STOCHASTIC ELECTRON ACCELERATION IN SHELL-TYPE SUPERNOVA REMNANTS

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ABSTRACT

We study the stochastic electron acceleration by fast-mode waves in the turbulent downstream of weakly magnetized collisionless astrophysical shocks. The acceleration is most efficient in a dissipative layer, and the model characteristics are determined by the shock speed, density, magnetic field, and turbulence decay length. The model explains observations of shell-type supernova remnants RX J1713.7−3946 and J0852.0−4622 and can be tested by observations in hard X-rays with the HXMT and NuSTAR or γ-rays with GLAST.

Subject headings: acceleration of particles — MHD — plasmas — shock waves — turbulence

Online material: color figures

1. INTRODUCTION

Recent observations of high-energy emission from a few shell-type supernova remnants (STSNRs) advance our understanding of the underlying physical processes significantly (Enomoto et al. 2002). These observations not only give high-quality emission morphology and spectra (Uchiyama et al. 2003; Aharonian et al. 2004; Cassam-Chenaï et al. 2004; Hiraga et al. 2005; Takahashi et al. 2008a), which lead to tight upper limits on thermal X-rays, but also discover X-ray variability on a timescale of a few months (Uchiyama et al. 2007), which is intimately connected to processes near the shock front (SF). The radio to X-ray emissions are produced by relativistic electrons through the synchrotron processes. The nature of the TeV emission is still a matter of debate. The major challenges to the hadronic scenario include (1) the lack of correlation between the TeV emission and molecular cloud distribution; (2) the requirement of very high mean target gas density, supernova explosion energy, and/or proton acceleration efficiency, which implies efficient magnetic field amplification in the context of the diffusive shock acceleration and strong suppression of the electron acceleration (Lucek & Bell 2000; Berezhko & Völk 2006; Butt et al. 2008); and (3) the good correlation between the X-ray and TeV emission (Plaga 2008; Aharonian et al. 2006, 2007a), which suggests a leptonic scenario. Because secondary leptons from hadronic processes will not produce as much radiation power as γ-rays from neutral pion (π0) decays (Aharonian et al. 2007b), the fact that the X-ray luminosity is slightly higher than the TeV luminosity requires that most leptons be accelerated from the background plasma.

Compared to the hadronic scenario, the leptonic scenario has much fewer parameters and is well-constrained by the observed radio, X-ray, and TeV spectra. In the simplest one-zone model, the TeV emission is produced through the inverse Compton (IC) scattering of the Galactic interstellar radiation by the same relativistic electrons producing the radio and X-rays (Porter et al. 2006). The TeV spectrum can be used to infer the high-energy electron distribution. One then needs to fit both the synchrotron spectral shape and flux level by adjusting the magnetic field alone. TeV observations always see spectral softening with the increase of the photon energy, suggesting a spectral cutoff (Aharonian et al. 2007a, 2007b). There are also indications that the synchrotron spectrum cuts off in the hard X-ray band (Takahashi et al. 2008a, 2008b; Uchiyama et al. 2007). A single-power-law model with an exponential cutoff for the electron distribution has difficulties in reproducing the broad TeV spectrum. Models with more gradual cutoffs lead to much better fits. Relativistic particle distributions with gradual high-energy cutoffs are a natural consequence of the stochastic acceleration (SA) by plasma waves in magnetized turbulence (Park & Petsio 1995; Becker et al. 2006; Stawarz & Petsio 2008). The particle distribution is determined by the interplay among the acceleration, cooling, and escape processes with the cutoff shape determined by the energy dependence of the ratio of the acceleration to cooling or escape rate. Given the low magnetic fields derived from leptonic models, the radiative cooling time is much longer than the remnant lifetime. The gradual high-energy cutoff has to be caused by the balance between the acceleration and escape processes.

We consider the evolution of weakly magnetized turbulence in the downstream of strong nonrelativistic shocks and study the SA by fast-mode waves. The large-scale turbulence cascades following the Kolmogorov phenomenology (Boldyrev 2002; Jiang et al. 2008) until it reaches the scale, where the eddy turnover time becomes comparable to the period of fast-mode waves. The Iroshnikov-Kraichnan phenomenology prevails on even smaller scales. The collisionless damping of plasma waves sets in at the coherent length of the magnetic field, where the eddy turnover time is comparable to the period of Alfvén waves. Fast-mode waves propagating nearly parallel to the mean magnetic field can survive the transit-time damping (TTD) by the thermal background particles, and accelerate some electrons through cyclotron resonances to a power-law high-energy distribution, which cuts off at the energy, where the particle gyroradius reaches the coherent length of the magnetic field (Liu et al. 2006). Acceleration of higher energy electrons by the nearly isotropic fast-mode waves on larger scales leads to a gradual cutoff. SA of electrons can naturally account for emissions from the STSNRs with a weak magnetic field. Future observations in the hard X-ray and MeV bands may test the model. The particle acceleration model is described in § 2 and applied to SNRs RX J1713−3946 and J0852.0−4622 to derive quantitative results in § 3. Section 4 gives the conclusions.
2. TURBULENCE CASCADE, WAVE DAMPING, AND STOCHASTIC PARTICLE ACCELERATION

The collisionless nature of weakly magnetized shocks implies difficulties in converting the free energy in the system into heat and energetic particles instantaneously at a narrow SF. Instead, the high magnetic Reynolds number implies strong turbulence. Previous studies of diffusive shocks focus on its roles in making high-energy particles crossing the SF repeatedly. The turbulence can also carry a significant fraction of the released free energy and decay gradually. The turbulence generation mechanism is not well understood. It is usually attributed to the Weibel instability (Weibel 1959) of particle streamlines and may also depend on the properties of the upstream plasma. We here assume that the turbulence is isotropic and generated at a scale $L$ with an eddy speed $u$. The free energy dissipation rate is then given by $Q = C_{\parallel} \rho u^2 L$, where $C_{\parallel} \sim 1$ is dimensionless. For strong nonrelativistic shocks with the shock frame upstream speed $U$ much higher than the speed of the upstream fast-mode waves $v_p = (v_p^2 + 5v_F^2/3)^{1/2}$; mass, momentum, and energy conservation across the SF require $U^2 = 5v_F^2 + 5u^2 + 2v_p^2 + U^2/16$, where the Alfvén speed $v_A = (B^2/4\pi \rho)^{1/2}$, the isothermal sound speed $v_s = (\rho \rho)^{1/2}$, and $B, \rho, P$ are the magnetic field, mass density, and pressure, respectively. The shock structure can be complicated due to the presence of strong turbulence, and the speed $v_s, v_A$, and $u$ should be considered as averaged quantities over $\sim L$, which also corresponds to the energy dissipation scale. The excitation of plasma waves should be very efficient in the supersonic phase with $u \gtrsim v_A$, and the waves may prevail in the downstream, accelerating some particles in the background plasma to very high energies.

The eddy turnover speed and time are given respectively by $v_{edd}(k) = 4\pi W(k) k^3$ and $\tau_{edd}(k) = 2\pi/C_{\parallel} k v_{edd}$, where $W(k) = (4\pi)^{-1} (2\pi Q \rho)^{1/2} k^{-11/3}$, $k = 2\pi/l$, and $l$ are the isotropic turbulence power spectrum, wavenumber, and eddy size, respectively. At the turbulence generation scale $k_m = 2\pi l/w$, $v_{edd} = u$, and the total turbulence energy is given by $\int W(k) k^2 d k = (3/2) w^2$. The MHD wave period is given by $\tau_p(k) = 2\pi/v_p k$. Then the transition scale from the Kolmogorov to Kraichnan phenomenology $k_c$ occurs at $\tau_p(k_c) = \tau_{edd}(k_c)$, which gives $k_c = (C_{\parallel} w v_p^2)^{1/3} k_m$. For $k > k_c$, the turbulence spectrum in the inertial range is given by

$$W(k) = (4\pi)^{-1} (v_p/C_{\parallel})^{1/2} u^{3/2} k_{m}^{1/2} k^{-7/2}. \tag{1}$$

The collisionless damping starts at the coherent length of the magnetic field $l_d = 2\pi/k_B$, where the period of Alfvén waves $2\pi/k_B v_A$ is comparable to the eddy turnover time, $\tau_{edd} = (C_{\parallel} u v_p)^{1/3} k_m$. Both $v_A$ and $v_p$ may be lower than $u$ near the SF. If $u \gg v_A, v_p$ at the SF, fast, slow, and Alfvén waves can all be excited, and the consequent plasma heating quickly leads to $v_A \gtrsim v_p$. Fast-mode waves with the highest phase speed should dominate the SF processes. We assume that they dominate the energy dissipation as the turbulent flow moves away from the SF and $v_p$ evolves from $v_A$ to $u$ at the same time. For a fully ionized hydrogen plasma with isotropic particle distributions, which are reasonable in the absence of strong large-scale magnetic fields, the TTD rate is given by (Stix 1962; Quaataer 1998; Petrov et al. 2006)

$$\Lambda_p(\theta, k) = \frac{(2\pi k_B)^{1/2} \sin^2 \theta}{2(m_p + m_e) \cos \theta} \left[ (T_m, m_p)^{1/2} \exp \left( -\frac{m_p \omega^2}{2k_B T_m k_{||}^2} \right) + (T_r, m_p)^{1/2} \exp \left( -\frac{m_p \omega^2}{2k_B T_r k_{||}^2} \right) \right]. \tag{2}$$

where $T_p, T_r, m_p, m_e, \theta, \omega, and k_{||} = k \cos \theta$ are the electron and proton temperatures, masses, angle between the wave propagation direction and mean magnetic field, wave frequency, and parallel component of the wavevector, respectively. The first and second terms in the brackets on the right-hand side correspond to damping by electrons and protons, respectively. For weakly magnetized plasma with $v_e < v_p$, proton heating always dominates the TTD for $\omega^2/k_{||}^2 \sim v_p^2 < 2k_B T_p/m_p$ (Quaataer 1998). If $v_p$ does not change dramatically in the downstream, the continuous heating of background particles through the TTD processes makes $T_p \rightarrow (m_p/m_e) T_r$, since the heating rates are proportional to $(mT)^{5/2}$, where $m$ and $T$ represent the mass and temperature of the particles, respectively. Parallel propagating waves with sin $\theta = 0$ are not subject to the TTD. Obliquely propagating waves are damped efficiently by the background particles. Although the damping rates for waves propagating nearly perpendicular to the magnetic field with $\cos \theta = 0$ are also low, these waves are subject to damping by magnetic field wandering (Petrov et al. 2006). The turbulence power spectrum cuts off sharply when the damping rate becomes comparable to the turbulence cascade rate $\Gamma = \tau^{-2} (\tau_{edd}^{-1} + \tau_p^{-1}) \approx \tau_{edd}^{-1} \tau_p^{-1}$ (Jiang et al. 2008). One can define a critical propagation angle $\theta_c(k)$, where $\Lambda_p(\theta_c, k) = \Gamma(k)$. Then for $k a_T = C_{\parallel} m_p / u^2$ with $C_{\parallel} < 3/5$ for $v_p < u$, equations (1) and (2) give

$$\frac{v_p^2}{C_{\parallel}^2 u^2 v_p k^2} \approx \sin^2 \theta_c \exp \left( -\frac{v_p^2}{2C_{\parallel} u^2 \cos^2 2\theta_c} \right),$$

where the electron pressure and damping have been ignored. For $v_e \ll U$, we have $3v_p^2 / 5 = v_p^2 \approx C_{\parallel} u^2 = [3C_{\parallel}^2 / (2C_{\parallel} + 1)] U^2$. The scattering and acceleration times of relativistic electrons by these nearly parallel propagating waves are energy independent and given, respectively, by $\tau_{es} = 2\pi k_B s_{es}$ and $\tau_{es} = (3c^2 v_p^2 s_{es})$ (Liu et al. 2006; Blandford & Eichler 1987), where $c$ is the speed of light. The electron escape time from the large-scale eddies is given by $\tau_{esc} = L/c = 4c^3 s_{esc}$, and the spectral index of the accelerated electrons in the steady state is given by $p = (9/4 + \tau_{esc} / \tau_{es})^{3/2} - 1/2 = (9/4 + 12c^2 v_p^2 / C_{\parallel} u^2 v_p^2) - 1/2$. The maximum energy that electrons can reach is given by $\gamma m_e c^2 = 2\pi q B k_B = 7.26 (p + 0.5)^{5/2} [2C_{\parallel} / (C_{\parallel} + 1)]^{1/2} (L/10^{18})^2 (c/B) [c/10^4] (s/10^{-10})$ cm, where $U_0 = U/0.015c$ and $q$ is the elementary charge units. The ratio of the dissipated energy carried by nonthermal particles to that of the thermal particles should be greater than $\eta = \theta_c^2 / \theta_c^2 = 2\pi v_p^2 / 2\pi k_B (\pi^2 / 4 k_B) = e / 2.72$, since the isotropic turbulence with $k < k_c$ can also accelerate particles with the Lorentz factor $\gamma \geq \gamma_c$.

Electrons with even higher energy interact with the nearly isotropic fast mode waves at $k < k_c$. The corresponding accel-
...The thick solid line corresponds to the fiducial model with $\beta = 0.5$, $p = 1.85$, $\gamma = 7.2 \times 10^6$, and $B = 12.0 \mu G$. The thin solid line with $p = 2.0$, $B = 12.0 \mu G$, and $\gamma = 4.4 \times 10^7$ is the best for an electron distribution with an exponential cutoff, i.e., $\beta = 1.0$. The fiducial model gives a better fit to the X-ray and TeV spectra. The inset is an enlargement of the TeV spectrum to validate the usage of the steady-state solution of the particle kinetic equation and implies $1/2 \frac{d}{dt} \frac{E_{\gamma}}{E_{\gamma}} = 0.5$. The solid lines are for $p = 1.85$ and $\beta = 0.25$ and 1.0. The thin solid lines are for $\beta = 0.5$ and $p = 1.7$ and 2.0. [See the electronic edition of the Journal for a color version of this figure.]

3. RESULTS

Several STSNRs have been observed extensively in the radio, X-ray, and TeV bands. X-ray observations with Chandra, XMM-Newton, and Suzaku, and TeV observations with HESS have made several surprising discoveries that challenge the classical diffusive shock particle acceleration model. The SNR RX J1713.7−3946 is about $t = 1600$ years old (Wang et al. 1997) with a radius of $R \approx 10$ pc and a distance of $D \approx 1$ kpc. By fitting its broadband spectrum with the above electron distribution and background photon field given by Porter et al. (2006) through the synchrotron and IC processes, we find that $p = 1.85$, $B = 12.0 \mu G$, $\gamma, m_c^2 c^2 = 3.68$ TeV, and the total energy of electrons with $\gamma > 1800$ (Liu et al. 2006; Fan et al. 2008), $E_e = 3.92 \times 10^{44}$ erg (Fig. 1). The thin line shows the best fit with $\beta = 1$, whose TeV spectrum is not broad enough to explain the HESS observations. In the hard X-ray band, our model predicted emission spectrum is significantly harder and brighter, which gives better fit to the Suzaku observations and can be tested with future HXMT and NuSTAR observations. Our model also predicts a higher MeV flux testable with future GLAST observations. Figure 2 shows the dependence of the spectrum on $p$ and $\beta$.

The distribution of relativistic electrons obtained above implies $L = 2.34 \times 10^{47} (C_{sp}/C_{sp} + 1)^{-1/2} \gamma U_t^{-1} \text{ cm}$, and the electron density $n_e = 2.00 \times 10^{-3} (C_{sp} + 1)^{5/2} C_{sp}^{1/2} C_{e}^{1/2} U_t^{-1/2} \text{ cm}^{-3}$, where $Y = 0.1$ is the helium abundance. For the steady-state solution to be applicable, $C_{sp} < 0.824 \frac{C_{sp}^{1/2}}{C_{sp}^{1/2}} < 0.64$. For $C_{sp} = C_{sp} = 0.5$, we have $L = 4.04 \times 10^{47} U_t^{-1} \text{ cm}$, $n_e = 1.33 \times 10^{-3} \gamma U_t^{-1/2} \text{ cm}^{-3}$, $t_{\text{ed}} = 138 \gamma t_{\text{yr}}$, $t_{\text{ed}} = 42.3 U_t^{-1/2} \text{ yr}$, and $\eta = 0.384 U_t^{-1/2}$. The energy carried by the magnetic field $E_B = B^2 R^2 / 2 = 2.77 \times 10^{44} U_t^{-1} \text{ erg}$. The short coherent length of the magnetic field $2\pi/k_{\text{p}} = \gamma, m_e^2 c^2 / q B = 1.02 \times 10^{15} \text{ cm}$ is consistent with the low level of polarization from most part of the radio image (Lazendic et al. 2004). If $k_{\text{p}}$ does not change significantly in the downstream, the distance that relativistic electrons diffuse through the $t = 1600 \text{ yr}$ lifetime of the remnant is about $(t/t_{\text{ed}})^{5/2} L = 0.805 \text{ pc}$. Most of them are therefore trapped near the SF. The total kinetic energy carried by the supersonic dissipative layer in the lab frame is $K = 2\pi R^2 L (1 + 2Y) n_e (3U_t/4)^{1/2} = 7.34 \times 10^{46} U_t^{-3/2} \text{ erg}$. The corresponding gas mass $M = 2k/(3U_t/4)^{1/2} = 1.29 \times 10^{32} U_t^{-3/2} g$, thermal energy $E = 6\pi R^2 L m_e (1 + 2Y) v = 3.67 \times 10^{42} U_t^{-3/2} \text{ erg}$, gas temperature $k_{\text{p}} T_g = 2m_e (1 + 2Y) E / 3M (1 - Y) = 26.4 U_t^0 \text{ keV}$, where we have assumed that the temperatures of protons and helium ions are the same and $T_{\text{p}} = T_g$.

The acceleration efficiency of $\gamma \eta$ is not sensitive to $p$ and $U$ and therefore does not change significantly through the evolution of the SNR, implying that a thermal energy of $E / \eta \sim 1.02 \times 10^{43} U_t^{-3/2} \text{ erg}$, which is 27.8 times higher than $E_e$, be produced in accompany with the electron acceleration. The mass of the shocked thermal plasma is about $E, M = \eta E_e = 5.7 \times 10^{40} U_t^{-3/2} M_{\odot}$.
3.58 × 10^{33} U_0^{-5/2} g = 1.79 U_0^{-5/2} M_6. These suggest that most of the relativistic electrons were accelerated more than τ_{shock} ~ 100 yr before outside the current dissipative layer. If the shocked plasma is uniformly distributed within a shell of 4.5 pc (Aharonian et al. 2006), the electron density is 1.74 × 10^{-3} U_0^{-5/2} cm^{-3}, which is comparable to that inferred from X-ray observations (Cassam-Chenai et al. 2004). However, X-ray observations of the northwest outer edge of the remnant suggest \( U < 0.015c \) (Uchiyama et al. 2007). So the model-predicted thermal X-ray emission likely exceeds the observed value. The northwest edge is interacting with dense molecular clouds. It is possible that the shock speed in a low-density medium is still greater than the observed upper limit of 0.063 \( U_0^{-5/2} \) erg, \( M_6 = 2.01 \times 10^{-3} U_0^{-5/2} \) g, \( E = 5.71 \times 10^{-4} U_0^{-5/2} \) erg, \( k_B T_p = 26.4 \times U_0^{-5/2} \) keV, and \( E_{\gamma} = B^2 R^2 L/2 = 4.53 \times 10^{-6} U_0^{-5/2} \) erg. These two remnants are very similar.

4. CONCLUSIONS

The steady-state distribution of relativistic electrons accelerated by fast-mode waves in a downstream turbulent dissipative layer with a thickness of weakly magnetized collisionless astrophysical shocks can be approximated as

\[
\frac{f}{c} \propto \gamma^{-p} \exp \left(-\frac{\gamma/\gamma_c}{2}\right),
\]

where \( p = (9/4 + 2^2/3 c^2 \gamma_b^2 S U_{}^{10/3})^{1/2} - 1/2, \gamma_c \), \( m_c c^2 \), the ratio of the dissipated energies going into nonthermal and thermal particles is \( \sim \eta = (e^{2.96/5 \pi} / 5) \gamma_b^{3/2} U_{}^{5/2} U_c^{1/2} \). For \( p > 0.92 \), the turbulence decay time is shorter than the particle acceleration time and one needs to consider the time-dependent solution. The accelerated electrons in the STSNRs give excellent fits to the broadband spectra of SNRs RX J1713.7−3946 and J0852.0−4622. The model attributes the recently observed variable X-ray features to inhomogeneity in the upstream and the spatial diffusion of relativistic electrons near the cutoff energy and predicts a harder X-ray spectrum and higher hard X-ray and \( \gamma \)-ray fluxes than leptonic models with a sharper cutoff in the electron distribution. Future high-sensitivity hard X-ray and \( \gamma \)-ray observations will test the model. The major uncertainty of the model is related to the turbulence generation mechanism near the collisionless SF. In this Letter, we assume that the turbulence is isotropic and has a characteristic eddy size of \( L \). Little is known about the evolution of anisotropic turbulence. The latter is valid as far as the free energy dissipation through shocks proceeds over the scale \( L \) instead of instantaneously at a narrow region (\( \ll L \)) near the SF as suggested in the diffusive shock models.

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