Mon. Not. R. Astron. Soc. 389, 1306–1310 (2008) doi:10.1111/j.1365-2966.2008.13632.x

Cosmic ray protons in the energy range 10^{16}–$10^{18.5}$ eV: stochastic gyroresonant acceleration in hypernova shocks?

Yi-Zhong Fan[1,2]⋆

[1]*Niels Bohr International Academy, Niels Bohr Institute, Copenhagen University, Blegdamsvej 17, DK-2100 Copenhagen, Denmark*
[2]*Purple Mountain Observatory, Chinese Academy of Sciences, Nanjing 210008, China*

Accepted 2008 June 23. Received 2008 June 23; in original form 2007 December 5

ABSTRACT

The hypernovae (HNe) associated with gamma-ray bursts (GRBs) may have a fairly steep energy–velocity distribution, i.e. $E(\geq\beta) \propto \beta^{-q}$ for $q < 2$ and $\beta \geq \beta_o$, where β is the velocity of the material and $\beta_o \sim 0.1$ is the velocity of the slowest ejecta of the HN explosion, both in units of the speed of light (c). The cosmic ray protons above the second knee but below the ankle may be accelerated by the HN shocks in the velocity range of $\beta \sim \beta_o$–$4\beta_o$. When $\beta \leq 4\beta_o$, the radius of the shock front to the central engine is very large and the medium decelerating the HN outflow is very likely to be homogeneous. With this argument, we show that for $q \sim 1.7$, as inferred from the optical modelling of SN 2003lw, the stochastic gyroresonant acceleration model can account for the spectrum change of high-energy protons around the second knee. The self-magnetized shock acceleration model, however, yields too steep a spectrum which is inconsistent with the observation unless the medium surrounding the HN is a free wind holding up to a radius ∼1–10 kpc.

Key words: acceleration of particles – supernovae: general – cosmic rays – supernova remnants – gamma-rays: bursts.

1 INTRODUCTION

One of the most widely suggested sources of cosmic rays (CRs) is supernova (SN) remnants (see Hillas 2005 for a review). Since 2001, more and more researchers have noted that hypernovae (HNe), in particular those associated with gamma-ray bursts (GRBs), may play an important role in interpreting the CR spectrum above its first knee,[1] i.e. $\sim 3 \times 10^{15}$ eV (Dermer 2001a,b; Erlykin, Wibig & Wolfendale 2001; Sveshnikova 2003; Wick, Dermer & Atoyan 2004; Wang et al. 2007; Budnik et al. 2008). This is reasonable since the average velocity and the total kinetic energy of hypernova (HN) outflows are much larger than those of normal SNe; so is the energy of accelerated particles. A reliable interpretation of the CR spectrum up to $\sim 10^{18}$ eV thus should take into account the different energies and types among SNe (Sveshnikova 2003).

How to accelerate protons up to an energy $\sim 10^{18.5}$ eV in HN blast waves? Dermer (2001b) suggested that gyroresonant stochastic acceleration might play such a role (see fig. 10 of Dermer 2001a for a quantitative plot). Other authors (Erlykin et al. 2001; Wang et al. 2007; Budnik et al. 2008) considered the self-magnetized acceleration model put forward by Bell & Lucek (2001), in which the magnetic field of the upstream region has been significantly amplified by CRs. Considering the energy distribution of the HN outflow,[2] Wang et al. (2007) and Budnik et al. (2008) suggested that with the second model the CR proton spectrum steepening around the second (first) knee could be reproduced. In this work, we point out one potential limit of such an interpretation and show that the gyroresonant stochastic acceleration model does not suffer from that problem.

This paper is arranged as follows. In Section 2, we discuss the energy–velocity distribution of HN outflows and the medium profile surrounding the HN outflows. We find that for HN outflows with a fairly steep energy–velocity distribution, when $\beta_o \leq \beta \leq 4\beta_o$ that may play the main role in accelerating the CR protons above the second knee but below the ankle (see the discussions below equations 8 and 16), the radius of the shock front to the central engine is very large and the medium decelerating the HN outflow is likely to be homogeneous. In Section 3, we calculate the change of the CR spectrum around the second knee which is caused by the energy–velocity distribution of HN outflows, and compare the results with the CR spectrum observation so as to constrain the models. In Section 4, we summarize our results with some discussions.

⋆E-mail: yzfan@pmo.ac.cn

[1] The second knee and the ankle in the CR spectrum are at $\sim 3 \times 10^{17}$ and $\sim 3 \times 10^{18}$ eV, respectively.

[2] See Berezhko & Völk (2004) and Ptuskin & Zirakashvili (2005) for the influence of the energy distribution of normal SN outflows on the spectrum of accelerated CRs.

2 THE HYPERNOVA OUTFLOW: ENERGY–VELOCITY DISTRIBUTION AND THE MEDIUM INTO WHICH IT EXPANDS

2.1 Energy–velocity distribution of hypernova outflows

HNe, especially those associated with GRBs/X-ray flashes (XRFs), are distinguished for the broad lines in their spectra, indicating very high expansion velocity of the ejecta. The modelling of optical light curves and spectra, in principle, can reconstruct the energy–velocity distribution of the outflows. However, no reliable constraint can be given on the $\beta > 0.3$ outflow by optical data even if that part had some optical depth since the current optical modelling is not fully relativistic (Deng, private communication). In SN 1998bw, SN 2003dh, SN 2003lw and SN 2006aj, strong photospheric velocity evolution is evident (Hjorth et al. 2003; Mazzali et al. 2006; Soderberg et al. 2006). The optical modelling of SN 2003lw showed that material moving faster than $(0.1, 0.2)c$ was $\sim(1.4, 0.1)\,\mathrm{M}_\odot$, respectively (Mazzali et al. 2006), implying a fairly steep initial kinetic energy distribution $E(\geq\Gamma\beta) \propto (\Gamma\beta)^{-1.7}$, where $\Gamma = (1-\beta^2)^{-1/2}$. But for other events, no result has been published. Soderberg et al. (2006) constrained the kinetic energy profile of HN outflows in a more speculative way. They used optical spectral data to probe the slowest ejecta in supernova explosions and employed radio observations to trace the fastest component of the outflow. They then took these two data points to estimate the energy–velocity distribution. Their results may be biased because the fast moving material identified by radio observations might be the decelerated GRB/XRF ejecta rather than the fastest component of the main SN explosion. If so, it is not a continuous distribution of matter between the two data points (Soderberg et al. 2006; Xu, Zou & Fan 2008).

Fairly speaking, observationally so far we do not have a reliable estimate of the initial kinetic energy–velocity distribution of (most) HN outflows in the velocity range of $\beta \sim 0.1$–0.5. Theoretically, the standard hydrodynamic collapse of a massive star (Tan, Matzner & McKee 2001) results in a kinetic energy profile of the SN explosion $E(\geq\Gamma\beta) \propto (\Gamma\beta)^{-5.2}$. Such a steep function, however, is inconsistent with the constraint from the optical data of SN 2003lw (Mazzali et al. 2006), for which a rough estimate gives $E(\geq\Gamma\beta) \propto (\Gamma\beta)^{-1.7}$. Motivated by this fact, we *assume* that all HNe associated with GRBs have a fairly steep energy distribution, which is generally written as $E(>\Gamma\beta) = A[\Gamma\beta/(\Gamma_\mathrm{o}\beta_\mathrm{o})]^{-q}$ for $\beta < 0.5$, where $\Gamma_\mathrm{o} = (1-\beta_\mathrm{o}^2)^{-1/2}$. For SN 1998bw/GRB 980425, SN 2003dh/GRB 030329, SN 2003lw/GRB 031203 and SN 2006aj/GRB 060218, optical modelling suggests $A \sim 0.2$–6×10^{52} erg and $\Gamma_\mathrm{o}\beta_\mathrm{o} \sim 0.04$–$0.1$ (Soderberg et al. 2006). The parameter q, however, is not reliably determined in most cases. For simplicity, we approximate $E(>\Gamma\beta) = A[\Gamma\beta/(\Gamma_\mathrm{o}\beta_\mathrm{o})]^{-q}$ as $E(>\beta) = A(\beta/\beta_\mathrm{o})^{-q}$ for $\Gamma \sim 1$.

In an explosion, the outmost, also the fastest, part of the SN outflow interacts with the medium first. When the fast component is decelerated by the medium, the slower part will catch up with the decelerating shock front. As a result, the total kinetic energy of the shocked medium increases and the deceleration of the shock is suppressed. In the quasi-similar evolution phase of the HN shock, the fastest component has swept enough medium and has got decelerated. A significant part of the initial kinetic energy of the HN material $E(\geq\beta)$ has been used to accelerate the medium to a velocity $\sim\beta$. So when we talk about the CR acceleration in the blast wave, the $E(\geq\beta)$ mentioned there actually represents the total kinetic energy of the shocked medium moving with a velocity β. For a medium taking the profile $n \propto R^{-k}$ $(0 \leq k \leq 3)$, the rest mass swept by the HN blast wave is $M_\mathrm{med} = \int_0^R 4\pi n m_\mathrm{p} R^2 \mathrm{d}R \propto R^{3-k}$. For $\beta > \beta_\mathrm{o}$, conservation of energy gives $E(\geq\beta) \approx M_\mathrm{med}\,\beta^2/2$, i.e.

$$\beta^{-(q+2)} \propto R^{3-k}. \tag{1}$$

With the relation that $R \sim \beta t$, the dynamics of the HN outflow is described by (for $\beta > \beta_\mathrm{o}$)

$$\beta \propto t^{-\frac{3-k}{5+q-k}}. \tag{2}$$

In the next section, we will show how q and k influence the CR spectrum.

2.2 The medium into which the hypernova outflow expands

As shown in equation (2), the dynamics of a HN shock depends on the medium profile sensitively. Here, we review the medium profiles of all four GRB-HN events, based on the GRB and/or HN afterglow modelling. For GRB 980425, the medium is wind-like (i.e. $k \sim 2$). The afterglow modelling favours an unusual small $A_* \sim 0.01$–0.04 (Li & Chevalier 1999; Waxman 2004), where $A_* \equiv (\dot{M}/10^{-5}\,\mathrm{M}_\odot\,\mathrm{yr}^{-1})(v_\mathrm{w}/10^8\,\mathrm{cm\,s}^{-1})^{-1}$, $\dot{M}$ is the mass-loss rate of the progenitor and v_w is the velocity of the stellar wind. For GRB 030329, the circumburst medium is found to be homogeneous (i.e. $k \sim 0$), as shown in many independent investigations (Frail et al. 2005; Pihlström et al. 2007; van der Horst et al. 2008; Xue et al. 2008). For GRB 031203, after modelling the radio data, Soderberg et al. (2004) got a constant $n \sim 0.6\,\mathrm{cm}^{-3}$ (cf. Ramirez-Ruiz et al. 2005). For GRB 060218, the high-quality radio data support the homogeneous medium model with $n \sim 100\,\mathrm{cm}^{-3}$ (Fan, Piran & Xu 2006; Soderberg et al. 2006). As such, we have no compelling evidence for a wind-like medium surrounding most GRBs, even for those associated with HNe. The physical reason is not clear, yet. A post-common envelope binary merger model (e.g. Fryer, Rockefeller & Young 2006) or a fast motion of the Wolf–Rayet star relative to the interstellar medium (ISM) (van Marle et al. 2006) may be able to solve this puzzle.

Actually, a free wind medium, supposed to surround the progenitor, is unlikely to be able to keep such a profile up to the radius:

$$R_\mathrm{dec}(\beta_\mathrm{o}) \sim 3 \times 10^{22} \left(\frac{M_\mathrm{ej}}{10\,\mathrm{M}_\odot}\right) A_{*,-1}^{-1}\,\mathrm{cm}, \tag{3}$$

where M_ej is the rest mass of the GRB-associated HN ejecta and R_dec is the deceleration radius. This is because during their evolution, massive stars lose a major fraction of their mass in the form of a stellar wind. The interaction between this stellar wind and the surrounding ISM creates a circumstellar bubble (e.g. Wijers 2001; Ramirez-Ruiz et al. 2001; Dai & Wu 2003; Chevalier, Li & Fransson 2004; van Marle et al. 2006). The analytical calculation suggests that the free wind of a Wolf–Rayet star usually terminates at (Chevalier 2007)

$$R_\mathrm{t} = 5.7 \times 10^{18} \left(\frac{v_\mathrm{w}}{10^3\,\mathrm{km\,s}^{-1}}\right)\left(\frac{p/k}{10^5\,\mathrm{cm}^3\,\mathrm{K}}\right)^{1/2} A_{*,-1}^{1/2}\,\mathrm{cm}, \tag{4}$$

where p is the pressure in the shocked wind and k is the Boltzmann constant. This is confirmed by observations of Wolf–Rayet nebulae, such as NGC 6888 and RCW 58, which also have radii of the order of a few pc (Gruendl et al. 2000). Here, we take the numerical example given in fig. 1 of Chevalier et al. (2004) to show that the HN outflow is mainly decelerated in the ISM-like medium region. In their numerical example, $n \sim 0.5R_{18}^{-2}\,\mathrm{cm}^{-3}$ for $R_{18} < 1.2$. The total mass of the free wind medium is thus $\sim 7 \times 10^{-3}\,\mathrm{M}_\odot \ll M_\mathrm{ej}$. Here and throughout this work, the convention $Q_\mathrm{x} = Q/10^\mathrm{x}$ has been adopted in cgs units.

We therefore conclude that the medium is most likely to be ISM-like at the radius where the HN outflow has been decelerated to $\beta < 4\beta_o$. This can also be understood as follows. One can infer from equation (1) that for $k = 2$ the outflow component with $\beta > \beta_o$ will decelerate at a radius $R_{dec}(\beta) \approx (\beta/\beta_o)^{-(2+q)} R_{dec}(\beta_o)$. So for $\beta \sim 4\beta_o$ and $q \sim 2$,

$$R_{dec}(4\beta_o) \sim 4 \times 10^{-3} R_{dec}(\beta_o) \sim R_t.$$

3 SPECTRUM OF COSMIC RAY PROTONS: OBSERVATION AND INTERPRETATION

3.1 A new CR proton component in the energy range of 10^{16}–$10^{18.5}$ eV

The spectrum of protons steepens suddenly at the first knee by a factor of

$$\Delta\gamma(\mathrm{I}) \sim -2.1.$$

In view that the spectra of heavier particles would steepen at higher energies, the likely interpretation of the steepening of all CRs at the first knee is the sudden decline of the light particles such as H and He (see Hillas 2005; Hörandel 2008, and the references therein).

The proton CR spectrum before and after the second knee, after subtracting the modelled ‘Galactic’ component, can be roughly estimated as (Ulrich et al. 2004; Antoni et al. 2005; Hillas 2005; Hörandel 2008)

$$\frac{dN_{CR}}{dE_{CR}} \propto \begin{cases} E_{CR}^{-2.4} & \text{for } 0.1 < E_{CR,17} < 3, \\ E_{CR}^{-3.3} & \text{for } 3 < E_{CR,17} < 30, \end{cases} \tag{5}$$

which indicates the factor of spectral steepening is $\Delta\gamma(\mathrm{II}) \sim -0.9$. Above the ankle, the CR spectrum changes to $E_{CR}^{-2.7}$, so the factor of flattening is $\Delta\gamma(\mathrm{III}) \sim 0.6$.

The interpretations of spectral changes at the second knee and at the ankle are much less clear. Hillas (2005) interpreted them as a result of an extragalactic component with a spectrum $\propto E_{CR}^{-2.3}$ suffering losses by the interaction between cosmological microwave background radiation and starlight. In this work, we consider that the detected spectral change around the second knee is due to the energy–velocity distribution of HN outflows.

3.2 Theoretical interpretation

3.2.1 Self-magnetized shock acceleration model

In this model, the magnetic field of the upstream region is assumed to be amplified significantly by the CRs themselves (e.g. Bell & Lucek 2001).

With equation (2), we have the radius of the forward shock front as $R \propto t^{\frac{2+q}{5+q-k}}$. The maximum energy accelerated by the forward shock can be estimated by (Bell & Lucek 2001; Berezhko & Völk 2004; Ptuskin & Zirakashvili 2005)

$$E_{max}(k, q) \sim Z\beta eBR \propto t^{\frac{k-4+q(2-k)/2}{5+q-k}} \propto \beta^{\frac{4-k-q(2-k)/2}{3-k}}, \tag{6}$$

where $B \propto \beta R^{-k/2}$ is the magnetic field in the upstream of the shock, which is of the same order as that of the shocked medium.

In the ISM case (i.e. $k = 0$), we have

$$E_{max}(0, q) \propto \beta^{(4-q)/3}, \tag{7}$$

while in the wind case (i.e. $k = 2$),

$$E_{max}(2, q) \propto \beta^2. \tag{8}$$

Here, we do not present the numerical coefficient of $E_{max}(k, q)$ because Wang et al. (2007) and Budnik et al. (2008) have already shown that for typical parameters, $\beta \sim \beta_o \sim 0.1$ is high enough to accelerate protons up to $\sim 10^{17}$ eV regardless of k. In a stellar wind medium, the HN shock front with $\beta \sim 4\beta_o$ can accelerate protons up to $\sim 3 \times 10^{18}$ eV. Therefore, the CR protons above the second knee but below the ankle are mainly accelerated by the HN shock in the velocity range of $\sim\beta_o$–$4\beta_o$.

To get an estimate of the spectrum of the accelerated particles, following Berezhko & Völk (2004) and Ptuskin & Zirakashvili (2005) we assume: (i) the particles with an energy E_{max} escape the shock immediately; (ii) the total energy of the accelerated particles at an energy $E_{CR} = E_{max}(\beta)$ is proportional to $E(\geq\beta)$. In view of the relations $E(\geq\beta) \propto [E_{max}(\beta)]^{-3q/(4-q)}$ for $k = 0$ and $E(\geq\beta) \propto [E_{max}(\beta)]^{-q/2}$ for $k = 2$, we have

$$\frac{dN}{dE_{CR}} \propto \begin{cases} E_{CR}^{-(2+\delta)-\frac{3q}{4-q}} & \text{for } k = 0, \\ E_{CR}^{-(2+\delta)-\frac{q}{2}} & \text{for } k = 2, \end{cases} \tag{9}$$

where $\delta \sim 0.4$ is introduced to account for the proton spectrum in the energy range of 10^{16}–3×10^{17} eV. As $\beta \leq \beta_o, E(\geq\beta) \propto \beta^0$ if the energy loss of the HN shock is ignorable. The accelerated proton spectrum should be $\propto E_{CR}^{-(2+\delta)}$. This explains why there is a spectrum change around the second knee if $E_{max}(\beta_o) \sim 3 \times 10^{17}$ eV.

With $\delta = 0$, to match the detected proton spectrum $dN/dE_{CR} \propto E_{CR}^{-3.3}$, one has to have $q \sim 2.6$, which is very close to that of SN 2003lw and SN 1998bw reported in Soderberg et al.[3] (2006). Therefore, Wang et al. (2007) concluded that the self-magnetized shock acceleration model could account for the spectrum data. However, a few puzzles have to be solved before accepting this argument: (i) if $\delta = 0$, some novel effects are needed to interpret why the proton spectrum departs from E_{CR}^{-2} significantly in the 10^{16}–3×10^{17} eV range. The authors also need to explain why these effects, if any, disappeared in the $3 \times (10^{17}$–$10^{18})$ eV range. (ii) A wind profile holding to a radius $\sim$1–10 kpc is crucial for their argument. If the medium is ISM-like when the outflow gets decelerated to $\beta < 0.4$, Wang et al. (2007)'s approach would yield a spectrum

$$\frac{dN}{dE_{CR}} \propto E_{CR}^{-5} \tag{10}$$

for $q \sim 2$, which is too steep to be consistent with the data. We take this puzzle as a potential limit of their interpretation.

Let us investigate whether a specific wind bubble can solve this puzzle. We assume that the free wind profile is terminated at a radius $\sim R_t \sim 10^{19}$ cm and is followed by an ISM-like shell. Suppose that the shell is so massive that the deceleration of the whole HN outflow occurs at $R \sim R_t \sim$ constant, we have $E_{max} \sim Z\beta eBR \propto \beta B$. If the shell is not dense enough to form a strong reverse shock, i.e. the forward shock velocity decreases continually rather than abruptly, then $B \propto \beta n^{1/2}$. As a result, we have $E_{max} \propto \beta^2 n^{1/2}$ and $dN/dE_{CR} \propto E_{CR}^{-(2+\delta)-\frac{q}{2}}$, provided that the CR protons in the energy range of $\sim 3 \times (10^{17}, 10^{18})$ eV are mainly accelerated in the shocked shell. Though such a possibility is attractive, the request that the reverse shock does not form is hard to satisfy. This is because at a radius $\sim R_t \sim 10^{19}$ cm, the number density of the wind medium $n_w \sim 3 \times 10^{-4}\ \mathrm{cm^{-3}}\ A_{*,-1}\ R_{t,19}^{-2}$. On the other hand, the assumption that $4\pi R_t^3 n_t m_p \sim M_{ej}$ requires that $n_t \sim 1\ \mathrm{cm^{-3}}\ (M_{ej}/10\,\mathrm{M_\odot}) R_{t,19}^{-3}$.

[3] Please see Section 2.1 for the discussion of uncertainty of the q obtained by their method.

So we have a density contrast $n_t/n_w \sim 10^3$. The forward shock expanding into the dense shell will have a pressure $\sim\beta^2 n_t m_p c^2/3$, which is much higher than that of the shocked wind medium ($\sim\beta^2 n_w m_p c^2/3$). A pressure balance will be established by a strong reverse shock penetrating into the shocked wind medium. Therefore, the forward shock velocity is much smaller than $\beta\mid_{\rm shocked\ wind\ medium}$ and cannot accelerate protons to an energy $\sim 10^{18}$ eV. The reverse shock with a velocity $\beta_r \approx \beta\mid_{\rm shocked\ wind\ medium}$ plausibly plays a more important role in accelerating high-energy CR protons. The shocked wind medium has only a very small mass (relative to $M_{\rm ej}$). The reverse shock gets weak after penetrating into the dense HN outflow which has a density comparable to n_t. Then, the forward shock velocity increases and significant CR acceleration in the forward shock front is possible. A detailed numerical calculation, like Ptuskin & Zirakashvili's (2005), is needed to draw further conclusions.

3.2.2 *Gyroresonant stochastic acceleration model*

The maximum energy-gain rate due to the stochastic Fermi acceleration for a marginally relativistic shock can be estimated as (Dermer 2001b)

$$\frac{dE_{\rm CR}}{dR} \approx \frac{\varepsilon_{\rm turb}(\upsilon-1)}{2^{3/2}} ZeB_*\beta^2 \left(\frac{2^{1/2}E_{\rm CR}}{ZeB_* f_\Delta R\beta}\right)^{\upsilon-1}, \tag{11}$$

where Z is the atomic number, $\varepsilon_{\rm turb}$ is the ratio of plasma turbulence to the shock energy density, $B_* \approx 0.4\, n^{1/2}\varepsilon_B^{1/2}$ Gauss, $f_\Delta \sim 1/12$ is the ratio of the width of the swept medium by the shock to R (Dermer & Humi 2001) and υ is the spectrum index of the turbulence ($\upsilon = 5/3$ for Kolmogorov turbulence and 3/2 for Kraichnan turbulence).

Dermer (2001b) took a $\beta \sim$ constant, integrated equation (11) over R, then got $E_{\max}(R)$. However, currently β evolves with R. As shown below, the smaller the radius, the larger the β and the higher the $E_{\max}$. Very energetic CRs can be accelerated at early times but cannot be accelerated continually because of the adiabatic cooling. Taking into account the adiabatic cooling effect, equation (11) takes the new form

$$\frac{dE_{\rm CR}}{dR} \approx \frac{\varepsilon_{\rm turb}(\upsilon-1)}{2^{3/2}} ZeB_*\beta^2 \left(\frac{2^{1/2}E_{\rm CR}}{ZeB_* f_\Delta R\beta}\right)^{\upsilon-1} - \frac{E_{\rm CR}}{R}. \tag{12}$$

Now $E_{\max}$ can be estimated by setting $\frac{dE_{\rm CR}}{dR}=0$, then we have

$$E_{\max} \approx \left[\frac{\varepsilon_{\rm turb}(\upsilon-1)\beta}{2f_\Delta}\right]^{1/(2-\upsilon)} \frac{ZeB_* f_\Delta R\beta}{\sqrt{2}}. \tag{13}$$

ISM-like medium. In this case, we have

$$E_{\max}({\rm ISM}) \sim Z n_0^{1/2}\epsilon_{B,-1}^{1/2} R_{19} \begin{cases} 10^{16}\ {\rm eV}\ (\frac{\varepsilon_{\rm turb}}{0.5})^3\beta_{-1}^4(12f_\Delta)^{-2} & (\upsilon=5/3),\\ 10^{17}\ {\rm eV}\ (\frac{\varepsilon_{\rm turb}}{0.5})^2\beta_{-1}^3(12f_\Delta)^{-1} & (\upsilon=3/2). \end{cases} \tag{14}$$

The energy conservation $4\pi R^3\beta^2 n m_p c^2/3 \approx E(>\beta)$ yields $R \approx 10^{19}\ {\rm cm}\, A_{52.7}^{1/3}\beta_{-1}^{-(q+2)/3} n_0^{-1/3}$. Combining with equation (14), we have

$$E_{\max}({\rm ISM}) \propto \beta^{\frac{5-\upsilon-q(2-\upsilon)}{3(2-\upsilon)}}, \tag{15}$$

i.e. $E_{\max}({\rm ISM}) \propto \beta^{(7-q)/3}$ for $\upsilon = 3/2$ and $\propto \beta^{(10-q)/3}$ for $\upsilon = 5/3$, both are sensitive to β.

Wind medium. In the termination wind shock model, the stellar wind profile may hold up to a distance $\sim 10^{18}$ cm (e.g. Chevalier et al. 2004). In this case, $n = 3\times 10^{35} A_* R^{-2}\ {\rm cm}^{-3}$. Now $B_* \approx 0.2\, A_{*,-1}^{1/2}\epsilon_{B,-1}^{1/2} R_{17}^{-1}$ Gauss and

$$E_{\max}({\rm wind}) \sim Z A_{*,-1}^{1/2}\epsilon_{B,-1}^{1/2} \begin{cases} 2\times 10^{15}\ {\rm eV}\ (\frac{\varepsilon_{\rm turb}}{0.5})^3\beta_{-1}^4(12f_\Delta)^{-2} & (\upsilon=5/3),\\ 2\times 10^{16}\ {\rm eV}\ (\frac{\varepsilon_{\rm turb}}{0.5})^2\beta_{-1}^3(12f_\Delta)^{-1} & (\upsilon=3/2). \end{cases} \tag{16}$$

As shown in equations (14) and (16), for $\varepsilon_{\rm turb} \sim 0.5$ and $\upsilon = (3/2, 5/3)$, at $\beta \sim \beta_o \sim 0.1$, we have $E_{\max} \sim (10^{17}, 10^{16})Z$ eV. Below we focus on the case of $\upsilon = 3/2$, because in the case of $\upsilon = 5/3$ the request of $E_{\max}(\beta_o) \sim 3\times 10^{17}$ eV is more difficult to satisfy. For $\beta \sim 0.5$, the stochastic gyroresonant acceleration is able to accelerate protons to $\sim 10^{19}$ eV (see also Dermer 2001a). The accelerated particle spectrum is thus ($\upsilon = 3/2$)

$$\frac{dN_{\rm CR}}{dE_{\rm CR}} \propto \begin{cases} E_{\rm CR}^{-(2+\delta)-\frac{3q}{7-q}} & \text{for } k=0,\\ E_{\rm CR}^{-(2+\delta)-\frac{q}{3}} & \text{for } k=2. \end{cases} \tag{17}$$

As shown in Section 2.1, the main deceleration of the HN outflow is very likely to be in an homogenous medium. The accelerated protons have a spectrum $dN/dE_{\rm CR} \propto E_{\rm CR}^{-(2.4+3q/(7-q))}$. To match the observation $\Delta\gamma({\rm II}) \approx -3q/(7-q) \sim -0.9$, we need

$q \sim 1.6.$

This is surprisingly close to the value ~ 1.7 that is inferred from the optical modelling of SN 2003lw. The detailed optical modelling of more GRB-associated HN explosions is highly needed to better constrain q and then confirm or rule out our interpretation.

If GRB-associated HNe expand into a wind bubble-like medium, a flatter CR spectrum in the higher energy range would appear. At a small radius (say, $<10^{18}$ cm), the medium is free wind-like and the accelerated particle spectrum is $\propto E_{\rm CR}^{-(2+\delta)-q/3}$, which then gets steepened by a factor of $q(2+q)/[3(7-q)] \sim 0.4$ for $q \sim 1.6$ after entering the ISM-like medium. Such a flattening seems not enough to match the observation $\Delta\gamma({\rm III}) \sim 0.6$. So CRs above the ankle may be mainly from AGNs, as indicated by the recent analysis of the correlation of the highest energy CRs with nearby extragalactic objects by the Pierre Auger Collaboration (Abraham et al. 2007).

The rate of local GRB-associated HNe only accounts for $\sim$(0.1–0.5) per cent of that for all local SNe (Della Valle 2006; Soderberg 2007). The typical energy of these HNe, however, is 50 times larger than that of the normal SNe. Roughly, we expect that a fraction $\sim$10 per cent of CR protons at 3 PeV could be attributed to GRB-associated HNe. It is enough to match the observation (Ulrich et al. 2004; Hörandel 2008). So the CR proton spectrum in the energy range of 10^{16}–$10^{18.5}$ eV may be quantitatively interpreted.

4 DISCUSSION AND SUMMARY

The particle acceleration in marginally relativistic HN shocks is discussed. The GRB-associated HN outflows are assumed to have a fairly steep energy distribution against their velocities, i.e. $E(\geq\beta) \propto \beta^{-q}$ for $q \sim 1.7$, as inferred from the optical modelling of SN 2003lw (see Section 2.1 for details). A significant fraction of a HN's kinetic energy is carried by the material moving with a velocity $>\beta_o(\sim 0.1)$, driving an energetic shock wave into the surrounding medium. The cosmic ray protons above the second knee but below the ankle may be accelerated by the HN shocks in the velocity range of $\beta \sim (1$–$4)\beta_o$. To satisfy this velocity bound, the HN outflows associated with GRBs must have reached a very large radius where the surrounding medium is very likely to be

ISM-like (see Section 2.2 for details). With this argument, the self-magnetized shock acceleration model adopted in Wang et al. (2007) would yield a much steeper spectrum that is inconsistent with the observation unless the medium surrounding the HN is a free wind holding up to a radius $R_{\rm dec} \sim 10\,{\rm kpc}\ (M_{\rm ej}/10\,{\rm M_{\odot}})A_{*,-1}^{-1}$. Such a request seems difficult to satisfy. A highly speculative solution is outlined in Section 3.2.1.

In this work, we find that for $q \sim 1.6$, the stochastic gyroresonant acceleration model can account for the spectrum change of high-energy protons around the second knee (see Section 3.2.2 for details). As a consequence, the stochastic gyroresonant acceleration mechanism in a relativistic GRB forward shock may account for part of the ultrahigh-energy CRs ($\sim 10^{20}$ eV), as suggested in Dermer (2001b, 2007) and Dermer & Humi (2001). A typical

$$q \approx \frac{-7\Delta\gamma({\rm II})}{3-\Delta\gamma({\rm II})} \sim 1.6 \text{ for } \Delta\gamma({\rm II}) \sim -0.9,$$

if confirmed in future optical modelling of the GRB-associated HN explosions, will be crucial evidence for our current speculation.

ACKNOWLEDGMENTS

We thank an anonymous referee and Charles D. Dermer for constructive comments, and Jin-Song Deng, Zhuo Li and Dong Xu for communication/discussion. This work is supported by a (post-doctoral) grant from the Danish National Science Foundation, the National Natural Science Foundation (grant 10673034) of China and a special grant of the Chinese Academy of Sciences.

REFERENCES

Abraham J. et al., 2007, Sci, 318, 938
Antoni T. et al., 2005, Astropart. Phys., 24, 1
Bell A. R., Lucek S. G., 2001, MNRAS, 321, 433
Berezhko E. G., Völk H. J., 2004, A&A, 427, 525
Budnik R., Katz B., MacFadyen A., Waxman E., 2008, ApJ, 673, 928
Chevalier R. A., 2007, in Wijers R. A. M. J., Kaper L., van Eerten H. J., eds, The Next Decade of Gamma-Ray Burst Afterglows, preprint (arXiv:0706.0401)
Chevalier R. A., Li Z. Y., Fransson C., 2004, ApJ, 606, 309
Dai Z. G., Wu X. F., 2003, ApJ, 591, L21
Della Valle M., 2006, in Holt S. S., Gehrels N., Nousek J. A. eds, Proc. AIP Conf. Ser. Vol. 836, Gamma-Ray Bursts in the Swift Era, Sixteenth Maryland Astrophysics Conference. Am. Inst. Phys., New York, p. 367
Dermer C. D., 2001a, in Schlickeiser R., ed., Proc. 27th International Cosmic Ray Conference. Hamburg, Germany, p. 72
Dermer C. D., 2001b, in Schlickeiser R., ed., Proc. 27th International Cosmic Ray Conference. Hamburg, Germany, p. 2039 (arXiv:astro-ph/0012490)
Dermer C. D., 2007, preprint (arXiv:0711.2804)
Dermer C. D., Humi M., 2001, ApJ, 556, 479
Erlykin A. D., Wibig T., Wolfendale A. W., 2001, New. J. Phys., 3, 18
Fan Y. Z., Piran T., Xu D., 2006, J. Cosmol. Astropart. Phys., 0609, 013
Frail D. A., Soderberg A. M., Kulkarni S. R., Berger E., Yost S. A., Fox D. W., Harrison F. A., 2005, ApJ, 619, 994
Fryer C. L., Rockefeller G., Young P. A., 2006, ApJ, 647, 1269
Gruendl R. A., Chu Y. H., Dunne B. C., Points S. D., 2000, ApJ, 120, 2670
Hillas A. M., 2005, J. Phys. G: Nucl. Part. Phys., 31, R95
Hjorth J. et al., 2003, Nat, 423, 847
Hörandel J. R., 2008, Adv. Space Res., 41, 442
Li Z. Y., Chevalier R. A., 1999, ApJ, 526, 716
Mazzali P. A. et al., 2006, ApJ, 645, 1323
Pihlstróm Y. M., Taylor G. B., Granot J., Doeleman S., 2007, ApJ, 664, 411
Ptuskin V. S., Zirakashvili V. N., 2005, A&A, 429, 755
Ramirez-Ruiz E., Dray L. M., Madau P., Tout C. A., 2001, MNRAS, 327, 829
Ramirez-Ruiz E., Granot J., Kouveliotou C., Woosley S. E., Patel S. K., Mazzali P. A., 2005, ApJ, 625, L91
Soderberg A., 2007, in Immler S., Weiler K. W., McCray R., eds, AIP. Conf. Ser. Vol. 937, Supernova 1987A: 20 years after: Supernovae and Gamma-Ray Bursters. Am. Inst. Phys., New York, p. 492 (arXiv:0706.3047)
Soderberg A. et al., 2004, Nat, 430, 648
Soderberg A. et al., 2006, Nat, 442, 1014
Sveshnikova L. G., 2003, A&A, 409, 799
Tan J. C., Matzner C. D., McKee C. F., 2001, ApJ, 551, 946
Ulrich H. et al., 2004, Eur. Phys. J. C, 33, 944
van der Horst A. J. et al., 2008, A&A, 480, 35
van Marle A. J., Langer N., Achterberg A., Garcia-Segura G., 2006, A&A, 460, 105
Wang X. Y., Razzaque S., Mészáros P., Dai Z. G., 2007, Phys. Rev. D, 76, 083009
Waxman E., 2004, ApJ, 605, L97
Wick S. D., Dermer C. D., Atoyan A., 2004, Astropart. Phys., 21, 125
Wijers R. A. M. J., 2001, in Costa E., Frontera F., Hjorth J., eds, Gamma-Ray Bursts in the Afterglow Era. Springer-Verlag, Berlin, p. 306
Xu D., Zou Y. C., Fan Y. Z., 2008, preprint (arXiv:0801.4325)
Xue R. R., Tam P. H., Wagner S. J., Behera B., Fan Y. Z., Wei D. M., 2008, ApJ, submitted

This paper has been typeset from a TEX/LATEX file prepared by the author.