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ELECTRON ACCELERATION AND EFFICIENCY IN NONTHERMAL GAMMA-RAY SOURCES

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ABSTRACT

In energetic nonthermal sources such as gamma-ray bursts, active galactic nuclei, or galactic jets, etc., one expects both relativistic and transrelativistic shocks acompanied by violent motions of moderately relativistic plasma. We present general considerations indicating that these sites are electron and positron accelerators leading to a modified power-law spectrum. The electron (or $e^{\pm}$) energy index is very hard, $\propto \gamma^{-1}$ or flatter, up to a comoving frame break energy γ_*, and becomes steeper above that. In the example of gamma-ray bursts, the Lorentz factor reaches $\gamma_* \sim 10^3$ for $e^{\pm}$ accelerated by the internal shock ensemble on subhydrodynamical timescales. For pairs accelerated on hydrodynamical timescales in the external shocks, similar hard spectra are obtained, and the break Lorentz factor can be as high as $\gamma_\star \lesssim 10^5$. Radiation from the nonthermal electrons produces photon spectra with shapes and characteristic energies in qualitative agreement with observed generic gamma-ray burst and blazar spectra. The scenario described here provides a plausible way to solve one of the crucial problems of nonthermal high-energy sources, namely, the efficient transfer of energy from the proton flow to an appropriate nonthermal lepton component.

Subject headings: acceleration of particles — cosmic rays — galaxies: active — galaxies: jets — gamma rays: bursts — gamma rays: theory — shock waves

1. INTRODUCTION

Shocks and, more generally, systems of shocks, as well as turbulent flow downstream from such shocks, may be a common feature of a number of nonthermal gamma-ray sources. A recent example of interest is the dissipative relativistic fireball model of gamma-ray bursts (GRBs), which follows from very general energetic and observational constraints, independently of whether the sources are at cosmological or Galactic halo distances. This implies a highly relativistic outflow of matter and electromagnetic fields lasting on the order of seconds where both external and internal shocks are expected (e.g., Rees & Mészáros 1992, 1994; Narayan, Paczyński, & Piran 1992; Paczyński & Xu 1994), providing a scenario in good qualitative agreement with major observational requirements (e.g., Mészáros, Rees, & Papanathanassiou 1994; Mészáros & Rees 1994). Under the cosmological assumption (which we henceforth assume), likely energy sources may be, for example, the coalescence of a compact binary (Narayan et al. 1992) or a failed supernova-like collapse (Woosley 1993). While the blast wave propagating into the external medium is highly relativistic, the reverse shock propagating back into the ejecta is likely to be only moderately relativistic. In addition, the irregular nature of the primary energy release results in the formation of a complex internal structure of faster and slower portions of the flow leading to internal shocks having moderate Lorentz factors ($\sim$1) in the comoving frame of the wind. Hydrodynamic (Waxman & Piran 1994) or MHD (Thompson 1994) turbulence may be expected in such scenarios and could play a role in the flow dynamics. In active galactic nucleus (AGN) jets one also expects internal shocks, as well as termination shocks, and, interestingly, the gamma-ray spectrum of blazars is qualitatively similar to that of GRBs. Such systems of shocks and turbulent regions provide an environment similar to those thought to lead to efficient particle acceleration (see, e.g., Blandford & Eichler 1987 and Jones & Ellison 1991 for a review).

In the case of GRBs, which we take as a generic example in this paper, the shocks are expected to energize an interaction region of spatial scale $\Delta \sim ct_{\rm var}\Gamma$ in the wind comoving system, where in the case of internal shocks, $t_{\rm var}$ is the timescale of energy release fluctuations (or, in the case of the reverse external shock, it is the light crossing time over the energy deposition region), and Γ is the mean bulk Lorentz factor of the flow. A multiple-shock structure is likely to arise as a result of the energy release fluctuations, the reflection and intersection of shocks crossing finite shells, or an inhomogeneous outflow. Strong, smooth relativistic MHD fluctuations produced by the irregular flow motions are expected to arise and interact with particles accelerated in such shocks. While details are uncertain, it is possible to consider some fairly general features of such flow collision regions (FCRs), encompassing an ensemble of internal shocks and developed turbulent motions, of maximum comoving length scale Δ and occurring at a lab frame distance $r_d \sim ct_{\rm var}\Gamma^2$ from the center of the event. Large-amplitude variations of the bulk velocity and the magnetic field can be expected on comoving length scales $l = \alpha\Delta$, with $\alpha < 1$ inversely proportional to the number of shocks. For the multiple secondary shock generation expected from flow collisions, we estimate $\alpha \lesssim 10^{-1}$.

In addition to such larger scale hydrodynamic substructure, we suggest that instabilities and nonlinear effects provide some

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sort of cascading process leading also to a wide spectrum of MHD fluctuations on smaller scales. Note that small-scale MHD motions may have nonrelativistic bulk velocities (unlike the large-scale motions containing most of the energy). Such violent systems are favorable sites for nonthermal particle acceleration. The nearest nonrelativistic analogs are the corotation interaction regions of the solar wind, which have been known for a long time to be an efficient source of low-energy cosmic rays (e.g., McDonald et al. 1974). Similar systems were suggested by Rees (1978) as sites for energetic particle acceleration in AGN jets.

In the next section we consider the process of nonthermal electron and positron spectrum formation in FCRs. While the specific examples and numerical values refer to cosmological gamma-ray bursts, similar considerations apply to galactic halo burst models and may also be of relevance to other astrophysical (e.g., AGN or galactic jet) sources involving relativistic flows and shocks.

2. NONTHERMAL LEPTON SPECTRA

We consider here a nonequilibrium process for the transformation of the power in the baryon bulk motions of the relativistic wind to nonthermal leptons and the temporal evolution of the lepton spectra. Turbulent plasma motions and shocks have been considered as generic sources of nonthermal particles since Fermi's pioneering work on statistical acceleration (e.g., Blandford & Eichler 1987). There are three important timescales in our problem. The first one is of the order of the cyclotron gyration time of relativistic $e^{\pm}$, which is characteristic of the fast preacceleration process taking place at shock fronts. The spectra of leptons accelerated by single shocks may be steeper than what is needed to explain directly gamma-ray burst photon spectra. Thus we consider the individual shocks as injection agents providing superthermal particles that can be further subject to diffusive acceleration through scattering on resonant fluctuations and large-scale MHD plasma motions. The latter, as shown below, can produce very hard lepton spectra on longer, subhydrodynamical and hydrodynamical timescales.

2.1. *Cyclotron Timescale*

Superthermal particles can be naturally extracted from the thermal pool by collisionless shocks, and this process operates on cyclotron timescales. The microscopic physics of particle acceleration in relativistic shocks is very complicated, and several fundamental aspects remain unclear, although some important results highlighting the distinctive features of such shocks have been obtained, e.g., Hoshino et al. (1992). These authors find that, for transverse relativistic shocks, pair acceleration to nonthermal energies occurs if the upstream flow contains ions carrying most of the energy flux, and they speculate that this may extend also to electron-proton plasmas. (Such conditions are typical of GRBs, where a proton component is expected and pairs are present but do not dominate the energy density). Hoshino et al. (1992) obtain a downstream nonthermal pair distribution $N \propto \gamma^{-2}$, where N is the number of particles within the range $d\gamma$, and their results indicate that 0.1–0.2 of the upstream baryon flow energy goes into magnetosonic waves that accelerate the nonthermal pairs. Microscopic simulations of relativistic quasi-parallel shocks are not yet available, but for nonrelativistic shocks in proton-electron plasmas an injection fraction $\zeta \sim 10^{-3}$ is typical (e.g., Giacalone et al. 1992).

2.2. *Subhydrodynamical Timescale*

A common attribute of any Fermi-type acceleration process is the isotropization of the fast lepton distribution due to the scattering by magnetic field fluctuations. From momentum conservation, one infers that to provide efficient particle scattering the magnetic field fluctuations must have an energy density comparable to that of the fast particles in the FCR rest frame. Resonant scattering of superthermal leptons will also be accompanied by stochastic acceleration with a typical timescale $\tau_a^{\rm st} \propto (c/u_{\rm ph})^2(\lambda/c)$, where $u_{\rm ph}$ is the phase velocity of the waves resonating with the scattered particle. For the mean free path $\lambda(\gamma)$ of a relativistic charged lepton within such a system of superthermal particles, the fluctuation spectrum can be approximated as being continuous on scales much larger than the electron gyroradius, as a result of strong dissipation. For a broad range of magnetic field fluctuations with spectral energy density $W(k) \propto k^{-\mu}$, one expects $\lambda \propto l(r_g/l)^{2-\mu}$ from standard quasi-linear theory (e.g., Blandford & Eichler 1987), where $r_g = 1.6 \times 10^3 \gamma B^{-1}$ cm is the electron gyroradius. The use of quasi-linear theory for the effects of fluctuations on scales $\leq l$ is justified because of the relatively small amplitudes of the resonant fluctuations (see also Hoshino et al. 1992) in the subhydrodynamic MHD regime. Thus, here we calculate the temporal evolution of the $e^{\pm}$ spectrum using a standard Fokker-Planck treatment.

This process leads to a nonthermal electron (or $e^{\pm}$) spectrum of the form $N(\gamma) \propto \zeta n \gamma^{1-\mu}$, which develops on a subhydrodynamic timescale of the order of a few $\tau_a^{\rm st}$, which for $u_{\rm ph} \sim c$ is

$$\tau_a^{\rm st} \sim (r_g/l)^{2-\mu} l/c. \tag{1}$$

Typically, $(r_g/l)^{2-\mu} \ll 1$ for the energies of interest (except for the case $\mu = 2$). From energy conservation, this energy spectrum extends up to

$$\gamma_* \sim [\gamma_i (m_p/m_e) \epsilon \zeta^{-1}]^{1/(3-\mu)}, \tag{2}$$

where γ_i is the initial proton (lepton) Lorentz factor, $\epsilon < 1$ is the portion of the total upstream power in baryons pumped into turbulent fluctuations, and ζ is the lepton injection fraction. The difference between the steeper γ^{-2} spectrum injected near the shock and the harder $\gamma^{1-\mu}$ spectrum produced by the MHD fluctuations arises because of the different spatial extent of the acceleration regions. In the former, the acceleration and escape timescales are comparable, $t_a \sim \kappa/u^2$, where u is the flow velocity and $\kappa \sim v\lambda$ is the diffusion coefficient, with $v \sim c$ the particle velocity, while $t_{\rm esc} \sim \delta^2/\kappa$, where $\delta \sim (v/u)\lambda \sim \kappa/u$ is the width of the particle acceleration region near the shock, so $t_a^{-1} t_{\rm esc} \sim u^2\delta^2/\kappa^2 \sim 1$. For scattering by MHD fluctuations between the shocks, however, the acceleration time is much shorter than the escape time: $t_a \sim (r_g/l)^{2-\mu}(l/c) \sim \lambda/c$, while $t_{\rm esc} \sim \Delta^2/\kappa \sim (\Delta/\lambda)^2(\lambda/c) \gg (\lambda/c)$.

Since typical MHD turbulent spectra have indices μ with $1.5 \leq \mu \leq 2$, the resulting particle spectra are $\propto \gamma^{-1}$ or flatter. We conclude that a substantial portion of the power in turbulence and some portion $\sim\epsilon$ of the upstream baryon power in the flow is transferred to the charged leptons at energies near γ_*. An estimate of ϵ and ζ would depend on uncertain details, such as the ratio of the matter to antimatter content in the flow, wave modes supported by the plasma, etc.

In the absence of other information, we will adopt here $\zeta \sim 10^{-3}$ and $\epsilon \sim 0.1$, compatible with the numerical results of Giacalone et al. (1992) and Hoshino et al. (1992). Taking as an example $\mu = 1.5$, an estimate of the break energy for the above typical parameters of the model gives $\gamma_* \sim 3 \times 10^3$.

Synchrotron losses of relativistic pairs have a timescale $\tau_{\rm syn} \approx 5 \times 10^8 B^{-2}\gamma^{-1}$ s, if B is measured in gauss. Thus for the formation of a hard branch of the pair spectrum up to the Lorentz factor γ_* by resonant acceleration, the scale l must satisfy the condition $l \leq 10^{(3\mu+12)/(\mu-1)} B^{\mu/(1-\mu)} \gamma^{(\mu-3)/(\mu-1)}$ cm. This implies, for $B = 10^4$ G and $\gamma_* \sim 10^3$, the scale $l \leq 10^7$ cm if the turbulent fluctuations have an index $\mu = 2$, and $l \leq 10^{12}$ cm if $\mu = 1.5$.

On the subhydrodynamical timescale, the distribution of nonthermal pairs for $\gamma \gg \gamma_*$ will be highly intermittent, with nonthermal pairs at these energies concentrated in the vicinity of the shocks, since they lose their energy before being mixed within the FCR. The synchrotron photon spectrum of the system beyond the break might be dominated by the brightest spots from some particular shock or the superposition of contributions from a few shocks. The resulting spectral shape just beyond the break may be rather complicated, reflecting with some modifications the injection spectrum.

2.3. *Hydrodynamical Timescales*

In addition to the above effects, one can also expect acceleration from processes occurring on the longer hydrodynamical comoving timescales of the order of l/c. The electric fields induced by turbulent motions of plasmas carrying magnetic fields on different scales lead to statistical energy gains of the superthermal charged particles. For nonrelativistic MHD turbulence, the particle energy change over a turbulent correlation length (or correlation time) is small, because the induced electric field is smaller than the entrained magnetic field. However, the distinctive feature of statistical acceleration by the *relativistic* MHD turbulence and shocks on larger scales expected in the FCR is the possibility of a substantial particle energy change over one correlation scale, because the induced electric fields are no longer small. In this case a Fokker-Planck approach cannot be used. Instead, we argue here that it is possible to calculate the energy spectra of nonthermal particles within FCRs using an integrodifferential equation that is a generalization of the Fokker-Planck approach (for details see the review by Bykov & Toptygin 1993, hereafter BT).

Consider charged test particles interacting with a wide spectrum of MHD fields and an internal shock ensemble produced by the colliding flows within generalized FCRs. In the wind comoving frame, we can assume the fluctuations on all scales up to $\sim\Delta$ (including the internal shock ensemble) to be nearly isotropic (for the latter, it is enough if they are forward-backward symmetric). The small mean free path λ of the superthermal particles leads to their isotropy in the frame of the local bulk velocity fluctuations. The assumed statistical isotropy of the bulk velocity fluctuations in the comoving frame of the wind results then in a nearly isotropic particle distribution, after averaging over the ensemble of internal shocks and accompanying motions on scales $\sim l$.

To calculate the spectrum of nonthermal leptons accelerated by the ensemble of internal shocks and large-scale plasma motions in the FCR (averaged over the statistical ensemble of large-scale motions), we use a kinetic equation for the nearly isotropic distribution function $N = \gamma^2 F$, which takes into account the non–Fokker-Planck behavior of the system (see Bykov 1991 and BT),

$$\frac{\partial F(\boldsymbol{r}, \xi, t)}{\partial t} = Q_i(\xi) + \int_{-\infty}^{\infty} d\xi_1 \, D_1(\xi - \xi_1)\Delta F(\boldsymbol{r}, \xi_1, t) + \left(\frac{\partial^2}{\partial \xi^2} + 3\frac{\partial}{\partial \xi}\right)\int_{-\infty}^{\infty} d\xi_1 \, D_2(\xi - \xi_1) F(\boldsymbol{r}, \xi_1, t). \quad (3)$$

Here $\xi = \ln(\gamma/\gamma_i)$, γ_i is the Lorentz factor of the injected particles, $Q_i(\zeta) \propto \zeta c n l^2$ is the rate of nonthermal particle injection, and n is the lepton number density in the FCR comoving frame. The kernels of integral equation (3) determining the spatial and momentum diffusion are expressed through correlation functions describing the statistical properties of the large-scale MHD turbulence and shock ensemble. Following the renormalization method, the Fourier transforms of the kernels $D_1^F(s)$ and $D_2^F(s)$ are solutions of a transcendental algebraic system of equations of the form $D_{1,2}^F = \Phi_{1,2}(D_1^F, D_2^F, s)$. Here s is a variable which is the Fourier conjugate of ζ. Equation (3) and the renormalization equations are valid only for particles with sufficiently small mean free paths $\lambda(\gamma) \ll \Delta$.

The crucial point is that the solution of equation (3) has a universal behavior, only weakly dependent on the complicated details of the turbulent system. The stationary solution to equation (3) with a monoenergetic injection rate Q_i has an asymptotic behavior of a power-law form, $N \propto Q_i \gamma^{-\sigma}$ (Bykov 1991), where $\sigma = -0.5 + [2.25 + \theta \mathrm{D}_1(0) D_2^{-1}(0)]^{0.5}$, and we took $\theta \sim (l/\Delta)^2$. For conditions typical of developed turbulence, the ratio of the rate of the scattering to the acceleration rate is $D_1(0)D_2^{-1}(0) < 1$ (see BT), and for $\theta < 1$ we obtain $\sigma \sim 1$. This hard γ^{-1} spectral behavior arises because the acceleration time $\tau_a \sim l/c \sim \alpha\Delta/c$ is much shorter than the escape time at the relevant energies, $\tau_{\rm esc} \sim \Delta^2/\kappa \sim \Delta^2/(lc) \sim \Delta/(\alpha c)$. The power needed to produce such a spectrum of nonthermal particles increases in proportion to $\gamma_{\rm max}$, so it is important to understand its temporal evolution.

In the test-particle limit, where the back-reaction of the accelerated leptons on the energy-containing bulk motions is negligible, we have $N(\gamma, t) \propto \zeta n \gamma^{-1}$ for $\gamma \leq \gamma_\star(t)$, where $\gamma_\star(t) = \gamma_i \exp(t/\tau_a^h)$ and

$$\tau_a^h \propto l/c \sim \alpha(\Delta/c) \quad (4)$$

is the typical hydrodynamical acceleration timescale (see, e.g., BT), with $\gamma_i \sim$ a few, $\alpha < 1$, and Δ is the comoving width of the region energized by shocks. From the energy-balance equation, when the value $\gamma_\star(t) \sim \gamma_i(m_p/m_e)\epsilon\zeta^{-1}$ is reached, the growth must saturate, and the resulting spectrum consists of two branches. One is the hard spectrum $N(\gamma) \sim \zeta n \gamma^{-1}$, for $\gamma \leq \gamma_\star$, where

$$\gamma_\star \sim \gamma_i(m_p/m_e)\epsilon\zeta^{-1}. \quad (5)$$

For the typical values of our problem, $\gamma_i \sim 1$ and $\zeta \sim 10^{-3}$, so $\gamma_\star \sim 10^5$ (but it could be even larger, since $\epsilon \sim 1$ for large-scale plasma motions).

3. ELECTRON ENERGIZATION EFFICIENCY

We outline here the application of the above acceleration scenario to the dissipative fireball model of gamma-ray bursts. It has been argued in the Introduction that one can expect violent flow collision regions to form in the dissipative portion of the fireball evolution. We do not go here into a detailed discussion of the radiation physics, nor do we attempt to model GRBs in detail. Rather, we concentrate on broadly generic examples using typical values of the relevant physical quantities.

In the fireball wind models, FCRs might occur around radii $r \sim ct_{\rm var}\Gamma^2 \sim 10^{12}$–$10^{13}$ cm with bulk Lorentz factors $\Gamma \sim 10^2$ and mean comoving field strengths $B \sim 10^4 B_4$ G (e.g., Rees & Mészáros 1994). We assume the acceleration to occur beyond the region where significant pair formation is expected, i.e., outside the photosphere, and take as numerical examples a lepton injection fraction $\zeta \sim 10^{-3}$, turbulence energy fraction $\epsilon \sim 10^{-1}$, and initial injection Lorentz factor $\gamma_i \sim 1$. We also assume a broad spectrum of MHD- or whistler-type fluctuations with index $\mu = 1.5$. The characteristic timescale of the fast injection process of leptons is a few microseconds in the FCR comoving frame. Then, on a subhydrodynamical timescale $\tau_a^{\rm st} \sim 10c^{-1}(lr_g)^{1/2} \sim 10c^{-1}(\alpha ct_{\rm var}\Gamma r_g)^{1/2}$, a hard lepton spectrum is established, with $N(\gamma) \propto \zeta n\gamma^{-0.5}$ for $\gamma \leq \gamma_* \sim [\gamma_i(m_p/m_e)\epsilon\zeta^{-1}]^{2/3}$. The comoving frame time $\tau_a^{\rm st} \lesssim$ a few milliseconds is enough to transfer a fraction $\sim\epsilon$ of the baryonic power to the accelerated leptons, with Lorentz factors $\gamma \sim \gamma_* \sim 10^3$. For steeper turbulent fluctuation spectra ($\mu \gtrsim 1.5$), the timescale is somewhat longer but still compatible with observational requirements. As they are accelerated, the leptons radiate a synchrotron spectrum of the form $\nu F_\nu \sim \nu^{1.25}$, up to a break near 0.1 MeV for the parameters used, and $\propto\nu^k$ above that, with k between 0 and -1 if rather weak shocks dominate in the FCR.

Nonthermal lepton acceleration on the longer, hydrodynamical timescales might be important in GRB external shocks around radii $r \sim 10^{16}$ cm (e.g., Mészáros et al. 1994). Typical magnetic fields in the reverse shock could be $\sim$10 G, while the hydrodynamic timescale is $\tau_a^h \sim \alpha\Delta/c \sim$ tens of seconds in the FCR comoving frame (and Γ^{-1} times shorter in the lab frame). This is enough to form a pair spectrum $N(\gamma) \propto \zeta n\gamma^{-1}$ up to Lorentz factors $\gamma_\star \sim 2 \times 10^5$, for these parameters, and a synchrotron spectrum with $\nu F_\nu \sim \nu$ and a peak energy near the 0.5 MeV range.

Recent *GRO* observations of blazars (McNaron-Brown et al. 1995) show clear evidence for broken power-law spectra peaked in the MeV range with shapes similar to gamma-ray burst spectra (e.g., Greiner et al. 1995). From the similarity of the acceleration scenarios expected in both, involving relativistic shocks and turbulence, one may speculate on the possible applicability of the above processes to explain blazar spectra.

In principle, an injected fraction $\zeta < 1$ of protons may also be accelerated by the same mechanisms. The maximum proton Lorentz factors would in this case be $\gamma_{p,*} \sim (\epsilon\gamma_i/\zeta)^{1/(3-\mu)}$ or $\gamma_{p,\star} \sim (\epsilon\gamma_i/\zeta)$, which for $\epsilon \sim 10^{-1}$ and $\zeta \sim 10^{-3}$, could be as high as 10^4. However, the fraction of postshock proton energy going into such a flat, nonthermal relativistic proton component is at most ϵ, comparable to the fraction of energy going into the nonthermal flat lepton spectrum.

The efficiency of the transfer of energy from the proton to the lepton component in these models is high, typically of order $e_{pe} \sim (\zeta/\gamma_i)(m_e/m_p)\gamma_e^{3-\mu} \lesssim \epsilon$ for subhydrodynamic, and $e_{pe} \sim (\zeta/\gamma_i)(m_e/m_p)\gamma_e \lesssim \epsilon$ for hydrodynamic acceleration, where γ_e can go up to γ_* or $\gamma_\star$, and $\epsilon \lesssim 1$ is the fraction of upstream proton energy converted into turbulence in the semirelativistic wind and reverse blast wave shocks, γ_i is the initial Lorentz factor, and ζ is the lepton injection fraction. This high efficiency is due to the very hard lepton spectra achieved with $t_{\rm acc} \ll t_{\rm esc}$, $N(\gamma) \propto \gamma^{-1}$ or flatter, which puts most of their energy near the upper break value γ_* or $\gamma_\star$. This is a significant fraction of the equipartition value between the accelerated leptons and the bulk of the shocked protons. Similar shocks and turbulent regions are likely to be present in AGN or galactic jets. If these are electron-proton jets, as opposed to electron-positron jets, most of the energy is in the protons (as for the GRB case), and a significant fraction of it should be channeled into the electrons. Since the leptons, because of their smaller mass, are responsible for most of the radiation, this mechanism of proton-electron energy sharing fulfills a major prerequisite for a high radiative efficiency in GRBs and other nonthermal gamma-ray sources.

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