Dissipation in Relativistic Outflows: A Multisource Overview

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Abstract: Relativistically expanding sources of X-rays and γ-rays cover an enormous range of (central) compactness and Lorentz factor. The underlying physics is discussed, with an emphasis on how the dominant dissipative mode and the emergent spectrum depend on these parameters. Photons advected outward from high optical depth are a potentially important source of Compton seeds. Their characteristic energy is bounded below by $\sim 1$ MeV in pair-loaded outflows of relatively low compactness, and remains near $\sim 1$ MeV at very high compactness and low matter loading. This is compared with the characteristic energy of $O(1)$ MeV observed in the rest frame spectra of many sources, including γ-ray bursts, OSSE jet sources, MeV Blazars, and the intense initial 0.1 s pulse of the March 5 event. Additional topics discussed include the feedback of pair creation on electron heating and the formation of non-thermal spectra, their effectiveness at shielding the dissipative zone from ambient photons, direct Compton damping of irregularities in the outflow, the relative importance of various soft photon sources, and the softening of the emergent spectrum that results from heavy matter loading. The implications of this work for X-ray and optical afterglow from GRB's are briefly considered. Direct synchrotron emission behind the forward shock is inhibited by the extremely low energy density of the ambient magnetic field. Mildly relativistic ejecta off axis from the main γ-ray emitting cone become optically thin to scattering on a timescale of $\sim 1$ day $(E/10^{52}$ erg$)^{1/2}$, and can be a direct source of afterglow radiation.

1 Introduction: Variety of Sources and Spectral Behavior

X-ray and γ-ray emission from relativistic outflows is powered by the conversion of bulk kinetic energy and Poynting luminosity, by a variety of possible mechanisms. However, the assumed values of key parameters such as the Lorentz factor $\gamma$ of the outflow, the compactness of the central source, as well as the optical depth and size of the dissipative zone, vary dramatically between the different classes of sources. For example, $\ell_c \sim 10^{15}$ and $\gamma \sim 10^2 - 10^3$ are inferred for cosmological γ-ray burst (GRB) sources vs. $\ell_c \sim 10 - 10^2$ and $\gamma \sim 3 - 30$ for Blazars (Fig. 1). This motivates a more global analysis of how dissipation of kinetic and magnetic...
energy is achieved, which provides some interesting new perspectives on particular sources.

![Diagram of Blazars, γ-ray burst sources and the initial 0.1 sec hard pulse of the March 5, 1979 burst]

**Fig. 1.** Blazars, γ-ray burst sources and the initial 0.1 sec hard pulse of the March 5, 1979 burst occupy distinct regions of the plane defined by source compactness $\ell_c$ and asymptotic Lorentz factor $\gamma_\infty$. GRB outflows may contain lower-$\gamma_\infty$ ejecta that manifests itself as softer tails, precursors and sub-pulses (T94, T96). The Lorentz factors of Blazar sources are only indirectly constrained by superluminal motion outside the γ-ray emitting region (cf. Wagner, these proceedings).

Another key point is that sources (or types of sources) exhibit different spectral states. GRBs are predominantly non-thermal, but some contain subluminous precursors, tails, and sub-pulses with distinctly thermal high energy cutoffs (Yoshida et al. 1989; Pendleton et al. 1996). Blazars are occasionally observed with emission peaked strongly at $\sim 1$ MeV, in distinction to the more usual extended power-law behavior (Bloemen et al. 1995; Blom et al. 1995). And the remarkable 5 March 1979 burst was initiated by an extremely intense $\sim 0.1$ sec flare whose luminosity exceeded that of the remainder of the burst by a factor $\sim 300$ and showed much more pronounced spectral softening (Fenimore et al. 1996).

Spectral states with sharp high energy cutoffs have a simple interpretation as the residue of an optically thick outflow. Indeed, at very high $\ell_c$, the outflow is self-shielding from the central radiation source, as well as from external (e.g. side-scattered) radiation. Radiation advected outward from large scattering depth then becomes an important source of seeds for extended non-thermal spectra – in addition to the more familiar optically thin synchrotron-self-Compton mechanism. This means that the inner boundary conditions on the flow are much more important than is usually supposed. And this raises an interesting question: are sources of relatively low compactness (e.g. Blazars) ever self-shielding in this manner?
It is intriguing to note, in this regard, that the characteristic energy $\sim 1$ MeV appears in a number of sources (of widely varying parameters): the spectral breaks observed in OSSE jet sources ($\sim 1 - 3$ MeV after compensating for cosmological redshift; e.g. McNaron-Brown et al. 1995) and in classical GRBs ($\sim 100$ keV $\sim 1$ MeV at peak luminosity, without compensating for redshift; Mallozzi et al. 1995); in MeV Blazars; as well as the initial hard spike of the March 5 burst ($\sim 300$ keV). Of course, the possibility that selection effects narrow the observed spectral break distribution should be considered carefully. This happens in the case of the GRB sources (Piran & Narayan 1996) only if the total burst energy is constrained to yield an inverse correlation between break energy and flux – in distinction to the strong positive correction observed within individual bursts. Most plausible cosmological GRB sources release as much as $10^{53} - 10^{54}$ erg, which allows for a fraction of very energetic and hard bursts.

After careful consideration of various photon sources, it turns out that an advected Wien peak maintains a characteristic energy of $\sim 1$ MeV over a wide range of central compactness, if the flow is sufficiently relativistic (Sects. 1.2; 3.1). Pair creation by photon collisions, $\gamma + \gamma \rightarrow e^+ + e^-$, which has traditionally been viewed as inimical to high energy gamma-ray production, can in fact play an essential role by i) increasing the efficiency of leptonic dissipative modes; ii) reducing the lower cut-off energy of the non-thermal pair distribution to $\gamma_{\text{min}} \sim 1$, which yields a break in the spectral distribution of Compton-upscattered Wien photons near the position of the original Wien peak; and iii) selecting non-thermal over thermal spectra at Comptonizing hotspots a high-$\gamma$ outflow (Sect. 3.3). Since the photon collision cross section is comparable to Thomson, feedback from pair creation works most effectively near $\tau_T \sim 1$. This contrasts with the inhomogeneous external Comptonization model (Blandford and Levinson 1995, hereafter BL95), where the non-thermal high energy continuum emerges well outside the scattering photosphere. Indeed, the position of the scattering photosphere in a pair-loaded outflow is sensitive to the amount of continuous heating, and pair creation can significantly broaden the transition zone between optically thick and thin flows (Sects. 1.2, 2.2).

Thus, by focussing on those aspects of the physics that are special to the large-\(\ell_c\) GRB regime, and then considering how these vary with compactness, interesting new insights can be obtained on both the GRB and Blazar problems. When constructing GRB models, it is sobering to realize that Blazars are still far from being understood, even with the much broader spectral information available.

In these notes, I will explore the following additional points:

- The dependence of the dominant dissipative mode on optical depth. At $\tau_T > 1$ direct Comptonization by bulk fluid motions is most effective; whereas non-thermal (e.g. Fermi) particle acceleration is a crucial ingredient of any radiative model at low $\tau_T$. Strong-wave acceleration can be excluded if enough scattering charges are present to generate an observable flux of Comptonized high energy photons.

- What is the relative importance of double Compton emission and cyclo-synchrotron emission as seeds for Comptonization (at large $\ell_c$)?

- How does the influence of geometrical effects (e.g. beaming) vary with $\gamma_\infty$?
High energy cutoffs to extended power-law spectra are extremely diagnostic. In GRB sources these are very poorly constrained. Measurements by EGRET in the 30 MeV - 10 GeV range indicate a significant decorrelation with the 1 MeV flux, with the high energy emission often being significantly delayed (Hurley et al. 1994).

1.2 The \( \ell_c-\gamma_\infty \) Plane

The variety of possible dissipative modes is neatly summarized in a two-dimensional plane labeled by (Fig. 2)

\[
\ell_c = \frac{L_{rel}\sigma_T}{4\pi m_e c^3 R_0^3}; \quad \gamma_\infty = \frac{L_{rel}}{\dot{\mathcal{M}} c^2}.
\]

The radius \( R_0 \) of the central engine is identified with the Alfvén radius or light-cylinder radius as appropriate. At risk of oversimplification I will usually assume that \( \gamma \) has attained the limiting value \( \gamma_\infty = L_{rel}/\dot{\mathcal{M}} c^2 \) due to matter loading \( \dot{\mathcal{M}} \) (when discussing delayed dissipation at large distances from the central engine).

The radius at which dissipation takes place is limited significantly by causality at large \( \gamma \), as long as \( \gamma(r) \) grows faster with radius than \( r^{1/2} \) near the base of the outflow. Thus, it is convenient to define the compactness \( \ell_{\Delta t} = \ell_c \times (R_0/c\Delta t) \) associated with variability on a timescale \( \Delta t \). A plausible value of the dissipative radius is \( \sim 2\gamma_\infty^2 c\Delta t \) for a variety of dissipative modes, including magnetic reconnection and MHD turbulence (Romanova \& Lovelace 1992, hereafter RL92; Thompson 1994, 1996, hereafter T94, T96; Levinson, these proceedings) and shocks powered by variations in \( \dot{\mathcal{M}} \) (Pacyzński \& Xu 1994, hereafter PX94; Rees and Mészáros 1994, hereafter RM94). The radius of the scattering photosphere of the wind is related to \( R_{\text{diss}} \) by a very strong function of \( \gamma_\infty \),

\[
\frac{R_{\tau=1}^{e-\rho}}{2\gamma_\infty^2 c\Delta t} \sim \frac{m_e}{m_p} \frac{\ell_{\Delta t}}{\gamma_\infty^5} \quad (n_p \gg n_e^+),
\]

when the scattering depth is dominated by the advected electron-ion contaminant. A non-thermal photon tail extending above energy \( m_e c^2 \) in the rest frame will greatly increase the density of scatterers, with the result

\[
\frac{R_{\tau=1}^{e+\pm}}{2\gamma_\infty^2 c\Delta t} \sim \frac{\ell_{\Delta t}}{\gamma_\infty^5 \gamma_\infty^5} \quad (n_e^+ \gg n_p)
\]

for a photon index \( \beta \sim -2 \) characteristic of GRBs. The effects of pairs are discussed further in Sect. 3.

Outflows with Lorentz factor \( \gamma_\infty < \gamma_{e-\rho} = (m_e/m_p)^{1/5} \ell_{\Delta t}^{1/5} \) dissipate well inside the electron-ion photosphere. They occupy the upper left portion of Fig. 1, and have a possible association with the soft X-ray precursors, tails and quasi-thermal sub-pulses of GRBs (T96). Outflows with \( \gamma_{e-\rho} < \gamma_\infty < \gamma_{e+} = \ell_{\Delta t}^{1/5} \) dissipate inside the the pair photosphere, if the high energy continuum extends up to \( \sim m_e c^2 \) in
Fig. 2. The variety of dissipative regimes in relativistic outflows as summarized in the \( \ell_c - \gamma_\infty \) plane. Flows with heavy matter loading and asymptotic Lorentz factor \( \gamma_\infty < \gamma_{e-p} \) occupy the top left region. They are a plausible source of soft subcomponents of GRBs. Flows with \( \gamma_{e-p} < \gamma_\infty < \gamma_{e\pm} \) can become pair loaded, and occupy the adjoining stripe. Flows with low matter loading, \( \gamma_\infty > \gamma_{e\pm} \), dissipate at low optical depth. Above the upper horizontal dashed line, local thermodynamic equilibrium is achieved at the base of the flow; whereas below this line the photon distribution is Wien. All of the flows in this region are assumed to undergo delayed dissipation after thermal pairs freeze out, at a radius where inhomogeneities (such as reconnecting magnetic fields and shocks) regain causal contact. By contrast, flows in the lower portion of the diagram are assumed to dissipate continuously and remain pair loaded out to the scattering photosphere. This is expected in collimated flows with lower asymptotic Lorentz factors (such as Blazars). Large scattering depths cannot be maintained below the lower horizontal dashed line.

the rest frame. And, outflows with \( \gamma_\infty > \gamma_{e\pm} \) dissipate at low scattering depth, independent of the efficiency of pair creation. The spectral consequences of variations in the matter loading are discussed further in Sect. 3.

Flows in which the irregularities maintain causal contact will undergo continuous heating while optically thick. An interesting example of such a flow is the low-\( \gamma \) sheath of a relativistic jet, in which the luminosity of entrained photons increases with radius as the kinetic energy of the higher-\( \gamma \) core of the jet is dissipated. Although the relation between scattering depth and internal temperature of the flow is sensitive to the high energy distributions of the pairs and photons, the mean energy of the emergent photon spectrum is regulated to near \( \sim \gamma m_e c^2 \). In the case where the photon distribution is Wien, the strong \( T \)-dependence \( n_{e\pm}/n_\gamma = (\pi/2)^{1/2}(T/m_e c^2)^{-3/2} \exp(-m_e c^2/T) \) of the equilibrium pair density guarantees that as soon as \( T \) drops much below \( m_e c^2 \) in the rest frame, the flow becomes optically thin. This yields a simple relation between the observed
(Lorentz-boosted) temperature and the temperature $T_0$ at the base of the flow,

$$T_{\text{obs}} = T_0 \left( \frac{L_\gamma}{L_\gamma 0} \right),$$

in terms of the growth of the photon luminosity from $L_\gamma 0$ to $L_\gamma$.

The pair density resulting from direct heating of the photons by MHD turbulence can be less sensitive to $T$. I assume that the bulk of the photon energy lies in a Wien bump, but that wave energy is excited in transient surges (e.g. by reconnection) with an equivalent temperature $T_w$ somewhat in excess of $m_e c^2$. Then a significant fraction of the wave energy is converted to photons above the pair production threshold. Balancing this source against pair annihilation, the scattering depth in the direction perpendicular to a jet (of opening angle $\theta$) is

$$\tau_{T \perp} = \frac{1}{2} n_e \sigma R \theta \sim (\gamma \ell_\gamma)^{1/2},$$

where $\ell_\gamma$ is the compactness (1) in the advected photons at radius $R$. At $\tau_T > 1 (\ell_\gamma > 1)$, freshly created pairs Compton cool before annihilating (the cooling timescale being shorter by a factor $\sim \tau_{T \parallel}^{-1}$) and carry a fraction $\sim \tau_{T \parallel}^{-1}$ of the total energy of the flow. This regulates the Compton parameter induced by mildly relativistic pairs to $y \sim \tau_{T \parallel}^{-1} \cdot \tau_T \sim 1$. The annihilation photons Compton downscatter off the cooled pairs to an energy $\sim m_e c^2 / \tau_T$ in the rest frame. The photon spectrum emerging at the pair photosphere then peaks at an energy $\sim \gamma m_e c^2$. This implies only a modest increase in the mean energy per photon along the jet, by a factor $\sim \gamma m_e c^2 / 3 T_0$, which is easily supplied from the high-\(\gamma\) core to the low-\(\gamma\) sheath.

Continuous heating by relativistic electrons at low optical depth has been considered by Sikora et al. (1997) as a model for the MeV Blazars. They relate the peak energy to the observed variability timescales, but since this energy is not directly tied to a microphysical scale, it could be expected to lie well below $\sim 1$ MeV in some sources. By contrast, direct Compton damping of mildly relativistic turbulence in an optically thick jet yields a Wien peak energy that is bounded below by $\sim 1$ MeV (although still dependent on bulk $\gamma$ and the energy transferred from the bulk motion to the photons).

### 2 Dissipative Mechanisms

#### 2.1 Sources of Free Energy

The internal sources of free energy in a relativistic flow can be broadly divided into two categories: those associated with radial and angular inhomogeneities. Both appear to be important in jet sources, and while both have also been considered in GRB sources, radial inhomogeneities are probably more important in the large-$\gamma$ context. Reconnection surfaces and variations in the ratio of particle pressure to magnetic pressure will lead to internal heating of strongly-magnetized outflows (RL92; T94), as will kinetic energy fluctuations in particle-dominated outflows (PX94; RM94). However, interactions with an external medium (Rees and Mészáros 1992) can become significantly non-spherical as the outflow in a $\gamma$-ray burst source decelerates (Sect. 4).
Which of these energy sources dominates depends on the strength of the magnetic field in the outflow and the radius of the dissipative zone. It appears that $\gamma_\infty \sim 100 - 300$ can be achieved in an outflow of luminosity $\sim 10^{51}$ erg s$^{-1}$ only if it is Poynting-flux dominated at the source. Indeed, any triggering scenario for a GRB that produces an object with the density of nuclear matter and a rotation period of $\sim 10^{-3}$ s plausibly involves magnetic fields as strong as $\sim 10^{15}$ G through dynamo amplification and thus leads to a rotationally-driven MHD outflow of luminosity $L_P \sim 10^{50} - 10^{51}$ erg s$^{-1}$ (T94; see also Duncan & Thompson 1992; Usov 1992, 1994; Vietri 1996; Mészáros and Rees 1997 for particular models). The alternative mechanism of $\nu - \bar{\nu}$ annihilation into $e^\pm$ pairs has an efficiency of $\sim 10^{-3}$ (Jaroszynski 1993), and so the neutrino luminosity required to power a $\gamma$-ray flux of $\sim 10^{51}$ erg s$^{-1}$ drives a mass flux (by absorption on nucleons) that exceeds the tolerable value by $\sim 10^6$ (cf. Duncan, Shapiro and Wasserman 1986).

This advected magnetic field almost certainly has an important effect on the dissipative mechanism. Quasi-perpendicular shocks are significantly weakened even when the magnetic field carries a few percent of the energy flux; this in turn steepens non-thermal particle spectra arising from first-order Fermi acceleration. Although particle acceleration at oblique relativistic shocks can be quite efficient (as discussed by Kirk, these proceedings), one expects that internal shocks in GRB outflows with $\gamma \sim 100 - 300$ are approximately radial. Thus, GRB models based on internal shocks may unfortunately require complicated departures from spherical symmetry.

### 2.2 Geometrical Effects: Pair Cocoons

The huge compactness of a GRB source ($\sim 10^{15}$) causes it to be self-shielding. Not only is the central engine hidden from the dissipative zone, but the optical depth $\tau_\perp$ perpendicular to the axis of the flow exceeds the parallel depth by the very large factor,

$$\frac{\tau_\perp}{\tau_\parallel} \sim \gamma^2 \theta \sim 10^{4-5} \theta,$$

for reasonable values of the opening angle $\theta$.

This raises the question: as $\ell_c$ decreases toward the values typical of AGN, at what point does the central engine become visible? One intriguing possibility is that the engine remains shielded in some Blazars, with the high-$\gamma$ cores of the jet being surrounded by a pair cocoon (Fig 3). Because Lorentz dilation causes $\tau_\parallel$ to decrease in proportion to $\gamma^{-2}(L_{jet}/\dot{M}_{jet}c^2)^{-1} \propto \gamma^{-3}$, the high-$\gamma$ core of a jet can become optically thin along its axis well inside the photosphere of the $\gamma \sim 1 - 2$ sheath. This lies at a radius

$$R_{\tau=1}^\pm \sim 2 \times 10^{18} \left(\frac{L_{jet}}{10^{47} \text{ erg s}^{-1}}\right) \left(\frac{\theta}{0.1}\right)^{-1} \text{ cm},$$

assuming that the sheath acquires a non-negligible fraction of the kinetic energy at this radius. All that is required (Sect. 1) is that the sheath be i) pair-loaded and optically thick at its base, and ii) continuously heated (e.g. by Kelvin-Helmholtz
instabilities with higher-\(\gamma\) material) at a sufficient rate to overcome the effects of adiabatic cooling. Condition ii) is plausibly satisfied in the cores of superluminal sources, and i) is satisfied if the outflow is strongly turbulent at its base (Sect. 3). Below this threshold heating rate, the position of the pair photosphere is very sensitive to the amount of heating.

Fig. 3. A sheath of low-\(\gamma\) material surrounding the high-\(\gamma\) core of a jet can remain pair loaded out to a large distance (6) from the central engine. The scattering photosphere of this pair cocoon will in general lie far outside the radius at which the core becomes transparent to high energy \(\gamma\)-rays.

2.3 Diverse Photon Sources

When the flow is self-shielding in this manner, advected quasi-thermal radiation becomes an important source of Compton seeds (T94, T96). Even though external radiation exerts a stronger drag force than internally generated radiation by a factor \(\gamma^2\) (Sikora, Begelman, & Rees 1994; BL95), it cannot penetrate the high-\(\gamma\) component of the flow. This also disfavors models involving acceleration and heating of the jet by central continuum radiation propagating along the jet axis (Dermer and Schlickeiser 1993; Marcowith et al. 1995). The main competing source of seed photons is then synchrotron radiation (e.g. Maraschi et al. 1992; Mészáros, Rees and Papathanassiou 1994, hereafter MRP94).

In quantifying the relative importance of these two radiation sources, one must consider separately thermal and non-thermal scatterers. The advected radiation characteristically has a much higher frequency than the cyclotron frequency, and if its energy density \(U_\gamma^a\) is comparable to \(B^2/8\pi\), then it will dominate the magnetic field as a coolant of thermal electrons, due to self-absorption near the cyclotron resonance. Parametrizing the frequency at which the cyclo-synchrotron radiation
becomes self-absorbed in terms of a cyclotron harmonic \( N_{sa} \), one deduces from Kirchoff’s law that C-S cooling rate is smaller than the Compton cooling rate off the advected radiation by the factor

\[
\frac{\nu_j \nu}{4\sigma_T n_e c(T/m_ec^2)U_{\gamma}} = \frac{2\alpha_{em} N_{sa}^3}{\pi} \left( \frac{B}{B_{QED}} \right) \frac{B^2/8\pi}{U_{\gamma}}, \quad \left( \nu = N_{sa} \frac{eB}{2\pi m_ec} \right) \tag{7}
\]

where \( \alpha_{em} = 1/137 \) and \( B_{QED} = 4.4 \times 10^{13} \) G. For example, one expects \( B/B_{QED} < 10^{-8} \) near the scattering photospheres of GRB outflows, so that C-S cooling can be neglected even if \( N_{sa} \sim 300 \).

Another effect of the advected radiation is to limit the Compton parameter

\[
dy/d\ln R \simeq 4(T/m_ec^2)(1 + 4T/m_ec^2)d\tau_T/d\ln R \text{ in optically thick regions of the flow.}
\]

The magnitude of this effect depends on the mean energy per photon \( \langle h\nu \rangle \), and the bulk Lorentz factor \( \gamma \). If the flow is photon rich, with \( \langle h\nu \rangle/\gamma \ll m_ec^2 \) in its rest frame, then the advected photons limit \( y \) to a value \( \sim \ln(\langle \nu \rangle/\nu_0) \), where \( \nu_0 \) is peak frequency before the advected photons undergo (delayed) reheating. This prevents very low energy C-S photons from being upscattered much in frequency. A stronger limit \( y \sim 1 \) applies in outflows that are continuously Compton heated by MHD waves and shocks (Sect. 1.2). Nonetheless, a much larger \( y \)-parameter will be maintained at the base of the outflow, where the photon flux has a net divergence (Sect. 3.1).

Synchrotron radiation from Blazar sources is usually ascribed to the same high energy electrons/positrons that are responsible for the X-ray/\( \gamma \)-ray emission. The high energy cutoff of the synchrotron peak has been ascribed to a suppression of the non-thermal particle density inside the scattering photosphere (e.g. Levinson 1996), but mildly relativistic pairs are only suppressed by a factor \( \sim \tau_T^{-1} \). Thus, the particle spectrum must itself steepen considerably at \( \tau_T > 1 \). Rather large minimum non-thermal Lorentz factors \( \gamma_{min} \) are sometimes conjectured in GRB sources, sufficient to place the synchrotron peak energy directly in the MeV range (e.g. MRP94). However, synchrotron absorption becomes much more important in that context, if the outflow becomes sufficiently pair loaded that \( \gamma_{min} \sim 1 \) (Sect. 3.3).

Advected radiation can also be the dominant coolant at external relativistic shocks. Consider a shell of relativistic matter of initial width \( c\Delta t \). After the shell becomes optically thin to scattering, the radiation moves ahead of the shell, but continues to overlap inside a radius \( R > 2\gamma^2 c\Delta t \). The point is that when the external medium is cold (with a sound or Alfven speed much less than \( c \)), the external magnetic field \( B_{ex} \) that is swept up and compressed in between the forward shock and the contact discontinuity typically has a much smaller energy density than the radiation, by a factor

\[
\frac{B^2/8\pi}{U_{\gamma}} \sim 2 \times 10^{-10} \left( \frac{B_{ex}}{3 \times 10^{-6} \text{ G}} \right)^2 \left( \frac{L_{\gamma}}{10^{50} \text{ erg s}^{-1}} \right)^{-1} \left( \frac{\gamma}{10^2} \right)^8 \left( \frac{\Delta t}{10 \text{ s}} \right)^2 \tag{8}
\]

at \( R = 2\gamma^2 c\Delta t \) (assuming a compression factor of 7). This strongly suggests that direct synchrotron radiation from the forward shock is not the mechanism primarily
responsible for delayed optical and X-ray emission from GRBs. Direct radiation from a low-$\gamma$ component of the ejecta is considered in Sect. 4.

2.4 Dissipation at $\tau_T > 1$

MHD waves, turbulent motions and shocks can in principle tap a significant fraction of the energy of the outflow. A periodic excitation of frequency $\omega_k$ involving a displacement $\xi_k$ of the fluid deposits its energy directly in the photons via Compton drag when the scattering depth across $\xi_k$ is less than $\sim c/\omega_k \xi$ (T94). For example, relativistic Alfvén waves have a damping time

$$t_{\text{drag}} \omega_k \sim \frac{(\delta B)^2_k/8\pi}{U_\gamma},$$

(9)

which is shorter than the wave period as long as the photon gas has a higher energy density than waves in the relevant range of wavenumbers (Thompson and Blaes 1997, hereafter TB97). This mechanism is particularly effective near the scattering photosphere, and is also effective even at large $\tau_T$, as long as higher wavenumber turbulence is generated via a turbulent cascade. The wave amplitude generally decreases with wavenumber in such a cascade, with the result that $(\delta B)_k^2/8\pi \ll U_\gamma$ at dissipative scales even if $(\delta B)_k^2/8\pi \sim U_\gamma$ at the outer scale. Shocks also transfer energy to the photon fluid via compression and direct first-order Fermi acceleration (Blandford and Payne 1981), although energy transfer is slowed significantly when the photon pressure becomes comparable to the material ram pressure ahead of the shock.

2.5 Dissipation at $\tau_T < 1$

The high energy emission in Blazars covers a wide range of energies, up to 10 GeV (or higher in the TeV sources) and hence must be powered by non-thermal particle distributions. The most familiar possibilities are first-order Fermi acceleration at shocks (Blandford and Eichler 1987) and electrostatic acceleration (e.g. RL92). Photo-pion production on protons can more easily accomodate the TeV sources (Mannheim 1993), but requires a supplementary $\sim$ MeV emission mechanism in sources with soft high energy spectral states, because of the much greater cross-section for $\gamma+\gamma \rightarrow e^\pm$.

The relevant physics is treated in sufficient depth elsewhere that I will focus on two questions here.

1. Does reconnection deposit energy primarily in thermal or non-thermal particles, and in electrons or ions? This problem is far from being understood from first principles, but Solar flares do provide clear evidence that the efficiency of electron acceleration can (at least in a non-relativistic plasma) be quite high. However, Type III radio bursts (which are powered by flare particles that escape the Sun along open magnetic field lines) also provide direct evidence that most of the flare energy is dissipated in the form of bulk heating of electrons to energies of 10-100
keV, with only a small fraction being deposited in a relativistic, non-thermal tail (Lin 1990).

2. What is the effectiveness of electrostatic and strong-wave acceleration in a relativistic, Comptonizing medium, as compared to shock acceleration? This question highlights a crucial difference between the large scale relativistic outflows associated with Blazars and GRBs, and a laboratory system such as a Tokomak (or even smaller scale astrophysical systems such as coronal loops and arcades). A magnetic field $B$ with gradient scale $\ell_B = B/|\nabla B|$ requires a minimal charge density $n_{c,\text{min}} = B/4\pi e \ell_B$ to support the associated current; otherwise the displacement current cannot be neglected and charges are accelerated to relativistic energies. The ratio of the actual electron density to $n_{c,\text{min}}$ can be expressed in terms of the scattering depth across $\ell_B$,

$$\frac{n_e}{n_{c,\text{min}}} \sim \frac{\tau_T}{\alpha_{\text{em}}} \left(\frac{B}{B_{\text{QED}}}\right)^{-1}. \quad (10)$$

This works out to $n_e/n_{c,\text{min}} \sim 10^{13}$ for parameters appropriate to inhomogeneous external Compton Blazar models (BL95). It has been hypothesized that the boundary layers of jets may be partially evacuated and sites for strong-wave acceleration (Bisnovatyi-Kogan & Lovelace 1997), but in fact the degree of evacuation must be extraordinary for such a mechanism to be important.

Another approach to this problem is to re-express $B$ in terms of the plasma $\beta_e = 8\pi n_e T/B^2$,

$$\frac{n_e}{n_{c,\text{min}}} \sim \beta_e \left(\frac{B}{B_{\text{QED}}}\right) \left(\frac{T}{m_e c^2}\right)^{-1} \frac{\ell_B}{\ell_{m_e}}. \quad (11)$$

where $\ell_{m_e} = 4 \times 10^{-11}$ cm is the Compton wavelength of the electron. This is $n_e/n_{c,\text{min}} \sim 10^4 \beta_e$, $10^8 \beta_e$, and $10^{21} \beta_e$ for parameters appropriate to Tokomaks, Solar flares, and Blazar jets. This suggests that much higher wavenumber distortions of the magnetic field are required to provide efficient electron acceleration through reconnection in relativistic outflows. By contrast, direct Comptonization of a background photon fluid is more effective in jets and GRBs due to the higher scattering depth (T94).

3 Spectral Consequences

Most modelling of high energy emission from relativistic outflows ignores the inner boundary condition on the outflow. The flow is hypothesized to dissipate outside the pair annihilation radius, and flow conditions interior to that radius are assumed not to influence the emergent high energy spectrum. Seed radiation for Comptonization is assumed to originate in a central accretion disk exterior to the volume of the outflow.

We have already seen, however, that the emergent spectrum can be dominated by advected radiation if the outflow is optically thick at its base (T94, T96).
Moreover, the inner boundary conditions are very well defined in flows of large central compactness, such as GRBs ($\ell_c \sim 10^{15}$). The mean photon energy emerging from the flow is $\langle h\nu \rangle \sim L_P/\dot{N}_\gamma$, when the asymptotic Lorentz factor lies near the critical value $\gamma_{\pm}$ (or $\gamma_{-\gamma}$ if pairs are absent). Thermalization is rapid near the base of the flow, the photon gas is very close to black body, and $\langle h\nu \rangle$ is directly related to the effective temperature at the light cylinder,

$$\langle h\nu \rangle \sim T_{eff} = 0.8 \left( \frac{L_\gamma}{10^{50}} \right)^{1/4} \left( \frac{P}{10^{-3} \text{ s}} \right)^{-1/2} \text{MeV}. \quad (12)$$

This is remarkably close to the observed range of spectral break energies, after allowing for cosmological redshift.

In this regime, the ratio of photon luminosity $L_\gamma$ to (ordered) Poynting luminosity $L_P$ at the base of the wind is a key parameter. It should be emphasized that the baryon loading is tolerably small only if $L_\gamma < 10^{-2}L_P$ at the neutrinosphere. In other words, a key requirement of this model is that the wind be reheated from $L_\gamma \ll L_P$ to $L_\gamma \sim L_P$ well outside the neutrinosphere. This is plausibly accomplished by a MHD cascade to high wavenumber, even though photons and pairs are tightly coupled on macroscopic scales (Sect. 2.4; TB97).

Expression (12) also leads to an interesting question: how does the mean energy per photon change as one decreases the central compactness? As we now show, the trend of decreasing mean photon energy with decreasing compactness in fact can be reversed, with $\langle h\nu \rangle$ approaching $\sim m_e c^2$ at $\ell_c \sim 10^2$.

### 3.1 Direct MeV Wien Peak

A magnetized outflow that is strongly turbulent will trigger a cascade to high wavenumber that *must* dissipate inside the Alfvén radius. We look for an optically thick equilibrium state in which the dissipative zone is shielded from external photons and soft photons are generated internally. In applications to AGN the advected matter contributes negligible optical depth, and so the outflow is necessarily hot and pair loaded. Damping via resonant couplings between high wavenumber turbulence and cosmic ray particles has been considered by Dermer, Miller, & Li (1996). However, direct Compton drag of bulk turbulent motions is an effective damping mechanism at high compactness, typically at much lower wavenumbers (TB97). I now estimate the temperature $T_0$ of the flow at its base, focusing on two photon sources.

1. **Double Compton Emission.** This dominates bremsstrahlung emission when $T_0 \sim m_e c^2$ and the outflow is photon rich, $n_{\gamma} \gg n_e = n_{e^+} + n_{e^-}$ In a Wien photon gas,

$$\dot{n}_dC = \frac{16A}{\pi} \alpha_{em} n_e n_{\gamma} \sigma_{TC} \left( \frac{T_0}{m_e c^2} \right)^2,$$

where $A \simeq \ln(T/h\nu_{min})$ and the photon gas approaches a Planckian distribution at frequency $\nu_{min}$.

Double Compton dominates cyclotron emission when the magnetic field in the central engine exceeds $B_{QED} = 4.4 \times 10^{13}$ G, so that thermal pairs do not populate
excited Landau levels. Such strong fields have indeed been associated with SGR 0526-66, which emitted the March 5, 1979 superburst. The initial 0.1 s pulse of that burst appears to have approached a luminosity of $\sim 10^7$ times the Eddington luminosity (Fenimore et al. 1996), and has the appearance of an expanding pair fireball (Thompson & Duncan 1995; Fatuzzo & Melia 1996). If the outflow is driven by a magnetic field that is also strongly turbulent, then Compton drag can raise the photon energy density $3T n_0$ close to $B^2/8\pi$ near the light cylinder (TB97).

Equating $\dot{n}_d C$ with the photon loss rate, one deduces an equilibrium scattering depth

$$\tau_T = \frac{\pi}{16 A \alpha_{em}} \left( \frac{T}{m_e c^2} \right)^{-2}. \quad (14)$$

Re-expressing this in terms of the equilibrium pair density yields the relation between $T_0$ and compactness $\ell_c$ shown in Fig. 4.

![Fig. 4. Mean photon energy $\langle h\nu \rangle$ versus central compactness $\ell_c$ for flows in which double Compton emission (solid line) and cyclo-synchrotron emission (long-dashed line) is the dominant soft photon source. The minimum temperature for an optical thick flow (assuming a Wien photon distribution) is labelled by the short-dashed line. Note the breaks in the curves at the transition from Wien to black-body photon distributions.](image)

2. Cyclo-Synchrotron Emission in a Non-thermal Pair Plasma. Such a very high optical depth and compactness cannot be maintained in weaker magnetic fields. Cyclo-synchrotron photons are created rapidly, and their energy rapidly exponentiates. (For related calculations with non-thermal particle distributions, see Ghisellini, Guilbert, & Svensson 1988.) To show this, I parametrize by $N_e eB/m_e c$ the critical frequency at which Compton scattering increases the frequency of a C-S photon at the same rate as it is absorbed, $dy/dt = c \alpha_\nu$. From Kirchoff’s law,
\[
\frac{\dot{n}_{C-S}}{\dot{n}_{dC}} \sim \frac{2N_c^2}{A} \left( \frac{B^2 / 8\pi}{n_r m_e c^2} \right). \tag{15}
\]

Since \(N_c \gg 1\) for \(3T \sim m_e c^2\) (Mahadevan, Narayan, & Yi 1996) this yields \(\dot{n}_{C-S} \gg \dot{n}_{dC}\).

Is the energy released by a turbulent cascade at the base of the outflow dissipated at large or small optical depth? If even a tiny fraction of the bulk kinetic energy is carried by particles, then the wave energy must cascade to very high wavenumber before charges are accelerated either electrostatically – as the turbulence becomes charge starved – or by resonant interactions. The relative importance of these two effects is determined by the ratio

\[
\frac{k_z}{eB/m_e c^2} \sim \left( \frac{\ell}{\ell_m} \right)^{-1} \frac{\tau_T}{\alpha_{em}} \left( \frac{B}{B_{QED}} \right)^{-2},
\]

in the case of sheared Alfvén waves of wavenumber \(k_z\) (TB97). The cascade can, however, be cut off by Compton drag much closer to the outer scale. The MHD wave motions are only mildly relativistic, and so a large scattering depth is required to provide the Compton parameter \(\gamma \sim 10\) needed to upscatter the bulk of the C-S photons to high energy. The mean energy per photon is given by

\[
\frac{3T_0}{m_e c^2} \sim \frac{B^2 / 8\pi}{m_e c^2 \dot{n}_{C-S}(R/c)} \sim \frac{\pi}{32N_c^2 \alpha_{em}} \tau_T^{-1} \left( \frac{T_0}{m_e c^2} \right)^{-2}, \tag{16}
\]

where I approximate the bulk of the photon distribution as Wien. Assuming that thermal and turbulent velocities are comparable, one has \(T_0/m_e c^2 \approx y/8\tau_T\) and

\[
\tau_T \sim 13 \left( \frac{y}{10} \right)^{3/2} \left( \frac{N_c}{20} \right). \tag{17}
\]

This gives \(T_0/m_e c^2 = 0.095 (y/10)^{-1/2} (N_c/20)^{-1}\), as shown in Fig. 4.

This scattering depth lies above the thermal value, and so must be maintained by an extended non-thermal tail to the pair distribution function. This is naturally provided by MHD wave heating, as discussed in Sect. 1.2. The minimum central compactness needed to support this optically thick state is, on energetic grounds, \(\ell_c \sim 10\) (Fig. 2).

### 3.2 Dissipation at Large Compactness and Low \(\gamma_\infty\)

A much wider range of peak energies is possible if enough matter is advected by the outflow to make it optically thick. This is possible only for a rather higher compactness, \(\ell_c > m_p/m_e \sim 2000\), than is expected in Blazar sources, but well within the range of cosmological GRB sources (cf. Paczyński 1990). If the matter loading is heavy enough that \(\gamma_\infty < \gamma_{e-p} = 0.2\ell_5^{1/5}\), then dissipation on an observed timescale \(\Delta t\) occurs inside the \(e-p\) scattering photosphere. This provides a natural explanation for the soft tails and precursors to GRBs seen by Ginga, and the soft sub-pulses seen by BATSE (T94, T96).
Unlike pair-dominated outflows with negligible matter (Sect. 3.1) the emergent temperature is very sensitive to the amount of continuous heating. (In the case of spherical flows with radial inhomogeneities, this is equivalent to a broad power spectrum of inhomogeneities.) I further assume that the flow expands relativistically at its base, with $\gamma(r) \propto r^\alpha$ with $\alpha > \frac{1}{2}$. Then inhomogeneities in the flow fall out of causal contact, and dissipation on timescale $\Delta t$ is delayed to radius $R_{\text{dis}} \approx 2\gamma^2 c \Delta t$. The corresponding luminosity $L(R_{\text{dis}})$ is then depleted by adiabatic expansion out to the scattering photosphere. Assuming that photon-number changing processes freeze out before dissipation takes place, the emergent temperature is a very strong function of matter loading,

$$\frac{T_{\text{obs}}}{T_0} = 0.7 \frac{L(R_{\text{dis}})}{L(\gamma(R_0))} \left( \frac{\gamma}{\gamma_{\text{e-p}}} \right)^{10/3} \propto (\Delta t)^{2/3}. \tag{18}$$

The correction for adiabatic losses (the last factor in this expression) has a very strong dependence on $\gamma$, but a much weaker dependence on the (more directly observable) variability timescale $\Delta t$. This means that the power spectrum of inhomogeneities must cover a very wide range of frequencies for continuous heating to overcome the effects of adiabatic cooling in a radial flow. The situation is quite different in collimated flows, where energy can be continuously extracted from angular (e.g. Kelvin-Helmholtz) instabilities (Sect. 2.2).

The number of advected photons must also be compared with the number of fresh cyclo-synchrotron photons generated in the flow outside the central engine. Following Sect 3.1, this is

$$\frac{\dot{n}_{C-S} t}{n_\gamma} \sim \frac{8 \alpha_{em} N^2 a}{\pi \gamma} \frac{d\gamma}{d\ln t} \left( \frac{T}{m_e c^2} \right). \tag{19}$$

Here all quantities refer to the rest frame of the outflow. The advected photons suppress $\dot{n}_{C-S}$ at large $\tau_T$ by soaking up energy from the electrons and holding down $\gamma$. We conclude that $\dot{n}_{C-S}/n_\gamma$ is typically less than unity for high-$\gamma$ outflows.

### 3.3 $e^+ - e^-$ Amplifier in Relativistic Outflows

The efficiency with which the available energy (in shocks, MHD waves, and reconnecting magnetic fields) is deposited in high energy photons depends directly on the fraction $\varepsilon_e$ of the energy deposited in electrons and pairs. Pair creation (via photon collisions $\gamma + \gamma \rightarrow e^+ + e^-$) has traditionally been used to constrain the emission region in high energy sources (e.g. Cavallo & Rees 1978; Baring & Harding 1997), but one would like to emphasize an opposing point of view here: that the high energy photon flux from a relativistic outflow can be significantly

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1. In contrast to the model of MeV Blazars discussed in Sects. 2.2 and 3.1, in which the collimated, lower-$\gamma$ flow is assumed to be heated continuously by non-radial Kelvin-Helmholtz instabilities.
raised by pair creation, due to an increase in $\varepsilon_e$. Energization of positrons by high-harmonic proton synchrotron maser radiation behind a shock (Hoshino & Arons 1991) provides an example of such a leptonic acceleration mechanism.

Pair creation has, of course, been included for some time in models of BH accretion disk coronae (e.g. Stern et al. 1996 and references therein), but in that context the formation of a power law high energy continuum does not depend essentially on the presence of pairs. For example, direct Comptonization by MHD motions in the corona will effectively heat the photons, independently of the relative amounts of rest energy in baryons and leptons (T94). The situation is quite different in a relativistic outflow that expands sufficiently rapidly at its base that inhomogeneities fall out of causal contact. If these inhomogeneities cover a wide range of spatial frequencies $k_{\text{min}} < k < k_{\text{max}}$ (as is needed in GRB models to accomodate the broad power spectra of the bursts) then pair creation at wave number $k_{\text{max}}$ (radius $\sim 4\pi\gamma^{-2}k_{\text{max}}^{-1}$) will increase the radius of the scattering photosphere by a factor

$$\frac{R_{\tau=1}^e}{R_{\tau=1}^p} \sim \left(\frac{m_p}{m_e}\right) \left(\frac{E_{\text{br}}}{\gamma m_e c^2}\right)^{2-\beta},$$

(20)

given a high energy photon index $\beta$. This pair amplifier allows a much wider range of wavenumbers to be dissipated directly by Compton drag, and hence causes regions of the wind with non-thermal high energy spectra to be significantly brighter than regions with thermal spectra (T96).

At this point I should distinguish between pair amplification that is linear (the probability of energization $P_e$ of a newly created pair is the same as that of a seed electron) and non-linear (the pairs feed back on $P_e$). For example, the amplifier operating at a shock is non-linear if fresh pairs all act as suprathermal seeds for first-order Fermi acceleration; whereas it is linear if the pairs cool down to the temperature of the background thermal plasma before interacting resonantly with plasma waves.\(^2\) The corresponding leptonic efficiencies are

$$\varepsilon_e = \frac{P_e + 2n_{e^+/n_p}}{P_p + P_e + 2n_{e^+/n_p}} \quad \text{(non-linear)};$$

$$= \frac{P_e(1 + 2n_{e^+/n_p})}{P_p + P_e(1 + n_{e^+/n_p})} \quad \text{(linear)},$$

(21)

where $P_p$ is the probability of energization of a proton and $\varepsilon_e = P_e/(P_e + P_p)$ in the absence of pairs.

The pair amplifier operating at Comptonizing hotspots (Sect. 3.4) is non-linear in different sense: pair creation regulates the high energy index $\beta$ to the appropriate value to yield a Thomson depth $\tau_T \sim \frac{1}{3}(T/m_e c^2)^{-1}$ within individual hotspots. This second-order Fermi acceleration mechanism therefore yields a power-law high energy spectrum over a wider range of matter loadings ($\gamma_{e-p} < \gamma_\infty < \gamma_{e^\pm} = 4.5\gamma_{e-p}$) than does synchrotron cooling of shock-accelerated pairs.

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\(^2\) The gyroperiod of a relativistic electron is orders of magnitude shorter than its cooling time if the magnetic field contributes an appreciable fraction of the pressure of the outflow.
Pair amplification via photon collisions is also non-local: photons upscattered above energy $m_e c^2$ in one hotspot will raise the pair density in another portion of the flow. A nice example is provided by an expanding shell of matter and photons of radial width $c \Delta t$. The bulk Lorentz factor of the photon gas is not limited by the inertia of the matter when $\gamma > \gamma_e \pm$. Outside a radius $\sim 2\gamma^2 c \Delta t$ photons with energies greater than $m_e c^2$ in the wind frame stream ahead of the forward shock, and sidescatter against seed electrons. The pair density grows exponentially (at first) inside a radius $\sim \ell c R_0$, until $n_e^+ / n_p$ reaches unity and the material ahead of the shock is accelerated to a limiting Lorentz factor $\sim \ell_1^1 / 2^\gamma \gamma_\infty$ (for a hard incident photon spectrum with $\beta = -2$). The feedback of pair creation on the structure of a relativistic shock is an interesting problem that has not been properly addressed.

One immediate spectral consequence of the pair amplifier is a suppression of the minimum leptonic Lorentz factor $\gamma_{\text{min}}$, and hence a suppression of the minimum synchrotron frequency $E_{\text{sync}}(\text{min}) \sim \gamma_{\text{min}}^2 eB / m_e c$. In fact, $\gamma_{\text{min}} \rightarrow 1$ as the inertia of the pairs becomes comparable to that of the protons. To give an example of the potential importance of this effect, consider a Poynting-flux dominated outflow that approaches its limiting Lorentz factor $\gamma_\infty$. The high energy photon index is taken to be $\beta = -2$ out to a rest frame energy $m_e c^2$. Near the scattering photosphere of the wind, the cyclotron energy is

$$\hbar \frac{eB}{m_e c} = \frac{8\pi e c^2}{\sigma_T(2L_P c)^{1/2}} \gamma_\infty^3 = 0.04 \left(\frac{\gamma_\infty}{300}\right)^3 \left(\frac{L_P}{10^{51} \text{ erg s}^{-1}}\right)^{-1/2} \text{ eV}. \quad (22)$$

The efficiency of electron acceleration can be increased by pair creation, but at the cost of suppressing $E_{\text{sync}}(\text{min})$ far below the observed range of break energies in GRB spectra. As a result, the primary emission process must be inverse Compton.

### 3.4 Delayed Inhomogeneous Comptonization: Broken Power-law Spectra with a Thermal Photon Source

Let us now consider the photon spectrum that results from delayed reheating of a relativistic outflow at large $\ell_c$, outside the electron-ion photosphere ($\gamma_\infty \sim \gamma_{e-p}$). At large $\ell_c$, the photons are adiabatically cooled in between the central engine and the causal contact radius, where they are reheated to a luminosity $L_\gamma \sim (\delta B / B)^2 L_P$ (when the outflow is Poynting-flux dominated). The mean photon energy is restored to a value

$$\langle \hbar \nu \rangle \sim 0.7 L_\gamma / L_\gamma 0, \quad (23)$$

given that photon number is conserved at this radius (Sect. 3.2).

As before, we consider a broad power spectrum of inhomogeneities, $k_{\text{min}} < k < k_{\text{max}}$; the corresponding (radial) size of a hot spot is $\Delta \sim \pi / k$. Let us suppose that

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\(^*\) The feedback of pair creation on the formation MeV breaks in high energy synchro-Compton cascades above accretion disks has been considered by Done, Ghisellini and Fabian (1990).
wavenumbers $k_{\ast} < k < k_{\text{max}}$ dissipate inside the electron-ion photosphere, and $k_{\text{min}} < k < k_{\ast}$ outside. Individual hotspots are assumed to release their energy when the causal propagation distance $R/2\gamma_\infty^2$ begins to exceed $\epsilon^{-1} \cdot \Delta$.

Hotspots with $k > k_{\ast}$ dissipate when the scattering depth of the flow is $\tau_T = k/k_{\ast}$ (due to seed electrons). The scattering depth across an individual spot $\tau_T^{\text{spot}}$ is smaller by $\epsilon$. The resultant spectrum is Wien when $\tau_T^{\text{spot}} \gg 1$. When $\tau_T^{\text{spot}} < 1$ (but the flow itself is still optically thick) cold seed photons escape the spot before being upscattered, and one may use the standard loss-probability formalism (Shapiro, Lightman, & Eardley 1976). Since the seed photons have adiabatically cooled by a factor $\sim (2\gamma_\infty)^{-2/3}(kR_0/2\pi)^{2/3}$, the accumulated $y$-parameter required to upscatter them is large, and the resulting photon index is

$$\alpha = \frac{1}{2} - \sqrt{(9/4) + (4/y)} \simeq -1. \quad (24)$$

This power law distribution [extending up to a mean energy (23)], with a superimposed Wien peak at energy (23), is the net result of this first stage of Comptonization. It compares favorably with the low energy spectra of GRBs (e.g. Cohen et al. 1996).

As dissipation continues at wavenumber $k \sim k_{\ast}$, Compton drag regulates the $y$-parameter to a value near unity. Hotspots with temperature $T_w \sim m_e c^2$ will upscatter photons above energy $\langle h\nu \rangle$ in a non-thermal tail that extends to the pair creation threshold in the wind rest frame. This in turn greatly amplifies the number of scattering charges, since photons greatly outnumber electrons in the outflow (by a factor $\sim (2\gamma_\infty)^{-2/3}(kR_0/2\pi)^{2/3}$), the accumulated $y$-parameter required to upscatter them is large, and the resulting photon index is

$$\alpha = \frac{1}{2} - \sqrt{(9/4) + (4/y)} \simeq -1. \quad (24)$$

This mechanism does not require fine-tuning of the Lorentz factor if the range of wavenumbers is broad, $k_{\text{max}} \gg k_{\text{min}}$. Nonetheless, one expects that $\gamma_\infty$ is a strong function of time in any GRB source involving an optically thick neutron

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4 As long as $\epsilon \ll 1$ photons are able to diffuse freely between spots, but not escape the wind entirely.
torus or neutron star that emits neutrinos (T94). The neutrino luminosity plausibly passes through the critical value at which the neutrino driven mass-loss rate is \( \dot{M} = L_P/c^2\gamma_e-p \); indeed the total Poynting luminosity \( L_P \) from a centrifugally supported torus is limited to \( L_P \sim 10^{51} \) erg s\(^{-1} \) in this manner.

A related model uses strong shocks to directly accelerate the photons via the first-order Fermi process (Blandford and Payne 1981). If the photons pass successively through several strong shocks separated by adiabatic cooling, then it can be shown that the number index converges to a value \(-1\) (Melrose & Pope 1993) up to an energy \( \sim L_P/\dot{N}_\gamma \).

### 3.5 A Hybrid Model: Comptonization of an MeV Bump by Non-thermal Pairs

An advected Wien photon gas with temperature \( T_0 \sim m_e c^2 \) can seed Comptonization by non-thermal pairs below the scattering photosphere. If the distribution of relativistic pairs has a lower cutoff \( \gamma_{\text{min}} = O(1) \), then the resultant spectrum breaks in the MeV range. Such a low cutoff results from a high energy pair cascade in a compact photon source (BL95), and results even for steeper pair spectra if pair creation feeds back on the leptonic acceleration efficiency to load the outflow heavily with pairs, \( n_e m_e c^2 \sim (\delta B)^2/8\pi \) (Sect. 3.3).

A further benefit of heavy pair loading is that the outflow becomes photon-starved when \( 3T_0 \) approaches \( m_e c^2 \) in the rest frame, so that the advected bump is strongly depleted by Compton upscattering above \( \sim 1 \) MeV. The resultant break energy is then (for \( \nu F_\nu \sim \text{const} \) above 1 MeV)

\[
h\nu_{br} \sim \frac{L_\gamma}{\dot{N}_\gamma} \left[ \ln \left( \frac{h\nu_{\text{max}}}{m_e c^2} \right) \right]^{-1} \sim 3T_0 \left( \frac{L_\gamma}{L_\gamma^0} \right) \left[ \ln \left( \frac{h\nu_{\text{max}}}{m_e c^2} \right) \right]^{-1}.
\]

Indeed, a narrow bump near 1 MeV appears to be the exception rather than the rule in Blazar spectra, although PKS 0208-512 does provide a spectacular exception (von Montigny et al., these proceedings).

As a model for Blazar spectra, this has a number of advantages over models involving i) direct Comptonization of photons from the central source (Dermer and Schlickeiser 1993); and ii) Comptonization of side-scattered photons (Sikora, Begelman, & Rees 1994; BL95). First, a Comptonized UV bump (Sikora et al. 1997) is avoided because the flow is self-shielding (Sect. 2.2); second, the advected Wien peak has a high enough temperature that the photon source is depleted during creation of the power-law \( \gamma \)-ray spectrum; and, third, observations of both MeV power-law breaks and isolated MeV bumps in Blazar spectra are directly tied to the electron rest energy. The duality between these two spectral states is ascribed to the presence or absence of a strong non-thermal \( e^\pm \) component.

Comptonization of an advected MeV bump can be powered either by thermal or non-thermal particles. Is it reasonable to expect that quasi-thermal motions should be the dominant Compton heat source in GRB sources, but non-thermal particles in Blazars? The key difference between these sources, aside from the central compactness, appears to be the degree of relativistic expansion. Shocks
in a GRB outflow should (locally) more closely approximate spherical surfaces, with the result that first-order Fermi acceleration is strongly suppressed in the relativistic limit (cf. Kirk, these proceedings).

4 Conclusions: Optically Thick vs. Thin Sources

We have studied dissipation in relativistic outflows that are sufficiently compact ($\ell_c > 10$) to be optically thick at the center. Advedted MeV radiation provides an interesting new source of Compton seeds in this regime. Interacting with bulk turbulent motions and non-thermal pairs near the scattering photosphere, it can manifest itself either as a MeV Blazar, or as an extended power law state when most of the available energy is converted to non-thermal pairs. The absence of a prominent MeV bump in most Blazar sources can be explained since the outflow is photon starved. The spectral signature is expected to be different when $\ell_c < 1 - 10$, or when the advected radiative flux is low. The high energy spectral break energy can cover a wider range of frequencies when synchrotron photons are the dominant seeds (MRP94; Ghisellini, these proceedings; Takahara, these proceedings).

Electron-positron pairs play a crucial role here by i) maintaining a large scattering depth near the base of the outflow and shielding the high-$\gamma$ core of a jet from ambient radiation; ii) maintaining the mean energy of the advected radiation above $\sim 1 \text{ MeV}$; iii) enhancing the efficiency of leptonic dissipative modes; and iv) reducing the minimum energy of the non-thermal pair population to $\gamma_{\text{min}} \sim 1$, which keeps the minimum energy of the Comptonized MeV photons in the MeV range. In a large-$\gamma$ (GRB) outflow, pairs also feed back on the emergent spectrum by expanding the scattering photosphere, and thus greatly increasing the range of wavenumbers that are damped by Compton drag off advected radiation.

Heavy Matter Loading and GRB Afterglow. One should also consider the effects of matter opacity in GRB outflows with extremely high central compactness. Although a core of the outflow (e.g. near the rotation axis of the central engine) must attain very high $\gamma_\infty \sim 100 - 300$, material off axis may not. Material expanding with $\gamma_\infty \sim 1 - 2$ becomes optically thin to scattering on a timescale $\sim 1$ day $(E/10^{52} \text{ erg})^{1/2}$. This is comparable to the timescale on which the optical afterglow detected from GRB970508 reached a maximum (Bond 1997). This is telling, since direct synchrotron emission between the contact discontinuity and the forward shock probably is strongly suppressed due to the relative weakness of the ambient magnetic field (Sect. 2.2). A further motivation for simultaneous optical observations of GRBs comes from the observation that the minimum frequency $N_eeB/m_e c \sim (10^2 - 10^3) eB/m_e c$ for optically thin cyclo-synchrotron emission lies near $\sim 1 \text{ eV}$ (Sect. 3.3) near the $e^\pm$ scattering photosphere and at a bulk Lorentz factor of $\sim 10^2$.

5 For example, if the central engine is a rapidly-rotating neutron star or neutron torus, then neutrino emission can easily power mass loss rates as high as $(10^{51} \text{ erg s}^{-1}/c^2 \sim 10^{-3} \text{ g s}^{-1}$.
**GRB Time Profiles: FRED vs. Chaotic.** The light-curves of GRBs show a bewildering variety of shapes (Meegan et al. 1996), but at least two main classes can be identified: bursts with smooth, asymmetric pulses (‘Fast Rise and Exponential Decay’ or FRED); and more ‘chaotic’ bursts in which narrow and wide pulses are often superimposed and in which the asymmetry of individual pulses is usually less clearly defined.

Chaotic bursts are most easily explained if the dissipation in a burst is driven by local physics within the outflow (PX94; T94; RM94; Sari & Piran 1997), rather than by interaction with an external medium. This leads to a simple discriminant between the two classes: FRED bursts arise from shells of ejecta that come into causal contact before the γ-rays escape, and vice versa for chaotic bursts. However, if the FRED bursts are also powered by local physics on a scale smaller than the width of the shell of ejecta, then the smoothness of the lightcurves indicates that the γ-ray emitting zone lies at scattering depths $\tau_T > 1$. Given the lack of a clear spectral distinction between the two classes, one reaches the same conclusion for chaotic bursts. Indeed, the transition zone between large and small $\tau_T$ can be considerably broadened by pair creation, and pairs are most effective at enhancing leptonic dissipative modes near the scattering photosphere (Sect 3.3). In sum, this leads to the following simple model: a burst is smooth (FRED) or chaotic depending on whether the scattering photosphere lies outside or inside the causal contact radius $2\gamma_\infty^2 c\Delta t$.

**Range of Temporal Frequencies and GRB Soft Tails.** A flow will, in general, be variable on a range of timescales $\Delta t$ and so dissipation can occur over a range of optical depths. The emergent spectrum varies considerably depending on whether most of the available energy resides at long timescales or short.

In this regard, it is interesting to note that the extended soft bump following GRB 870303 (a chaotic burst) detected by Ginga had a quasi-thermal cutoff (Yoshida et al. 1989), whereas the extended soft emission in GRB 960720 (a FRED burst) detected by BeppoSAX is closer to an extended powerlaw (with the possibility of a cutoff at $\sim 30$ keV; Piro et al. 1997). This spectral difference could be explained if the ejecta that produced the soft tail of GRB 870303 dissipated at $\tau_T^{-p} \gg 1$ while they were causally disconnected from the primary pulse of ejecta. High energy cut-offs to GRB afterglow are in general very diagnostic: the high energy spectral index is attracted to $\beta = -2$ below an energy $\sim (\gamma/\tau_T)m_e c^2$ (T94), and so the presence of a high energy spectral break yields information about a combination of bulk Lorentz factor and scattering depth (see also Baring & Harding 1997).
References

Baring, M.G. & Harding, A.K. 1997, Ap. J. 481, L85
Bisnovaty-Kogan, G.S. & Lovelace, R.V.E. 1997, preprint
Blandford, R.D. & Payne, D.G. 1981, M.N.R.A.S., 194, 1041
Blandford, R.D. & Eichler, D. 1987, Phys. Rep., 154 1
Blandford, R.D. & Levinson, A. 1995, Ap. J. 441, 79 (BL95)
Bloemen, H., et al. 1995, A. A., 293, L1
Blom, J.J., et al. 1995, A. A., 298, L33
Bond, H.E. 1997, IAU circular 6654; see also circulars 6655-6658
Cavallo, G. & Rees, M.J. 1978, M.N.R.A.S., 183, 359
Cohen, E., Katz, J.I., Piran, T., Sari, R., Preece, R.D., & Band, D.L. 1996, preprint.
Dermer, C.D. & Schlickeiser, R. 1993, Ap. J., 416, 458
Dermer, C.D., Miller, J.A., & Li, H. 1996, Ap. J. 456, 106
Done, C., Ghisellini, G., & Fabian, A.C. 1990, M.N.R.A.S., 245, 1
Duncan, R.C., Shapiro, S.L., & Wasserman, I. 1986, Ap. J. 309, 141
Duncan, R.C. & Thompson, C. 1992, Ap. J. 392, L9
Fatuzzo, M. & Melia, F. 1996, Ap. J. 464, 316
Fenimore, E.E., Klebesadel, R.W., & Narayan, R. 1996, Ap. J., 460, 964
Ghisellini, G., Guilbert, P.W., & Svensson, R. 1988, Ap. J., 334, 467
Haardt, F. & Maraschi, L. 1993, Ap. J. 413, 480
Hoshino, M. & Arons, J. 1991, Phys. Fluids B, 3, 818
Hurley, K., et al. 1994, Nat., 372, 652
Jaroszynski, M. 1993, A. A., 43, 183
Levinson A. 1996, Ap. J., 459, 520
Lin, R.P. 1990, in Basic Plasma Processes on the Sun, ed. E.R. Priest & V. Krishan, p. 467
Mahadevan, R., Narayan, R. & Yi, I. 1996, Ap. J. 465, 327
Mallozzi, et al. 1995, Ap. J. 454, 597
Mannheim, K. 1993, A. A. 269, 67
Maraschi, Ghisellini, & Celotti, 1992, Ap. J. 397, L5
Marcowith, A., Henri, G., & Pelletier, G. 1995, M.N.R.A.S., 277, 681
McNaron-Brown, et al. 1995, Ap. J. 451, 575
Meeghan, C.A., et al. 1996, Ap. J. 106, 65
Melrose, D.B. & Pope, M.H. 1993, Proc. Ast. Soc. Aust., 10, 222
Mészáros, P., Rees, M.H. & Papathanassiou, H. 1994, Ap. J. 432, 181 (MRP94)
Mészáros, P. & Rees, M.J. 1997, Ap. J. 482, 29
Paczyński, B. 1990, Ap. J., 363, 218
Paczyński, B. & Xu, G. 1994, Ap. J. 427, 708 (PX94)
Papathanassiou, H. & Mészáros, P. 1996, Ap. J. 471, L91
Pendleton, G.N., et al. 1996, in proceedings of the Third Huntsville Symposium on
Gamma-Ray Bursts, ed. C. Kouveliotou, M.S. Briggs & G.J. Fishman, p. 228
Piran, T. & Shemi, A. 1993, Ap. J. 403, L67
Piran, T. & Narayan, R. 1996, in proceedings of the Third Huntsville Symposium on
Gamma-Ray Bursts, ed. C. Kouveliotou, M.S. Briggs & G.J. Fishman, p. 233
Piro, L., et al. 1997, preprint
Rees, M.J. & Mészáros, P. 1992, M.N.R.A.S., 258, 41
Rees, M.J. & Mészáros, P. 1994, Ap. J. 430, 93 (LM94)
Romanova, M.M. & Lovelace, R.V.E. 1992, A. A., 262, 26
Sari, R. & Piran, T. 1997, Ap. J., 485, 270
Shapiro, S.L., Lightman, A.P., & Eardley, D.M. 1976, Ap. J., 204, 187
Sikora, M., Begelman, M.C., & Rees, M.J. 1994, Ap. J. 421, 153
Sikora, M., Madejski, G., Moderski, R., & Poutanen, J. 1997, Ap. J. 484, 108
Stern, B., Poutanen, J., Svensson, R., Sikora, M., & Begelman, M.C. 1995, 449, L13
Thompson, C. 1994, M.N.R.A.S., 270, 480 (T94)
Thompson, C. & Duncan R.C. 1995, M.N.R.A.S., 275, 255
Thompson, C. 1996, in proceedings of the Third Huntsville Symposium on Gamma-Ray
Bursts, ed. C. Kouveliotou, M.S. Briggs & G.J. Fishman, p. 802 (T96)
Thompson, C. & Blaes, O. 1997, submitted to Phys. Rev. D (TB97)
Usov, V.V. 1992, Nat., 357, 472
Usov, V.V. 1994, M.N.R.A.S., 267, 1035
Vietri, M. 1996, Ap. J. 471, L95
Yoshida, A., et al. 1989, P.A.S.P., 41, 509

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