_{\text{dec}}$, $L_{\text{EC}} \sim 1.1 \times 10^{-12}\text{ergcm}^{-2}\text{s}^{-1}$, for typical parameters used.

3.3 Attenuation of high-energy photons by soft photons in PWB

As the PWB has plenty of soft photons, we must consider whether the high-energy photons suffer strong attenuation due to pair-production reaction while propagating through the PWB. The optical depth of a high-energy photon with energy $E_\gamma$ due to $\gamma - \gamma$ absorption writes $\tau_{\gamma\gamma} = \int_{\nu_0}^{\nu_{\text{max}}} \sigma(\nu) n(\nu) \Delta \nu$, where $h\nu_0 \equiv 2(m_e c^2)^2 / E_\gamma$ and $\nu_{\text{max}}$ is the frequency of the maximum energy photon in the bubble, $n(\nu) \simeq L(\nu) / h \nu 4\pi R_b^2 c$ is the number density of photons with energy $h \nu$ ($L(\nu)$ is the luminosity per unit frequency of the PWB), and $\sigma(\nu)$ is the pair production cross section of $e^- e^+$.

The formula of the PWB luminosity $L(\nu)$ has been derived by Chevalier (2000) (see his Eq. (9)), given by

$$L(\nu) = \frac{1}{2} \left( \frac{p-2}{p-1} \right)^{p-1} \frac{64\epsilon_e^2}{(2\pi m_e c)^2} L_{p-2}^{(p-2)/4} R_s^{p-2} R_w^{-2} \nu_{\text{dec}}^{-(p+2)/4} L_{w}^{(p+2)/4} \nu^{-p/2},$$

for $\nu > \nu_m^b > \nu_c^b$ and $L(\nu) \propto \nu^{-2}$ for $\nu_m^b > \nu > \nu_c^b$. Setting $\tau_{\gamma\gamma} = 1$, after a simple algebraic computation, we get $h\nu_0 \simeq 0.18\epsilon_{0.5} \gamma_{w,5}^{0.86} \gamma_{5}^{0.73} t^{-0.14}_{sd,1} \Delta E_{\text{rot,53}}^{-0.045} v_{b,0.1}^{-0.09}$ for $p = 2.2$. Therefore the high-energy cutoffs due to pair production with soft photons in PWB lie at

$$E_{\text{cut}} \simeq 5\text{TeV} \epsilon_{0.5} \epsilon_{B,0.5}^{-0.18} \gamma_{w,5}^{-0.86} \gamma_{5}^{0.73} t^{-0.14}_{sd,1} \Delta E_{\text{rot,53}}^{-0.045} v_{b,0.1}^{-0.09}$$

for $p = 2.2$. So we can safely conclude that the synchrotron GeV photons can reach us without significant attenuation due to pair production through interaction with soft photons in the PWB.

4 SUMMARY

GRBs could arise within PWBs in supernova-like scenarios (Vietri & Stella 1998; Wang et al. 2000) of GRBs, in which SN explosion precedes GRB by a time from several months to

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tens of years and the pulsar emits a relativistic wind, making a wind bubble while shocking against the ambient medium. After the spin-down time, the central pulsar may collapse to a black hole or convert to a strange star, forming the GRB fireball. In this paper, we studied the early afterglow high-energy emissions from GRB blast wave expanding into such a pulsar wind bubble. We find that, owing to that the PWB is composed of hot $e^+e^-$ pairs with high randomly-moving Lorentz factors, the majority of the GRB-shocked PWB pairs have such high thermal Lorentz factors that their characteristic synchrotron frequencies lie at GeV bands. Our calculation shows that the attenuation of GeV photons by soft photons in PWB due to pair-production reaction is negligible. Strong power-law decaying GeV emissions during the early afterglow phase is naturally expected in this scenario and it may provide a plausible explanation for the long-time ($\sim 5400$ s) GeV photons from GRB940217. The predicted power-law decaying behavior of the flux is distinguished from the afterglow electron-IC mechanism (Zhang & Mészáros 2001) as the latter predicts a delayed hump in the light curves.

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