e\(^\pm\) PAIR LOADING AND THE ORIGIN OF THE UPSTREAM MAGNETIC FIELD IN GRB SHOCKS

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ABSTRACT

We investigate here the effects of plasma instabilities driven by rapid e\(^\pm\)\,-pair cascades, which arise in the environment of \(\gamma\)-ray burst (GRB) sources as a result of back-scattering of a seed fraction of the original spectrum. The injection of e\(^\pm\)\,-pairs induces strong streaming motions in the ambient medium. One therefore expects the pair-enriched medium ahead of the forward shock to be strongly sheared on length scales comparable to the radiation front thickness. Using three-dimensional particle-in-cell simulations, we show that plasma instabilities driven by these streaming e\(^\pm\)\,-pairs are responsible for the excitation of near-equipartition, turbulent magnetic fields. Our results reveal the importance of the electromagnetic filamentation instability in ensuring an effective coupling between e\(^\pm\)\,-pairs and ions, and may help explain the origin of large upstream fields in GRB shocks.

Subject headings: gamma rays: bursts — instabilities — magnetic fields — plasmas — shock waves

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1. INTRODUCTION

More than three decades ago, it was pointed out that \(\gamma\)-rays produced in sufficiently luminous and compact astrophysical sources would create e\(^\pm\) pairs by collisions with lower energy photons: \(\gamma\gamma \rightarrow e^+ e^-\) (Jelley 1966). This mechanism both depletes the escaping radiation and also changes the composition and properties of the radiating gas through the injection of new particles. An approximate condition for such pair creation to become significant is that a sizable fraction of the radiation from the object be emitted above the electron mass energy, \(\epsilon \gg m_e c^2\). In this case, the photons are beamed and is alleviated when the radiating source itself expands at a relativistic speed (Piran 1999). In this case, the photons are beamed into a narrow angle \(\theta \sim 1/\Gamma\) along the direction of motion, and, as a result, the threshold energy for pair production within the beam is increased to \(\epsilon \sim \Gamma\). The nonthermal spectrum of GRB sources is therefore thought to arise in shocks that develop beyond the radius at which the relativistic fireball has become optically thin to \(\gamma\gamma\) collisions (Piran 1999). However, the observed spectra are hard, with a significant fraction of the energy above the \(\gamma\gamma \rightarrow e^+ e^-\) formation energy threshold, and a high compactness parameter can result in new pairs being formed outside the originally optically thin shocks responsible for the primary radiation (Thompson & Madau 2000). Radiation scattered by the external medium, as the collimated \(\gamma\)-ray front propagates through the ambient medium, would be decollimated, and, as long as \(l \gtrsim 1\), absorbed by the primary beam. An e\(^\pm\)\,-pair cascade can then be produced as photons are back-scattered by the newly formed e\(^\pm\)\,-pairs and interact with other incoming seed photons (Thompson & Madau 2000; Beloborodov 2002, 2005; Mészáros et al. 2001; Ramirez-Ruiz et al. 2002; Li et al. 2003; Kumar & Panaitescu 2004).

In this paper, we consider the plasma instabilities generated by rapid e\(^\pm\)\,-pair creation in GRBs. The injection of e\(^\pm\)\,-pairs induces strong streaming motions in the ambient medium ahead of the forward shock. This sheared flow will be Weibel-like unstable (Medvedev & Loeb 1999), and if there is time before the shock hits, the resulting plasma instabilities will generate sub-equipartition quasi-static long-lived magnetic fields on the collisionless temporal and spatial scales across the e\(^\pm\)\,-pair-enriched region. This is studied in \(\S\, 4\) using three-dimensional kinetic simulations of monoenergetic and broadband pair plasma shells interpenetrating an unmagnetized medium. The importance of the electromagnetic filamentation instability in providing an effective coupling between e\(^\pm\)\,-pairs and ions is investigated in \(\S\, 2\), while a brief description of the numerical methods and the initial models is given in \(\S\, 3\). The implications for the origin of the upstream magnetic field, in particular in the context of constraints imposed by observations of GRB afterglows, are discussed in \(\S\, 5\).

2. PAIR LOADING AND MAGNETIC FIELD GENERATION

Given a certain external baryon density \(n_b\) at a radius \(r\) outside the shocks producing the GRB primordial spectrum, the initial
Thomson scattering optical depth is $\tau \sim n_e \sigma_T r$, and a fraction $\tau$ of the primordial photons will be scattered back, initiating an $e^\pm$-pair cascade. Since the photon flux drops as $r^{-2}$, for a uniform (or decreasing) external ion density, most of the scattering occurs between $r$ and $r/2$, and the scattering and pair formation may be approximated as a local phenomenon.

Consider an initial input GRB radiation spectrum of the form $F(\epsilon) = F_0(\epsilon/\epsilon_\text{th})^{-\alpha}$ for $\epsilon > \epsilon_\text{th}$, where $\epsilon_\text{th} \sim 0.2 - 1.0$ is the break energy above which the spectral index $\alpha \sim 1 - 2$ (for the present purposes, the exact low-energy slope is unimportant). Radiation scattered by the external medium is therefore decollimated, and then absorbed by the primary beam. The impact of such a process strongly depends on how many photons each electron is able to scatter (Beloborodov 2002);

$$\lambda_T \sim 4 \times 10^8 \left( \frac{r_T}{10^{15} \text{ cm}} \right)^2 \left( \frac{L}{10^{51} \text{ erg s}^{-1}} \right) \text{ cm} < \Delta, \quad (3)$$

where $r_T = 1/(n_e \sigma_T) = c(F_0 \sigma_T)$ is the electron mean free path, and $n_e$ is the photon density. The photons forced out from the beam by the electrons can then produce $e^\pm$ pairs as they interact with incoming photons (Fig. 1), so that a large number of scatterings implies a large number of pairs created per ambient electron.

At some distance $\varpi \leq \Delta$ from the leading edge of the radiation front, the number of photons scattered by one ambient electron is $\sim \varpi/\lambda_T$, and a fraction $\sim \varpi/\lambda_{\gamma \gamma}$ of these photons are absorbed (Beloborodov 2002). One $e^\pm$ pair per ambient electron is injected when

$$\varpi = (\lambda_T \lambda_{\gamma \gamma})^{1/2} \approx \lambda_T \epsilon_\text{th}^{1/2} / [\Pi(\alpha)]^{1/2} \geq \lambda_T,$$

where $\Pi(\alpha) = 2^{-\alpha} (7/12)(1 + \alpha)^{-5/3}$ (Svensson 1987) and $\epsilon_\text{th}$ is the photon threshold energy for $e^\pm$ formation. For $1.5 \leq \alpha \lesssim 2$, one has $15 \lambda_T \lesssim \varpi \lesssim 25 \lambda_T$, so that pair creation substantially lags behind electron scattering (Beloborodov 2002).

Most of the momentum deposited through this process involves the side-scattering of very soft photons, which collide with hard $\gamma$-rays to produce energetic (and almost radially moving) pairs. A photon of energy $\epsilon_\gamma \ll 1$ that is side-scattered through an angle $\sim \theta_e$ creates a pair if it collides with another photon with energy exceeding $\epsilon_\text{th} \sim 4(\theta_e^2 \epsilon_\gamma)^{-1}$. The injected pair will be relativistic with $\gamma_+ \sim \epsilon_\text{th} \geq 1$. The distribution of injected pairs is therefore directly determined by the high-energy spectral index $\alpha$. This motivates our study in § 4 of radially streaming, relativistic pair plasma shells with a broadband kinetic energy distribution interpenetrating an unmagnetized medium.

This pair-dominated plasma, as long as its density $n_\perp \lesssim n_e (m_\text{p}/2m_\text{e})$, is initially held back by the inertia of its constituent ions, provided that the pairs remain coupled to the baryons (Thompson & Madau 2000; Beloborodov 2002). The latter is likely to be the case in the presence of weak magnetic fields (Thompson & Madau 2000). In the absence of coupling, the pair density would not exponentiate, mainly due to the $(1 - \beta)$ term in the scattering cross section, where $\beta = v/c$. Instabilities caused by pair streaming relative to the medium at rest, as we argue in § 4, are able to generate long-lived magnetic fields on the collisionless temporal and spatial scales, which are modest multiples of the electron plasma frequency, $\omega_\text{e} = (4\pi e^2 n_e/m_e)^{1/2}$, and the collisionless skin depth,

$$\lambda_\text{e} = \frac{c}{\omega_\text{e}} \sim 5 \times 10^5 \left( \frac{n_e}{1 \text{ cm}^{-3}} \right)^{-1/2} \ll \lambda_T < \lambda_{\gamma \gamma}, \quad (4)$$
Here, $n_e$ is the electron number density. Plasma instabilities therefore evolve faster than the time between successive scatterings, and much before the scattered photons are absorbed by the primary radiation (Beloborodov 2002). In the presence of a transverse magnetic field $B$, the pairs gyrate around field lines on the Larmor time, $\omega^{-1}_B = m_e c / (Be)$. The net momentum of the $e^\pm$ pairs is thus efficiently communicated to the medium (Frederiksen et al. 2004; Hededal & Nishikawa 2005). Magnetic coupling may dominate if $B > \omega_e$, which requires $B^2 / 4\pi > n_e m_e c^2$ (Beloborodov 2002). In the sections that follow, we present a quantitative discussion of the effects of plasma instabilities driven by rapid $e^\pm$-pair injection.

3. SIMULATION MODEL

Here we illustrate the main features of the collision of an $e^\pm$-pair plasma shell into an unmagnetized medium, initially at rest, using a modified version of the PIC code TRISTAN, first developed by Buneman (1993) and most recently updated by Nishikawa et al. (2005, 2006). The simulations were performed on an $85 \times 85 \times 640$ grid (the axes are labeled as $x$, $y$, and $z$) with a total of $3.8 \times 10^8$ particles with periodic ($x$-$y$ plane) and radiative ($z$-direction) boundary conditions. In physical units, the box size is $8.9 \times 8.9 \times 66.7 \ (c/\omega_e)^3$, and the simulations ran for $60\omega_e^{-1}$.

In the simulations, a quasi-neutral plasma shell, consisting of $e^\pm$ pairs and moving with a bulk momentum $u_z = \gamma_0 v_z / c$ along the $z$-direction, penetrates an ambient plasma initially at rest. Here, $\gamma_0$ is the initial Lorentz factor of the pairs, and $v_z$ is the bulk velocity of the shell along $z$. The $e^\pm$ pairs are continuously injected at $z = 2.6 \lambda_e$, where $\Lambda = \lambda_e / 9.6$ is the grid size. The ion-to-electron mass ratio of the ambient plasma is set to $m_I/m_e = 20$, with both plasma populations having a thermal spread with an electron rms velocity $\nu_0 / c = 0.1$. For completeness, a purely $e^\pm$-pair ambient plasma (i.e., with a mass ratio $m_I/m_e$ of 1) is also studied, with a thermal electron velocity $\nu_0 / c = 0.1$. The shell and the ambient plasma have a density ratio of 0.75. Two different bulk Lorentz factor configurations for the injected (cold) $e^\pm$ pairs are considered: a monoenergetic ($u_z = 12.5$, $\nu_0 / c = 0.01$) and a broadband distribution ($3.0 \leq u_z \leq 30.0$, $\nu_0 / c = 0.01$). Both distributions have similar kinetic energy contents and plasma temperatures (Fig. 2).

4. RESULTS AND INTERPRETATION

Much of our effort in this section is dedicated to the ability of $e^\pm$-pair loading to alter the physical parameters characterizing the ambient medium. Some of the questions at the forefront of our attention include the level of the (electro)magnetic field generated via plasma instabilities, as well as the saturated state of the particles and fields. We address all of these issues here. We consider both $e^\pm$-pair and electron-ion plasmas as possible ambient media. In the dimensional units used, the results for $e^\pm$-pair ambient plasmas are equally applicable to interpenetrating proton/antiproton plasma shells.

4.1. $e^\pm$-Pair Ambient Plasma

As pointed out in Silva et al. (2003), encountering the medium at rest, the incoming pairs are rapidly deflected by field fluctuations, which grow because of the two-stream instability (Medvedev & Loeb 1999; Pruet et al. 2001; Gruzinov 2001). A large number of oppositely directed current filaments are generated, as illustrated in Figure 3, which in turn will generate inhomogeneities in the magnetic field (predominantly in the $x$-$y$ plane). This configuration is unstable because opposite currents repel each other, whereas like currents are attracted to each other, and tend to coalesce and form larger current filaments. During the linear stage of the instability, the rapid generation of a large number of randomly distributed current and magnetic filaments is observed [with a correlation length $\Lambda_f \approx 2^{1/4} \gamma_0 (c/\omega_e) = 1.2 \lambda_e = 10 \Lambda$], which in turn is accompanied by an expeditious production of a strong magnetic field (Fig. 3). The magnetic field energy density reaches 9% of the initial total kinetic energy ($\epsilon_p$) in the broadband case, and 1.2% in the monoenergetic one [i.e., $\epsilon_B = \int (B^2 dv/8\pi) / \epsilon_p = 0.009$ and 0.012, respectively].

As the instability enters the saturation phase, these initially randomly oriented filaments begin to interact with each other and merge in a race in which larger electron channels consume smaller neighboring channels. In this manner, the transverse magnetic field grows in strength. The instability then saturates, and the energy in the magnetic field decays. The strong decrease in the magnetic field energy density is associated with a topological adjustment in the structure of the currents and fields (Silva et al. 2003), and is mostly linked to a reduction of the field’s volume filling factor. These distinctive features are present in both the monoenergetic and broadband cases (Fig. 3). However, there are some clear contrasts. For a plasma shell with a broadband distribution of initial Lorentz factors, transverse energy spreading happens over a variety of timescales. This can be
Fig. 3.—Growth of the two-stream instability at time $t = 59.8\omega_c^{-1}$ for the two different momentum distribution functions shown in Fig. 2. Here we show the average transverse magnetic field amplitude (in simulation units) in the $x$-$y$ plane as a function of $z$ (solid curves). The various panels show the $(x, y)$ components of the magnetic field and the $z$-component of the current density ($J_z$) in different subsections of the computational box. The top (bottom) panels are for a simulation in which the initial momentum distribution function magnetic field is monoenergetic (broadband). The color bar gives the amplitude of $J_z = |J_{\max} - J_{\min}|$ in simulation units for $J_{\max} = 20.04$ (25.57) for $z = 150$, $J_{\max} = 99.5$ (67.12) for $z = 350$, and $J_{\max} = 35.4$ (29.93) for $z = 520$. The expected correlation length, $\lambda_J$, of the randomly distributed current and magnetic filaments is plotted in the $z = 350$ subsections.
understood as follows. The instability arises from the free streaming of particles, with a corresponding linear growth rate scaling with $\gamma_0^{1/2}$ (Medvedev & Loeb 1999). Our experiments show that this does indeed happen; the continuous injection of a broadband distribution of momentum leads to a diversity of linear growth rates. As the field amplitude grows, the transverse deflection of particles increases (with decreasing $\gamma_0$), and thus free streaming across the field is suppressed preferentially for the mildly relativistic pairs. The magnetic field energy grows in the early stages by slowing down the pair plasma shell.

Saturation is then achieved by the combination of transverse energy spreading and the generation of near-equipartition magnetic fields (Silva et al. 2003). In the broadband case, transverse energy spreading happens more swiftly for mildly relativistic pairs, and, as a consequence, $B_{\text{sat}}$ is decreased. This follows directly from the scaling of $B_{\text{sat}}$ with $v_{\text{th}}^{-1}$ (or in this case, $v_{\text{th}}^{-1}$) (Medvedev & Loeb 1999). As can be clearly seen in Figure 4, the ambient $e^\pm$ pairs are more effectively heated in the broadband case. After saturation, which takes place earlier for mildly relativistic pairs, the energy stored in the magnetic field is transferred back to the plasma particles, leading to strong heating and the generation of a high-energy tail in the distribution (Fig. 2). The $e^\pm$ pairs are nonetheless expected to thermalize, given sufficient space and time (Spitkovsky 2006).

It is not surprising that the linear growth rate saturates at comparable timescales (Fig. 2), since in the broadband distribution, the energy-averaged Lorentz factor, $\gamma_0$, is not very different from $\gamma_0$ in the monoenergetic case. Figure 3 also shows a transient, electromagnetic precursor produced as the first injected $e^\pm$ pairs stream through the ambient medium unhampered. The saturation...
cause mutual attraction between currents, forcing like currents to approach each other and merge. As a result, the magnetic field grows in strength. This continues until the fields grow strong enough to deflect the much heavier ions (Frederiksen et al. 2004). Figure 5 shows that this happens over \(~20\) electron skin depths from around \(z = 300\). As illustrated in Figure 5, the ions stay clearly separated in phase space, and are only slowly heated. In the presence of ions, the incoming \(e^\pm\) pairs will drive higher levels of saturated \(B\)-field (Fig. 5) by a factor of \((m_i/m_e)^{1/2}\), albeit on a longer timescale; the magnetic field energy density in the monoenergetic case reaches 6.1% of the initial total kinetic energy in the electron-ion ambient plasma case, and 1.2% in the \(e^\pm\)-pair case. This is due to the massive ion bulk momentum, which constitutes a vast energy reservoir for particle heating.

In the case of an electron-ion ambient plasma, the ion channels are subjected to a growth mechanism similar to that of the positrons, although they grow at a slower rate. When ion channels grow sufficiently powerful, they begin to experience Debye shielding by the electrons, which by then have been significantly heated by scattering on the increasingly large electromagnetic field structures. The temporal development of the \(e^\pm\)-pair and ion channels is illustrated in Figure 6. The large random velocities of the electron population allow the concentrated ion channels to continue to sustain strong magnetic fields.

Because the total plasma momentum must be conserved, the electromagnetic field produced by the two-stream instability acquires part of the longitudinal momentum lost by the counterstreaming populations. Consequently, the magnetic field energy grows in the early stage by slowing down the \(e^\pm\)-pair plasma, as illustrated in Figure 2. Our experiments show that this also happens for \(e^\pm\)-pair shells injected into an electron-ion ambient plasma (Fig. 7). However, there are some clear differences in the evolution of the associated absorption of momentum between the broadband and the monoenergetic cases, as is argued in § 4.1. Figures 8 and 9 elucidate some of these differences.

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**Figure 5**—Longitudinal heating and acceleration, illustrated by changes in \(u_i\) for both injected (monoenergetic) \(e^\pm\) pairs (gray) and ambient plasma (black). The bottom (top) panel is for a simulation in which the ambient medium is composed of electrons and ions (\(e^\pm\) pairs). In the top panel, the ambient positrons are initially at rest, but are strongly accelerated by the jet. In the bottom panel, the ions, being heavier than the pairs, remain clearly separated in phase space, and are only slowly heated. Also shown are the average transverse magnetic field amplitudes (in arbitrary units) in the \(x-y\) plane as a function of \(z\) (solid curves). [See the electronic edition of the Journal for a color version of this figure.]

**Figure 6**—Longitudinal current densities in the counterstreaming plasmas. The various panels show the \(z\)-component of the current density, \(J_z\), in different subsections of the computational box for \(e^\pm\) pairs injected with a monoenergetic distribution. The small insets show the ion current in the same plane. The color bar gives an amplitude of \(J_i = \left|J_{max} - J_{min}\right|\) in simulation units for \(e^\pm\) pairs (ions): \(J_{max} = 20.04\) (0.1) for \(z = 150\), and \(J_{max} = 173.3\) (8.1) for \(z = 350\). The expected correlation length, \(\delta_j\), of the randomly distributed current and magnetic filaments is plotted in the \(z = 150\) and 350 subsections.
As can clearly be seen in Figure 8, transverse energy spreading occurs faster in the broadband case, and, as a consequence, the ambient ion beam is heated more promptly. The total magnetic energy then grows as the ion channels merge. The magnetic field associated with these currents also has a filamentary structure, as seen in Figure 9. The magnetic energy scales with the square of the electric current, which in turn grows in inverse proportion to the number of current channels. The net result is that the mean magnetic energy increases accordingly: the magnetic field energy density reaches 6.1% (4.7%) of the initial total kinetic energy in the monoenergetic (broadband) case.

The ambient ions retain distinct bulk speeds in shielded ion channels and thermalize much more slowly (Fig. 8). This is expected to continue as long as a surplus of bulk relative momentum remains in the counterstreaming plasmas. Figure 9 shows the extension of the ion currents and the corresponding magnetic field energy density. After the linear stage, the instability saturates, and the energy in the magnetic field decays, reflecting essentially the subsequent decrease in the field’s volume filling factor. In the monoenergetic case, the magnetic energy density drastically drops after about $\sim 30 \omega_e^{-1}$. The decrease in the magnetic field energy is, however, slower in the broadband shell. This is associated with the continuous transverse energy spreading of progressively faster moving $e^\pm$ pairs. The transverse energy spreading is less efficient for the faster moving component and, as a result, a large number of oppositely directed current filaments continue to be generated, even after the bulk of the flow enters the saturation phase. In the monoenergetic case, on the other hand, the magnetic energy density drastically drops after about $\sim 30 \omega_e^{-1}$. These fields, however, maintain a strong saturated level for at least the duration of the simulations. The final magnetic energy density level is still quite high, up to 2.75% (2.57%) of the initial total kinetic energy in the monoenergetic (broadband) case.

This is in contrast to the results reported by Frederiksen et al. (2004) for relativistically counterstreaming electron-ion plasma collisions; the magnetic field energy density continues to grow throughout their experiment, which lasts for about $480 \omega_e^{-1}$. In these experiments, however, both counterstreaming plasmas are composed of electrons and ions, which in turn significantly increases the relative bulk momentum. It is thus clear that the symmetry associated with the counterstreaming electron-ion populations is essential for the longevity of the ion current channels.

5. CONCLUSIONS

We present self-consistent three-dimensional simulations of the fields developed by the electromagnetic filamentation instability as a consequence of $e^\pm$-pair injection, which arise naturally in the environment of GRB sources as a result of back-scattering. Our results demonstrate that even in an initially unmagnetized scattering plasma, a small-scale, fluctuating, predominantly transversal, and near-equipartition magnetic field is unavoidably generated. These fields maintain a strong saturated level on timescales much longer than $\lambda_e$, at least for the duration of the simulations $\sim 60 (c/\lambda_e) \sim 0.1 (c/\lambda_T)$. The $e^\pm$ pairs are effectively scattered with the magnetic field, and thus effectively communicate their momentum to the scattering medium initially at rest. Our results indicate that the fields necessary to ensure that $e^\pm$ pairs remain coupled to the medium can be easily created via plasma instabilities. The next required step is to increase the ion-to-electron mass ratio in the scattering medium in order to determine the spatial spread and character of the particle coupling.
A question that has remained largely unanswered so far is what determines the characteristic strength of the upstream magnetic field afterglow shocks, which is inferred to extend to tenths, or even tens, of mG (Li & Waxman 2006). This is, of course, large in comparison with the magnetic fields presumed to be present in the interstellar medium, which are measured in μG. A sufficiently strong magnetic field may be present in a wind ejected by the GRB progenitor, but only if it is highly magnetized, which is not thought to be the case for the widely favored Wolf-Rayet stars (Li & Waxman 2006). The injection of $e^\pm$ pairs induces strong streaming motions in the ambient medium ahead of the forward shock. This sheared flow, as we demonstrated here, will be Weibel-like unstable, and, if $e^\pm$ loading continues, the resulting upstream flow will be strongly magnetized well after the shock hits the $e^\pm$-pair-enriched medium.

Since one expects the radiation front to lead the forward shock (Beloborodov 2002, 2005) by a small distance, $R/(4\Gamma^2)$, one expects the instability to be able to generate subequipartition magnetic fields, $B \sim 0.1(n_p/1\text{ cm}^{-3})^{1/2}$ G, in the scattering medium just before it is shocked. In this case, the constraint on the upstream field strength of $B \geq 0.05(n_p/1\text{ cm}^{-3})^{3/8}$ mG, imposed by X-ray afterglow observations (Li & Waxman 2006), is easily satisfied.

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