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maximum particle Energies by Fermi acceleration and the origin of cosmic rays above the knee

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Abstract. We derive the maximum accelerated particle energy from first-order and second-order Fermi acceleration at nonrelativistic and relativistic shocks for explosions taking place in a uniform surrounding medium. Second-order stochastic processes in relativistic flows are shown to be capable of accelerating cosmic rays to ultra-high energies. Cosmic rays above the knee of the cosmic ray spectrum can be accelerated by the second-order Fermi mechanism in relativistic flows, such as those occuring in gamma-ray bursts and unusual supernovae like SN 1998bw.

1 Introduction

Although there is a consensus that cosmic rays with energies less than the knee energy at $\approx 3 \times 10^{15}$ eV are accelerated by supernova remnant shocks, the origin of higher energy cosmic rays is unknown. Proposed sources of cosmic rays above the knee include galactic sources such as microquasars and pulsars, active galactic nuclei (Protheroe and Szabo, 1992), merger shocks in galaxy clusters (Kang Ryu, and Jones, 1996), dormant AGNs (Levinson and Boldt, 2001), and gamma-ray bursts (GRBs) (Milgrom and Usov, 1996; Dermer and Humi, 2001). Proposed sources of ultra-high energy ($\gtrsim 10^{19}$ eV) cosmic rays (UHECRs) include the termination shocks at the lobes of FRII jets (Rachen and Biermann, 1993), GRBs (Waxman, 1995; Vietri, 1995; Dermer, 2000), and exotic particles, such as monopoles and Z-bursts (for reviews, see Nagano and Watson (2000); Weiler (2001)).

Here we consider particle acceleration at nonrelativistic and relativistic shocks. We show that the sources of cosmic rays near and above the knee of the cosmic-ray spectrum could result from particle acceleration at the relativistic shocks produced by the baryon-dilute outflows of fireball transients and GRBs (Dermer, 2000; Dermer and Humi, 2001).The maximum energy that can be reached by a particle accelerated at a shock due to constraints upon diffusive transport and available time is derived, extending earlier treatments for nonrelativistic shocks by, e.g., Lagage and Cesarsky (1983), Drury (1983), and Baring et al. (1999).

In §2, the maximum particle energy is derived for the extremes of quasi-parallel (q-$\|$) and quasi-perpendicular (q-$\perp$) shocks (Jokipii, 1987; Baring et al., 1999). We find that the maximum particle energy is not greatly increased in the latter case due to a geometric constraint that applies in cases where a large scale toroidal magnetic field is absent. In §3, we consider the highest energy that particles can be accelerated by shock Fermi processes at a relativistic external shock, following Gallant and Achterberg (1999). In §4 we show that the highest energies that can be reached by Fermi processes in supernova or GRB-type explosions result from second-order gyroresonant stochastic acceleration in the shocked fluid of a relativistic blast wave (Waxman, 1995; Rachen and Mésźaros, 1998; Schlickeiser and Dermer, 2000; Dermer and Humi, 2001). This maximum particle energy E can exceed $\approx 10^{20}$ eV in GRB blast waves for reasonable parameter values. Cosmic rays with energies above the knee of the cosmic ray spectrum is proposed to originate from second-order Fermi acceleration at relativistic shocks (§5).

2 Nonrelativistic Shocks

We treat first-order Fermi acceleration of relativistic nonthermal particles with $v' \approx c$. Our notation is that $\beta(x)c = \beta c$ is the flow speed, and the blast-wave evolution is described by momentum $P = \beta\Gamma$. Primes refer to quantities in the comoving shock frame. In the primed frame stationary with respect to the shock, the upstream ($-$) flow approaches with speed $u_- = \beta_- c = 4\beta c/3$ and the downstream (+) flow recedes with speed $u_+ = \beta_+ c = \beta c/3$. The quantity $u = u_- - u_+ = \beta c$ is the speed of the shocked fluid and u_- is the speed of the shock as measured in the stationary upstream frame.

Let $p' = \sqrt{\gamma'^2 - 1}$ represent the dimensionless momen-

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tum of a particle with Lorentz factor γ' in the comoving shock frame. Then $\dot{p}'_{\rm FI} \cong \Delta p'/t_{cyc}$, where the cycle time t_{cyc} is given by the diffusive properties of the upstream and downstream regions. The average change in particle momentum over a complete cycle for relativistic test particles with $p' \gg 1$ is $\Delta p' = 4\beta p'/3$, when $\beta \ll 1$ (Schlickeiser, 1984; Gaisser, 1990), provided that there are scatterers in the upstream and downstream flow to isotropize the particles. In a one-dimensional flow,

$$t_{cyc} = \frac{4}{v'}\,(\frac{\kappa_-}{u_-} + \frac{\kappa_+}{u_+}) \;= \frac{4}{v'u_-}(\kappa_- + \chi\kappa_+), \quad (1)$$

where $\chi = u_-/u_+$ is the compression ratio and the diffusion coefficient $\kappa_\pm = \lambda_\pm v'/3 = \eta_\pm r_{\rm L} v'/3 = \eta_\pm r^o_{\rm L\pm} p'v'/3$, where $r_{\rm L\pm} = r^o_{\rm L\pm} p' = mc^2 p'/qB_\pm$ is the Larmor radius for a particle of mass m and charge q. Here we write diffusion coefficients in terms of the parameters $\eta_\pm$ that give the particle mean-free-paths scaled to the values implied by the Bohm diffusion limit evaluated for the local magnetic field B.

In q-$\|$ shocks, $B_-/B_+ \cong 1$. Hence

$$\dot{p}'_\| \cong \frac{u_-\beta}{r^o_{\rm L-}(\eta_- + \chi\frac{B_-}{B_+}\eta_+)} \;\lesssim \frac{u^2\chi}{cr^o_{\rm L-}(\chi^2-1)} \;\equiv \dot{p}'_{\|,max}. \quad (2)$$

The last relation defines the Bohm-diffusion-limit maximum acceleration rate for q-$\|$ shocks, assuming that the Bohm diffusion limit sets a lower limit to the diffusivity of a medium.

In q-$\perp$ shocks, a scattering event shifts a particle in the $\hat{x}$ direction by a mean distance $r_{\rm L}$; therefore $\lambda \to r_{\rm L}$. The effective particle drift speed is reduced by a factor η^{-1} because particles are confined within the gyro-radius size scale $r_{\rm L}$. Hence $\kappa_\perp \cong r_{\rm L} v'/3\eta$ or $\kappa_\perp \cong \kappa_\|/(1+\eta^2)$ (Jokipii, 1987), and

$$\kappa_{\perp\pm} = \frac{\eta_\pm}{1+\eta^2_\pm}\;\frac{r_{\rm L\pm}v'}{3}\;;\; r^o_{\rm L+} = \frac{B_-}{B_+} r^o_{\rm L-} = \frac{1}{\chi} r^o_{\rm L-}\;, \quad (3)$$

implying

$$\dot{p}'_\perp \lesssim \frac{u^2}{cr^o_{\rm L-}}\;\frac{\chi\eta}{\chi^2-1}\;\equiv \dot{p}'_{\perp,max}\;, \quad (4)$$

where $\eta = \min(\eta_-,\eta_+)$.

According to Jokipii (1987), to maintain an isotropic particle distribution requires $\eta \lesssim v'/u \sim 1/\beta$ for relativistic particles. If shock drift takes particles to regions of smaller obliquenesses, then the acceleration rate is reduced (Jokipii, 1987; Baring et al., 1999). Particles will drift along a q-$\perp$ shock to scatter in a region of smaller obliqueness when $\eta r_{\rm L} - /x \lesssim \theta_{1/2}$, where $\theta_{1/2}$ is a characteristic angle over which the shock obliquity changes. If the surrounding magnetic field is homogeneous, $\eta_- \lesssim \min[\beta^{-1}, \sqrt{x/r_{\rm L-}}]$. The condition $\eta \lesssim \sqrt{x/r_{\rm L-}}$ implies that $\eta \lesssim \beta^{-1/3}$.

The energy gain rates are more conveniently written in terms of the differential distance $dx = \beta_s\Gamma_s c dt' \to \beta_s c dt' = u_- dt' \cong u_- dt$ that the shock advances during the differential time element dt, as measured in the stationary frame of the external medium. From the above results, we find that the maximum energy gains per unit distance for quasi-parallel and quasi-perpendicular nonrelativistic shocks are, respectively,

$$\frac{dE'}{dx}|_{\|,max} \simeq \beta q B_-\;\;;\;\; \frac{dE'}{dx}|_{\perp,max} \simeq \beta^{2/3} q B_-\;\;. \quad (5)$$

The maximum particle energy in nonrelativistic supernova shocks is derived by considering a uniform medium with density n_0 that carries a homogeneous external magnetic field with intensity $B_- = B_0 = 10^{-6} B_{\mu G}$ G. Shock-wave evolution is assumed to follow the behavior of the coasting and adiabatic Sedov solutions. The evolution of the blast wave momentum $P(x) = \beta(x)\Gamma(x)$ in the adiabatic (Sedov) regime is described by the function

$$P(x) = \frac{P_0}{\sqrt{1+(x/x_d)^3}} \cong \begin{cases} P_0\;, & x \ll x_d \\ P_0(\frac{x}{x_d})^{-3/2}\;, & x_d \ll x \end{cases} \quad (6)$$

(Dermer and Humi, 2001), where $P_0 = \sqrt{\Gamma_0^2 - 1} = \beta_0\Gamma_0$ is the initial blast-wave Lorentz factor that defines the baryon loading. The deceleration radius

$$x_d \equiv [\frac{3(\partial E_0/\partial\Omega)}{\Gamma_0^2 m_p c^2 n_0}]^{1/3} \cong 2.1(\frac{m_\odot}{\Gamma_0^2 n_0})^{1/3}\;{\rm pc} \quad (7)$$

(Mészáros and Rees, 1993) is defined in terms of the directional energy release $m_\odot$ of the explosion, expressed in units of Solar rest-mass energy per 4π sr. Thus $m_\odot = 1 \Leftrightarrow \partial E_0/\partial\Omega = M_\odot c^2/4\pi = 1.4\times 10^{53}$ ergs sr^{-1}. For a nonrelativistic explosion, $x_d \cong \ell_{\rm S}$, the Sedov length (e.g., Sari et al., 1996), where $\ell_{\rm S} \equiv \Gamma_0^{2/3} x_d$.

The maximum particle energy for a particle accelerated from rest starting near the location of the explosion is found by integrating the energy gain rate per unit distance over the evolution of the blast wave radius. For q-$\|$ and q-$\perp$ nonrelativistic shocks, we obtain

$$E_{\|,max} \;\simeq\; 6\times 10^{15} Z\beta_0(\frac{m_\odot}{n_0})^{1/3} B_{\mu G}\;{\rm eV} \quad (8)$$

and

$$E_{\perp,max} \;\simeq\; 10^{16} Z\beta_0^{2/3}(\frac{m_\odot}{n_0})^{1/3} B_{\mu G}\;{\rm eV}\;, \quad (9)$$

respectively, where Ze is the particle charge.

3 Relativistic Shocks

It is simplest to work in the frame of the upstream medium to calculate the maximum energy of particles accelerated by a relativistic external shock. In the first cycle, a particle increases its energy by a factor Γ^2 but, due to kinematics of escape and capture, the particle increases its energy by only a factor ≈ 2 in successive cycles, as shown by Gallant and Achterberg (1999). This is because a particle is captured by the advancing relativistic shock when the particle is deflected by an angle $\theta_d \sim 1/\Gamma_s$, so that the upstream cycle time is $t_u \sim (\Gamma_s\omega_{\perp -})^{-1}$. Here the particle gyration frequency in the upstream medium is $\omega_{\perp -} = v/r_{\rm L-}$, where v is the particle speed and the particle deflection is determined by $B_{\perp -}$, the transverse component of the magnetic field in the upstream region. As a result, the nonthermal particle distribution function at a relativistic shock is highly anisotropic (Kirk and Schneider, 1987). As shown by Gallant and Achterberg (1999), the cycle time is dominated by the upstream transit time,

Because the energy increases by a factor $\cong 2$ during a single cycle, the momentum gain rate following the first cycle in relativistic shock acceleration is

$$\dot{p}_{rel} \simeq \frac{2p}{t_u} \simeq \frac{2cqB_{\perp -}\Gamma_s}{mc^2} \ , \qquad (10)$$

provided $\Gamma_s \gg 1$. If the particles are captured in the first cycle with Lorentz factor $\bar{\gamma}$, then the energy from the first cycle of Fermi acceleration in the stationary frame is $\approx \Gamma^2\bar{\gamma}mc^2$ (Vietri, 1995). After integrating the energy gain rate, we find that the maximum energy that can be achieved by relativistic first-order shock acceleration at an external shock is therefore

$$E_{rel,max} \simeq \bar{\gamma}\Gamma^2(x_0)mc^2 + 3\times 2^{3/2}qB_{\perp -}x_d\Gamma_0$$

$$\simeq [8\times 10^{13}\bar{\gamma}\Gamma_{300}^2 A + 10^{17}ZB_{\mu {\rm G}}(\frac{m_\odot \Gamma_{300}}{n_0})^{1/3}] \ \ {\rm eV} \ , \qquad (11)$$

whereAm_p is the particle mass, and $\Gamma = 300\Gamma_{300}$. From equation (11), we see that it is difficult to accelerate particles to energies larger than a factor $\sim 2^{1/3}\Gamma_0^{1/3}$ over $E_{\perp,max}$, given by equation (9), unless a pre-existing energetic particle distribution is found in the vicinity of the explosion.

4 Gyroresonant Stochastic Acceleration in Shocks

We derive the particle acceleration rate through stochastic gyroresonant processes from the expression $\dot{p}'_{\rm FII} \cong \Delta p'/t_{iso}$, where t_{iso} is the pitch-angle isotropization time scale in the comoving fluid frame. For relativistic hard sphere scatterers, the fractional change in momentum over time period t_{iso} is $\Delta p'/p' = 4p_A^2/3$ (Gaisser, 1990). The quantity p_Ac, which represents the dimensionless momentum of the scattering centers, reduces to the Alfvén speed in the weakly turbulent quasilinear regime. For relativistic shocks, the maximum value that p_A^2 can take is given by $p_A^2 \cong 2e_B/3$, as can be shown by considering the relativistic shock jump conditions (Blandford and McKee, 1976). The downstream magnetic field $B_+ \cong \max(\chi\Gamma B_{\perp -}, B_*\sqrt{\Gamma^2-\Gamma}\ )$, where $B_* \equiv (8\pi\chi n_0 m_p c^2 e_B)^{1/2} \cong 0.39(\frac{\chi}{4}n_0 e_B)^{1/2}$ G. The first term represents the compression of the upstream transverse magnetic field, and the second term defines the downstream field in terms of the equipartition field $B_{eq}({\rm G}) = 0.39(e_B n_0)^{1/2}$ $\sqrt{\Gamma^2-\Gamma}$ through the parameter e_B. For this convention, $\chi = 4$ for strong shocks whether or not the shocks are relativistic.

The particle pitch angle changes by $\delta B/B$ during one gyroperiod $t_{gyr} = r_{\rm L+}/c$, where B now refers to the mean downstream magnetic field in the shocked fluid. Thus $t_{iso} \cong t_{gyr}/(\delta B/B)^2$, noting that particles diffuse in pitch angle. Pitch angle changes due to gyroresonant interactions of particles with resonant plasma waves are described by the relation $(\delta B/B)^2 \approx \bar{k}W(\bar{k})/U_B$, where $U_B = B^2/8\pi$, and the turbulence spectrum $W(k) = W_0(k/k_{min})^{-v}$ for $k_{min} \leq k < k_{max}$. The index $v = 5/3$ for a Kolmogorov turbulence spectrum, and $v = 3/2$ if the Kraichnan phenomenology is used. The resonant wavenumber is normally assigned through the resonance condition $\omega - k_\parallel v_\parallel = \ell\Omega/\gamma$, but here we proceed by employing the simple resonance assumption $k \rightarrow 1/r_{\rm L}+$ (Biermann and Strittmatter, 1987). Assuming isotropy of forward and backward-moving waves gives the normalization $W_0 = \xi U_B(v-1)/2k_{min}$, where ξ is the ratio of plasma turbulence to magnetic field energy density. Hence $t_{iso}^{-1} \cong c\xi(v-1)(r_{\rm L+}k_{min})^{v-1}/(2r_{\rm L+})$, and

$$\dot{p}'_{\rm FII} \cong \frac{2}{3}p_A^2\ \xi(v-1)\ (\frac{c}{r^o_{\rm L+}})\ (r^o_{\rm L+}k_{min}p')^{v-1} \qquad (12)$$

(e.g., Dermer et al. (1996)).

The term $r^o_{\rm L+}k_{min}p'$ in the parentheses of equation (12) gives the comoving gyroradius in units of the inverse of the smallest turbulence wavenumber k_{min} found in the shocked fluid. If $k_{min} \sim 1/\Delta'$, where Δ' is the blast wave width, then the relation $r^o_{\rm L+}k_{min}p' \lesssim 1$ is an expression of the condition that the particle gyroradius is less than the size scale of the system (Hillas, 1984). When this condition is satisfied, $\dot{p}'_{\rm FII,max} \approx p_A^2 c/r^o_{\rm L+} \cong 2e_B c/(3r^o_{\rm L+})$.

By integrating the energy gain-rate for an adiabatic blast wave that evolves according to equation (6), we obtain the maximum energy-gain rates due to stochastic Fermi acceleration for nonrelativistic ($\Gamma - 1 \ll 1$) and relativistic ($\Gamma \gg 1$) shocks, given by

$$\frac{dE'}{dx}|_{FII,NR} = \frac{e_B\xi(v-1)}{2^{3/2}}\ qB_*\beta^2\ (\frac{2^{1/2}E'}{qB_*f_\Delta x\beta})^{v-1} \qquad (13)$$

and

$$\frac{dE'}{dx}|_{FII,ER} = \frac{2^{3/2}}{9}e_B\xi(v-1)\ qB_*\ (\frac{E'}{qB_*f_\Delta x})^{v-1}\ , \qquad (14)$$

respectively. Here f_Δ sets the scale for the smallest turbulence wave numbers in comparison with the inverse of the comoving size scale x/Γ of the blast wave width. From hydrodynamic considerations, $f_\Delta = 1/12$.

For an adiabatic blast wave that explodes in a uniform surrounding medium, the nonrelativistic expression (13) can be integrated to give the maximum energy

$$E_{max,FII,NR} \cong K_N\ qB_*f_\Delta x_d\beta_0\ . \qquad (15)$$

where

$$K_N \equiv [\frac{2^{v/2}}{4}\frac{e_B\xi}{f_\Delta}\ (v-1)(2-v)\beta_0 I_{2N}]^{1/(2-v)} \rightarrow$$

$$\begin{cases} 2\times 10^{-7}\ [\frac{e_B\xi\beta_{-2}(I_{2N}/6)}{f_\Delta}]^3\ , & \text{for } v = 5/3 \\ 4\times 10^{-5}\ [\frac{e_B\xi\beta_{-2}(I_{2N}/6)}{f_\Delta}]^2, & \text{for } v = 3/2 \end{cases} \ . \qquad (16)$$

Here $\beta_0 = 10^{-2}\beta_{-2}$, and I_{2N} is an integral that takes a maximum value $\simeq 6$.

The maximum energy of particles accelerated by second-order Fermi acceleration in a nonrelativistic blast wave is defined by the quantity

$$qB_* f_\Delta x_d \beta_0 \cong 8\times 10^{18} Ze_B^{1/2}(n_0)^{1/6} f_\Delta \beta_{-2} m_\odot^{1/3} \text{ eV} , \quad (17)$$

though the actual maximum energy is much less than this value when $\beta_0 \ll 1$, as can be seen from equation (16).

The relativistic expression (14) can be integrated to give

$$E_{max,FII,ER} \cong [\frac{2^{3/2}}{9}\frac{e_B\xi(v-1)}{f_\Delta}]^{\frac{1}{2-v}} qB_* f_\Delta x_d \Gamma_0 . \quad (18)$$

The maximum energy of particles accelerated by second-order Fermi acceleration in a relativistic blast wave occurs near the deceleration radius $x \approx x_d$. The basic scaling for this maximum energy is given by the term

$$qB_* f_\Delta x_d \Gamma_0 \cong 7.7\times 10^{20} Ze_B^{1/2} n_0^{1/6} f_\Delta (m_\odot \Gamma_0)^{1/3} \text{ eV} . \quad (19)$$

Note the sensitive dependence of the maximum energy given by equation (18) on the factor $e_B\xi/f_\Delta$. If $e_B\xi/f_\Delta > 1$ when the blast wave is near the deceleration radius, then stochastic acceleration processes in relativistic blast waves can in principle accelerate particles to ultra-high energies.

5 Discussion

The maximum particle energy that can be achieved through first-order Fermi acceleration of a particle initially at rest, for an adiabatic blast wave in a uniform surrounding medium, is given by

$$E_{max,1} \simeq 10^{16} ZB_{\mu \mathrm{G}} \beta_0^{2/3} (\frac{m_\odot \Gamma_0}{n_0})^{1/3} \text{ eV} . \quad (20)$$

Here we use the q-$\perp$ result for nonrelativistic blast waves, which gives the largest energies. The maximum particle energy that can be achieved through second-order Fermi acceleration in a relativistic blast wave is given by

$$E_{max,2} \simeq 8\times 10^{20} K_v Ze_B^{1/2} n_0^{1/6} f_\Delta (m_\odot \Gamma_0)^{1/3} \text{ eV} , \quad (21)$$

where $K_v = [2^{3/2}e_B\xi/9f_\Delta]^{1/(2-v)}$. The strong dependence of the maximum energy on β_0 makes second-order processes unimportant in nonrelativistic flows when $\beta_0 \ll 0.1$. As the flow speeds become marginally relativistic, first-order processes may accelerate particles to sufficiently high energies that second-order acceleration then starts to become important (Schlickeiser, 1984). At relativistic shocks, the second-order process can easily dominate.

Spectral models of GRBs within the external shock scenario (Mészáros and Rees, 1993; Chiang and Dermer, 1999) imply $\Gamma_{300} \sim 1$ and predict the existence of clean and dirty fireballs which have not yet been detected (Dermer et al., 1999). Mildly relativistic outflows are deduced from radio observations of the Type Ic SN 1998bw associated with GRB 980425 (Weiler et al., 2000; Kulkarni et al., 1998), so that hadronic cosmic ray acceleration to energies above the knee of the cosmic ray spectrum is possible from Type Ic SNe.

First-order Fermi processes in relativistic flows can accelerate particles to energies above the knee of the cosmic-ray spectrum. The smooth cosmic ray spectrum between the knee and the ankle suggests that the cosmic rays with energies above the knee originate from sources that accelerate particles to energies $\gtrsim 10^{19}$ eV. The second-order process in relativistic shocked fluids can accelerate particles to ultra-high energies to form both the galactic halo cosmic rays with energies between the knee and the ankle, as well as the metagalactic UHECRs. The GeV cosmic rays detected near Earth are probably a mixture of accelerated particles produced by many sources, but the cosmic rays formed near and above the knee may be produced by only one or a few energetic explosions (Erlykin and Wolfendale, 2000). GRBs are a prime candidate for the origin of these cosmic rays (Milgrom and Usov, 1996; Dermer and Humi, 2001).

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