THE OBSERVABLE EFFECTS OF A PHOTOSPHERIC COMPONENT ON GRB
AND XRF PROMPT EMISSION SPECTRUM

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ABSTRACT

A thermal radiative component is likely to accompany the first stages of the prompt emission of gamma-ray bursts (GRBs) and X-ray flashes (XRFs). We analyze the effect of such a component on the observable spectrum, assuming that the observable effects are due to a dissipation process occurring below or near the thermal photosphere. We consider both the internal shock model and a “slow heating” model as possible dissipation mechanisms. For comparable energy densities in the thermal and lepton components, the dominant emission mechanism is Compton scattering. This leads to a nearly flat energy spectrum ($\nu F_\nu \propto \nu^0$) above the thermal peak at $\approx 10^{10}–10^{11}$ keV and below $10^{10}–10^{11}$ MeV, for a wide range of optical depths $0.03 \leq \tau_{\gamma\gamma} \leq 100$, regardless of the details of the dissipation mechanism or the strength of the magnetic field. At lower energies steep slopes are expected, while above 100 MeV the spectrum depends on the details of the dissipation process. For higher values of the optical depth, a Wien peak is formed at 100 keV–1 MeV, and no higher energy component exists. For any value of $\tau_{\gamma\gamma}$, the number of pairs produced does not exceed the baryon-related electrons by a factor of larger than a few. We conclude that dissipation near the thermal photosphere can naturally explain both the steep slopes observed at low energies and a flat spectrum above 10 keV, thus providing an alternative scenario to the optically thin synchrotron-SSC model.

Subject headings: gamma rays: bursts — gamma rays: theory — plasmas — radiation mechanisms: nonthermal

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1. INTRODUCTION

The prompt emission from GRBs is believed to arise from the prompt dissipation of a substantial fraction of the bulk kinetic energy of a relativistic outflow, originating from a central compact object (for reviews, see Mészáros 2002; Waxman 2003; Piran 2005). The dissipated energy is converted to energetic electrons, which produce high-energy photons by synchrotron radiation and inverse Compton (IC) scattering. This model was found consistent with a large number of GRB observations (Band et al. 1993; Tavani 1996; Preece et al. 1998b). However, motivated by an increasing evidence for low-energy spectral slopes steeper than the optically thin synchrotron or synchrotron self-Compton (SSC) model predictions (Crider et al. 1997; Preece et al. 1998b; Frontera et al. 2000; Ghirlanda et al. 2003), an additional thermal component was suggested that may contribute to the observed spectrum (Mészáros & Rees 2000; Mészáros et al. 2002; Rees & Mészáros 2005). Indeed, thermal radiation originating from the base of the relativistic flow where the densities are high enough to provide (at least approximately) thermal equilibrium is an inevitable ingredient. This radiation is advected outward as long as the flow (presumably, jetted) material remains opaque, and eventually emerges highly collimated once the flow becomes optically thin, at a distance $r_{ph}$ from the core, which defines the photospheric radius. The comoving energy density of this thermal radiation is comparable to the energy density of the relativistic outflow, in the first stages of the flow.

The rapid variability observed in GRBs, if interpreted in the framework of the internal shocks model, implies that (at least some of) the internal collisions occur at radii smaller than the photosphere radius (Mészáros & Rees 2000; Rees & Mészáros 2005). A similar conclusion can be drawn if alternative energy dissipation mechanisms are considered, such as magnetic reconnection (Thompson 1994; Giannios & Spruit 2005). If the energy dissipation occurs below or near the photospheric radius, the thermal photons serve as seed photons to IC scattering by the relativistic electrons. For high optical depth to scattering by the electrons (and the created pairs), the optically thin synchrotron-SSC model predictions are not valid, and a different analysis is needed.

Saturated Comptonization of (synchrotron) soft photons was studied in the past in the context of X-rays emission from X-ray bursters (Sunyaev & Titarchuk 1980; Pozdniakov et al. 1983; Zdziarski 1986; Zdziarski et al. 1994). It was first suggested as a source of GRB prompt emission by Liang et al. (1997) and Liang (1997). Comptonization by thermal population of electrons was considered by Ghisellini & Celotti (1999) as an alternative source for GRB prompt emission. This model differs from the common “internal shock” model in that it assumes that the dissipated kinetic energy is continuously distributed among all the shell electrons, rather than being deposited into a small fraction of electrons that pass through the shock wave, and it assumes that the energy is equally distributed among electrons rather than following a power-law distribution. The implication of these models was further analyzed by Pe’er & Waxman (2004) (see also Ramirez-Ruiz 2005).

In a previous paper (Pe’er et al. 2005) we showed that clustering of the peak energy at 100 keV–1 MeV as observed in many GRBs (Brainerd 1998; Preece et al. 1998a) can be naturally explained as being due to a Compton-IC scattering balance, if the kinetic energy dissipation process occurs below the Thomson photosphere. This result was found to be independent of the nature of the dissipation process.

In this paper, we analyze in greater depth the effect of a photospheric term on the prompt emission from GRBs and XRFs, under different assumptions about the unknown values of the
free parameters of the emission models. In § 2 we show that the energy carried by the thermal photons is comparable to the dissipated energy for plausible values of the free parameters characterizing GRBs. This implies a large radiative efficiency, in contrast to conventional internal shock models. We analyze in § 3 the spectra obtained in the “slow heating” scenario and in § 4 the spectra obtained in the internal shock scenario. We show that these two different scenarios lead to very similar spectra for a large range of parameter values. In particular, we show here that a flat energy per decade spectrum $\nu F_{\nu} \propto \nu^{-1}$ arises naturally above a thermal peak at 30–100 keV. This provides a natural explanation for the ubiquitous $F_{\nu} \propto \nu^{-1}$ prompt emission spectra (Band et al. 1993; Preece et al. 2000) in the sub-MeV to mega–electron volt range. At lower energies, the spectrum depends on the details of the scenario considered; it is generally steeper than the typical synchrotron spectrum and can become flatter than the typical synchrotron spectrum and can become flatter above which the comoving optical depth along the jet falls to unity, lies outside the saturation radius $r_s$ of a central object of mass $M_\odot$, of the central object (e.g., black hole) is assumed, resulting in a Schwarzschild radius $r_g = 2GM/c^2 \approx 3 \times 10^6m_1$ cm. The comoving proton number density in the shocked plasma at $r_i$ is $n_p \approx L_0/(4\pi r_i^2c^2\Gamma_m^2m_e^2/\epsilon_e)$; therefore, the optical depth to Thomson scattering by the baryon-related electrons at this radius is

$$\tau_{pe} = \Gamma r_0 n_p \sigma_T = 100L_{52}\Gamma_2^{-5} \alpha^{-1}. \quad (2)$$

We thus find that for the parameters values that characterizes GRBs, indeed $r_i < r_{ph}$. We assume that a fraction $\epsilon_d$ of the kinetic energy is dissipated by the dissipation process, resulting in an internal energy density $u_{int} = L_0\epsilon_d/4\pi r_i^2c^2\Gamma^2$. Fractions $\epsilon_e$ and $\epsilon_B$ of this energy are carried by the electrons and the magnetic field, respectively; therefore, the ratio of the thermal photon energy density $u_{ph}$ to the electron energy density is

$$A \equiv u_{ph}/u_{el} = 0.44\Gamma_2^{-2/3}d_{-1}^{-1}e_e^{-0.5}, \quad (3)$$

where $\epsilon_d = 10^{-1}\epsilon_{d-1}$, $\epsilon_e = 10^{-0.5}\epsilon_{e-0.5}$, and $u_{el} \equiv \epsilon_e u_{int}$. Since the thermal photons provide the main photon reservoir for electron scattering, the ratio of synchrotron and IC emitted power by the electrons is

$$S \equiv \frac{P_{syn}}{P_{IC}} = \frac{u_B}{u_{ph}} = 1.1\Gamma_2^2d_{-1}\epsilon_e^{-0.5}, \quad (4)$$

where $\epsilon_B = 10^{-0.5}\epsilon_{B-0.5}$ and $u_B \equiv \epsilon_B u_{int}$. In order to determine the resulting spectrum, the particle energy distribution needs to be specified. In the analysis below, we consider two alternative scenarios: the slow heating scenario, and the more widely used internal shock scenario, in which power-law energy distribution of the injected electrons is assumed.

3. SLOW HEATING SCENARIO

In this scenario, as suggested by Ghisellini & Celotti (1999), the dissipated energy is continuously and evenly distributed among the shell’s electrons, during the comoving dynamical time of the dissipation process, $t_{syn} \equiv \Gamma \Delta t$. The electrons are assumed to have a Maxwellian distribution with normalized temperature $\theta_{el}$, and Comptonization is the main emission mechanism.

The electron temperature is determined by energy balance between the energy injection rate, $dE_e/dt = u_{el}/n_\text{el}t_{syn}$, and energy loss rate. For optical depth $\tau_{pe}$ not much larger than unity, the energy loss rate is given by $dE_e/dt = (4/3)c^3\sigma_T(\gamma_{\gamma\beta}^f)u_{ph}(1+S)$. Here $n_\text{el}$ is the number density of electrons in the plasma, $\gamma_{\gamma\beta}^f$ is the electron momentum, which for a relativistic Maxwell-Boltzmann distribution is related to the temperature via $\theta_{el} = \gamma_{\gamma\beta}^f/(1 + \beta^2)$. At steady state, the electron momentum is therefore given by

$$\gamma_{\gamma\beta}^f = \left[\frac{3}{4} \frac{1}{\tau_{pe}A(1+S)}\right]^{1/2}. \quad (5)$$

Since $A$ and $(1+S)$ are of order unity for parameters characterizing GRBs, $\gamma_{\gamma\beta}^f \approx \tau_{pe}^{-1/2}$.

The resulting spectrum above the thermal peak can now be approximated in the following way. For $\gamma_{\gamma\beta}^f \approx 1$ ($\tau_{pe} \leq 1$), the energy of a photon after $n_{sc}$ scattering is $\epsilon_n \approx (\gamma_{\gamma\beta}^f)^{2n_{sc}} \epsilon_0$, where

$$\theta = \frac{k_B}{m_e^2} \frac{T'}{c^2} = \frac{k_B}{m_e^2} \left(\frac{L_0}{4\pi r_i^2c^2}\Gamma_m \epsilon_{el} \epsilon_{B}\right)^{1/4} \left(2\Gamma_2^{-2/3}\right)$$

$$= 1.2 \times 10^{-3}L_{52}^2\Gamma_2^{1/3} \left(\alpha m_1\right)^{-1/2}, \quad (1)$$

where $k_B$ and $\alpha$ are Boltzmann’s and Stefan’s constants, respectively, $L_0 = 10^{52}L_{52} \text{ergs s}^{-1}$, $\Gamma = 10^2\Gamma_2$, and the characteristic mass $M \sim 10m_1 M_\odot$. Of the central object (e.g., black hole) is assumed, resulting in a Schwarzschild radius $r_g = 2GM/c^2 \approx 3 \times 10^6m_1$ cm. The recombining proton number density in the shocked plasma at $r_i$ is $n_p \approx L_0/(4\pi r_i^2c^2\Gamma_m^2m_e^2/\epsilon_e)$; therefore, the optical depth to Thomson scattering by the baryon-related electrons at this radius is

$$\tau_{pe} = \Gamma r_0 n_p \sigma_T = 100L_{52}\Gamma_2^{-5} \alpha^{-1}. \quad (2)$$
$\varepsilon_0$ is the initial photons’ energy. In the Thomson regime, the scattering rate is $dn_{ph}/dt \approx n_{ph,0}n_{el,COT} = n_{ph,0}\tau_{sync}/t_{dyn}$, where $n_{ph,0}$ is the number density of photons at the thermal peak. At the end of the dynamical time, the number density of photons that undergo $n_{el}$ scattering is therefore $n_{ph,nel} \approx n_{ph,0}\tau_{sync}$. These photons have energies in the range $\varepsilon_{nel} \ldots \varepsilon_{nuc} + d\varepsilon_{nel}$, therefore,

$$\log \left(\frac{n_{ph,nel}}{n_{ph,0}}\right) = \log \left(\frac{\varepsilon_{nel}}{\varepsilon_0}\right) \log (\gamma/\beta_T^2) = \log \left(\frac{\varepsilon_{nuc} + d\varepsilon_{nel}}{\varepsilon_0}\right) (1-1), \quad (6)$$

where $A$ is not much different from 1 and $S$ is not much larger than 1 assumed in the last equality, resulting in $(\gamma/\beta_T^2) \approx \tau_{sync}^{-1}$. Since $n_{ph,nel}$ is the number density of photons in the range $\varepsilon_{nuc} \ldots \varepsilon_{nuc} + d\varepsilon_{nel}$, we conclude that $\epsilon d\varepsilon \propto \varepsilon^{-\eta}$, or $\nu F_\nu \propto \nu^\eta$ above the thermal peak and below $\varepsilon_{max} = \gamma f_{m},e$. (2) in the frame.

For $\varepsilon_{sync} \ll 1$, the electron momenta at steady state $\gamma f_{sync} \gg 1$. An upper limit on the electron momenta $\gamma f_{sync} \sim \gamma f_{sync} \ll \gamma f_{sync} \ll \gamma f_{sync} \leq T \equiv (n_{ph,0}m_{ph}c^2 - (n_{el,0}m_{el}c^2 - 60n_{el,0}t_{sync} - 5p,0)$, where $(\gamma - 1)m_{el}c^2$ is the photon’s thermal energy and $\gamma f_{sync}$ is determined by the power-law index of the injected electrons.

We further explore the effect of a photospheric component on the observed spectrum in the more conventional internal shocks scenario, in which the dissipated energy is converted to acceleration of electrons high energies. We assume a fraction $\varepsilon_{sync} \leq 1$ of the electron number density $n_{el}$ is accelerated to a power-law energy distribution with power-law index $p$ in the range $\gamma f_{sync} \gamma f_{sync}$. A fraction $1 - \varepsilon_{sync}$ of the electrons are assumed to have a Maxwellian distribution with normalized temperature $\theta_\nu = \gamma f_{sync}/2$. Thermal photons are in the Klein-Nishina limit for Compton scattering by energetic electrons; therefore, $\gamma f_{sync}$ is determined by the equating the acceleration time and the synchrotron loss time,

$$\gamma f_{max} = \frac{2.3 \times 10^4 L_{52}^{-1/4} \epsilon_{d-1}^{-1/4} \epsilon_{0.5}^{1/4} \Gamma_2^{1/3}}{\theta_\nu^{1/2}} \frac{\gamma f_{sync}}{\gamma f_{sync}}$$

(9)

where $\alpha = 1.00$. The characteristic Lorentz factor of the accelerated electrons is

$$e_{d} \gamma f_{sync} (m_{ph}/m_{el}) = \frac{\epsilon_{d} \epsilon_{d} \gamma f_{sync} \gamma f_{sync}}{3/2 (1 - \epsilon_{d}) + \epsilon_{d} \epsilon_{d} \gamma f_{sync}} \approx 30 \epsilon_{d,0.5} \epsilon_{d,1} \gamma f_{sync}$$

(10)

where $\epsilon_{d} = 0.1$ was taken, $\log (\gamma f_{sync}/\gamma f_{sync}) \approx 7$, and a power-law index $p = 2$ of the accelerated electrons above $\gamma f_{sync}$ is assumed. Electrons having Lorentz factor $\gamma f_{sync}$ lose their energy by Compton scattering and synchrotron emission on timescale $t_{loss} = \gamma f_{sync} m_{el}c^2/(4/3) e \epsilon_{d,1} \gamma f_{sync} \gamma f_{sync} \theta_\nu (1 + S)$, and cool down to $\gamma f_{sync} \leq 1$ on timescale $t_{loss} = \gamma f_{sync}$. The ratio of the cooling time of electrons at $\gamma f_{sync}$ to the dynamical time is calculated using equations (2), (3), (4), and (10):

$$\frac{t_{loss} (\gamma f_{sync})}{t_{dyn}} = \frac{1}{(4/3) A (1 + S) (\gamma f_{sync}/\gamma f_{sync})} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}} \frac{1}{\gamma f_{sync}}$$

(11)

We thus conclude that for Lorentz factor $\Gamma < 2 \times 10^4$, $\gamma f_{sync}$ directly implies $t_{loss} (\gamma f_{sync}) < t_{dyn}$, and electrons lose all their energy during the dynamical time.

If the optical depth to scattering is low, $\gamma f_{sync} < 0.03(3/2) \gamma f_{sync}$, electrons at $\gamma f_{sync}$ maintain their energy during the dynamical time. The resulting IC spectrum in the range $\theta_{sync} c m_{el} c^2$, $\gamma f_{sync} c m_{el} c^2$ is $I_{\nu} \propto \nu^2$, or $\nu F_\nu \equiv \nu^2 d\nu/d\varepsilon \propto \varepsilon^2$. At higher energies, the spectral index is determined by the power-law index of the injected electrons. Above $\gamma f_{sync} c m_{el} c^2$ the Klein-Nishina effect significantly modifies the spectrum, and numerical treatment is required (see the numerical results in § 5).

For an optical depth $0.03(3/2) \gamma f_{sync} \leq 1$, electrons lose their energy and accumulate at $\gamma f_{sync} \sim 1$ at the end of the dynamical
time. The energy loss rate of electrons at Lorentz factor $\gamma < \gamma_{\text{char}}$ is $p(\gamma) \propto \gamma^2 \tau_{\text{opt}} (1 + S) \propto \gamma^2$; therefore, the electron distribution below $\gamma_{\text{char}}$ is $n_{\text{el}}(\gamma) \propto \gamma^{-2}$, and the resulting IC spectrum in the range $\delta m u c^2 \cdots \gamma_{\text{char}} m u c^2$ is $\nu \phi_\nu \equiv \nu^2 d\nu/d\nu \propto \nu^\alpha$ with $\alpha = 3/2$.

On top of the IC component there is the synchrotron component from the power-law accelerated electrons, with spectral index $\alpha = 1 - p/2$. The observed spectral index in this range is therefore expected to be $\alpha \approx 0.5$; the exact value depends on the details of the scenario, i.e., the values of $\epsilon_p$, $S$, and $p$. In §5 we present our numerical results in this case.

For higher optical depth, $\tau_{\nu_e} > 1$, electrons lose their energy by IC scattering and synchrotron emission and accumulate at $\nu_{\text{peak}}$. We have previously shown (Pe’er et al. 2005) that in this scenario the energy gain rate of electrons via direct Compton scattering energetic photons. We show in the Appendix that a factor of 2–3.

$e$ > $f m u c^2$ where $f \approx 3$ interacting with low-energy electrons having $\gamma \approx 1$, lose energy at an energy-independent rate, $d\nu/d\nu \approx c n_{\text{el}}(\nu u c^2)/2)$. The injection rate of energetic photons (via IC scattering and synchrotron emission) is approximated using energy considerations, $d\nu_{\text{inj}}(\nu > f m u c^2)/d\nu \approx u_{\text{inj}}(f m u c^2)^2/2)$. These photons lose energy by downscattering to energies below $f m u c^2$ on a timescale $\nu_{\text{inj}}(\nu u c^2)/d\nu = 2 f u_{\text{inj}}/\tau_{\nu_e}$; therefore, at steady state the number density of energetic photons is given by $n_{\text{ph}}(\nu > f m u c^2) = 2 u_{\text{inj}}/\nu_{\text{char}}$. This result implies that the energy gain rate of electrons via direct Compton scattering is equal to the energy gain rate in the slow heating mechanism, $d\nu/e dt \approx u_{\text{inj}}/\nu_{\text{char}}$. (see also Pe’er et al. 2005); therefore, for $\tau_{\nu_e}$ not much larger than unity, equation (5) characterizes the electron momenta at the end of the dynamical time in this scenario as well. We have previously shown (Pe’er et al. 2005) that in this case of high optical depth, the electrons accumulate at $\gamma_{\nu} \beta_{\nu} \approx 1$ with a very weak dependence on the values of the unknown parameters, and that for optical depth higher than a few tens a Wien peak is formed at 1–10 keV. Equation (7) provides the asymptotic value of the Wien peak energy for very high optical depth in this scenario as well. For $\tau_{\nu_e}$ of a few tens, the thermal peak is upscattered to energy in the range $\delta m u c^2 < \nu_{\text{peak}} < \nu_{\text{WP}}$ (see the numerical results in §5).

In contrast to the slow heating scenario, a large number of pairs can in principle be created, following the injection of particles to high energies. However, the following argument, based on the analysis of Pe’er & Waxman (2004), suggests that the number of pairs do not outnumber the baryon-related electrons by a factor of larger than a few, for any value of the optical depth $\tau_{\nu_e}$. From energy considerations, the pair injection rate is limited by $dn_e/d\nu \approx u_{\text{inj}}(\nu u c^2)^2/2)$. The pair annihilation rate is $\sim n_{\text{el}}^2 c \sigma_T$; therefore, the ratio of pairs to baryon-related electrons in steady state is

$$n_{\text{el}} \approx \left(\frac{\gamma_{\text{char}}}{\tau_{\nu_e}}\right)^{1/2}. \tag{12}$$

We therefore do not expect pairs to outnumber the baryon-related electrons by a number larger than a few for any value of the optical depth. The creation of pairs lowers the number density of energetic photons, thus lowering the value of $\gamma_{\nu} \beta_{\nu}$ by a small factor (see Pe’er et al. 2005).

5. NUMERICAL CALCULATIONS

5.1. The Numerical Model

In order to confirm the analytical calculations presented above and to derive the spectra resulting from different values of the free parameters, we calculated numerically the photon and particle energy distribution under the assumptions of §§3 and 4. In our numerical calculations we used the time-dependent numerical code presented in Pe’er & Waxman (2005). This code solves self-consistently the kinetic equations for $e^\pm$ and photons describing cyclosynchrotron emission, synchrotron self-absorption, direct and IC scattering, pair production and annihilation, and the evolution of a high-energy electromagnetic cascade. These equations are being solved during the dynamical time of the dissipation process.

In this version of the code, we assume the existence of a thermal photospheric component, which at the dissipation radius $r_i$ is characterized by time-independent luminosity $L_{\nu_i} = L_{\nu_0}(2\Gamma)^{-2/3}$ and temperature given by equation (1). A fraction $\epsilon_{\nu_0} = 0.1 \epsilon_{\nu_0} / \epsilon_{\nu_0} - 1$ of the kinetic energy is dissipated by the dissipation process. A fraction $\epsilon_{\nu_i}$ of the dissipated energy is carried by the electrons, and a fraction $\epsilon_{\nu_i}$ is carried by the magnetic field. In the slow heating scenario, the dissipated energy is assumed to be continuously distributed among the electrons in the two shells (the created pairs), which assume a Maxwellian distribution with normalized temperature $\theta_{\nu_i}(t)$. The electron temperature is determined self-consistently at each time step by balance of energy injection and energy loss. In the internal shock scenario, electrons are assumed to be injected by the shock waves at a constant rate. A fraction $\epsilon_{\nu_0}$ of the electron population is injected into a power-law energy distribution with power-law index $p$ between $\gamma_{\text{char}}$ and $\gamma_{\text{max}}$ (see eqs. [9] and [10]). The remaining fraction $1 - \epsilon_{\nu_0}$ of the injected electrons assume a Maxwellian distribution with temperature $\theta_{\nu_i} = \gamma_{\text{max}} / 2 \approx 20 \epsilon_{\nu_i} = 0.5 \epsilon_{\nu_0} / \epsilon_{\nu_0} - 1$.

5.2. Slow Heating Scenario

Numerical results of the observed spectra at the end of the dissipation process in the slow heating scenario are presented in Figure 1 for different values of the optical depth $\tau_{\nu_e}$. In producing the plots in this graph, we chose representative values of the free parameters, $\alpha = 1 / \Gamma = 10^{-2}$, $\epsilon_{\nu_0} = 10^{-1}$, and $\epsilon_{\nu_i} = \epsilon_{\nu_i} = 10^{-0.5}$. The differences in the plots are due to different intrinsic
luminosity, \( L_0 = 10^{46} - 10^{53} \) ergs s\(^{-1}\), which result in a wide range of optical depths, \( 10^{-2} \leq \tau_{\gamma_e} \leq 10^3 \) (see eq. [2]). A similar graph is obtained by assuming different value of \( \Gamma \) and constant luminosity, where variation in the value of \( \Gamma \) by a factor of 3 leads to variation in \( \tau_{\gamma_e} \) by a factor of 10.25.

In the case of low optical depth \( \tau < 1 \), one expects the thermal peak at the observed energy \( 3E \Gamma m_e c^2/(1 + z) \approx 10 \)–100 keV, caused by the advected thermal photons. In the scenario of \( \tau = 0.01 \), the characteristic Lorentz factor of the electrons is \( \gamma_{\gamma} \approx 10 \), and synchrotron radiation produces the peak observed at low energy, \( \gamma \approx 300 \) eV. For higher optical depth the characteristic electron Lorentz factor is lower, \( \gamma_{\gamma} \approx 1 \) (see eq. [5]), and the synchrotron emission occurs at lower energies, below the synchrotron self-absorption frequency (see the analysis of synchrotron self-absorption frequency as a function of the free parameters values in Pe’er & Waxman [2004]).

The high value of the characteristic Lorentz factor of the electrons in the scenario with \( \tau = 0.01 \) implies that two scatterings are required to upscatter thermal photons to \( \sim 100 \) MeV; therefore, Compton scattering of the synchrotron and thermal photons produces the wavy shape of the spectrum above 10 keV.

At moderate values of the optical depth, \( \tau \sim 0.1 \)–10, multiple Compton scattering produces the flat spectrum observed above the thermal peak, in the range 10 keV–10 MeV, in agreement with the predictions of § 3. At even higher values of the optical depth, \( \tau_{\gamma_e} \approx 100 \), a Wien peak is formed at sub-MeV energy. As explained in Pe’er & Waxman [2004], the high optical depth implies that in this case photons are not observed prior to a period of adiabatic expansion, during which 50%–70% of the photons’ energy is converted to kinetic energy (bulk motion) of the plasma. The observed Wien peak is therefore expected at \( \sim 300 \) keV (see also Pe’er et al. 2005).

5.3. Comparison of Scenarios: Low Optical Depth

Figure 2 shows comparison of the spectra obtained for the two dissipation scenarios considered (slow heating and internal shocks) and two different values of the magnetic field, in the scenario of very low optical depth, \( \tau_{\gamma_e} = 4 \times 10^{-3} \). For this value of the optical depth, only a small fraction of the electrons’ energy is radiated. If the magnetic field is weak (dash-dotted lines), the advected thermal peak at 10 keV is prominent. In the slow heating scenario (thin dash-dotted line) an IC peak at \( \sim 15 \) MeV, created by electrons having \( \gamma_{\gamma} \beta_e = 30 \) is expected. The thin solid line, representing the slow heating scenario with equipartition magnetic field, shows four peaks: the synchrotron peak at 300 eV, the thermal peak at 10 keV, and two peaks produced by IC scattering the low-energy peaks, at 300 keV and 10 MeV. The combined effect of these peaks is a wavy-shaped, nearly flat spectrum in the range 100 eV–100 MeV.

An internal shock scenario with weak magnetic field is shown by the thick dash-dotted line in Figure 2. Electrons injected with a power-law index \( p = 2 \) between \( \gamma_{\max} \) and \( \gamma_{\gamma} \) produce the high-energy component by IC scattering the thermal photons. Electrons with Lorentz factor \( \gamma < \gamma_{\gamma} = (400)^{-1} \sim 3 \times 10^2 \) cool by IC emission faster than the dynamical time; therefore, the resulting IC spectrum is nearly flat, \( v_F = c \gamma \delta m_e c^2/\delta c \sim c^3 \gamma_{\gamma} \), with \( \alpha = 0 \) below \( \Gamma^{-2} 0.6 \gamma_{\gamma} c^2/(1 + z) \approx 30 \) GeV. For energetic electrons with \( \gamma > \gamma_{\gamma} \), the thermal photons are in the Klein-Nishina regime, and as a result these electrons cool slower than the dynamical time, and the spectral slope of the IC spectrum at higher energies is \( \alpha = \frac{1}{2} \). If magnetic field is added (thick solid curve), electrons at all energies cool fast by synchrotron emission, and the resulting synchrotron spectrum is flat in the range 100 eV–10 GeV.

Figure 2—Time-averaged spectra obtained for low values of the optical depth, \( \tau_{\gamma_e} = 4 \times 10^{-3} \). Results for slow dissipation are shown by thin lines, and dissipation by shock waves with power-law index \( p = 2 \) and \( \epsilon_p = 1 \) are shown by thick lines. Solid lines are for high magnetic field, \( \epsilon_B = 10^{-0.5} \), and dash-dotted lines are for \( \epsilon_B = 10^{-4} \); \( \epsilon_p = 1 \) has cosmological parameters as in Fig. 1. [See the electronic edition of the Journal for a color version of this figure.]

5.4. Comparison of Scenarios: Intermediate and High Optical Depth

For low to moderate values of \( \tau_{\gamma_e} \), equation (5) provides a good estimate of the electron momenta. The ratio between the energy of the synchrotron peak \( \varepsilon_{\gamma_{\gamma}} \) and the energy of the thermal peak \( \varepsilon_{th} = 3h_\lambda/c^2 \) is given by

\[
\frac{\varepsilon_{\gamma_{\gamma}}}{\varepsilon_{th}} = \frac{3h \gamma_{th} \gamma_{\gamma}}{2m_e c^2} \frac{(\gamma_{\gamma} \beta_e)^2}{\varepsilon_{th}} = \frac{\pi a^3}{8} \frac{h}{k_B} \frac{\Gamma^{1/3}}{\tau_{\gamma_e}} \frac{(\epsilon_B \epsilon_p / \varepsilon_{th})^{1/2}}{A + 1} \frac{1}{\tau_{\gamma_e}}
\]

where \( \varepsilon_{th} = 10^{4.5} \) eV is the observed energy of the thermal peak. We therefore conclude that for \( \tau_{\gamma_e} \gg 0.1 \), the peak energy of the synchrotron component is much lower than the peak energy of the thermal component, and is below the self-absorption frequency for \( \tau_{\gamma_e} \geq 1 \). This conclusion is confirmed by the numerical results of the slow heating scenario (thin lines) with strong (solid lines) magnetic field, presented in Figures 3 and 4 for \( \tau_{\gamma_e} = 0.1 \) and 1 respectively. A low-energy synchrotron component at \( \sim 200 \) eV is prominent for \( \tau_{\gamma_e} = 0.1 \) (Fig. 3) and is much weaker for \( \tau_{\gamma_e} = 1 \) (Fig. 4).

Figures 3–5 provide a comparison of the spectra obtained for the two dissipation mechanisms, equipartition versus low magnetic field, and two different values of \( \epsilon_p \). The internal shock scenario is shown with the thick lines in Figure 3. The solid line of this scenario the the scenario of equipartition magnetic field and \( \epsilon_p = 0 \); i.e., the electrons are injected at \( \gamma_{\gamma_{\gamma}} \) and then cool to \( \gamma_{\gamma} \). Synchrotron emission dominates the spectrum below the thermal peak, at 0.1–10 keV, and IC scattering dominates at higher energies. The electron distribution below \( \gamma_{\gamma_{\gamma}} \) is \( \epsilon_p \) (\( \gamma_{\gamma} \)) \( \sim 2 \); therefore, the emitted radiation has spectral slope \( \alpha = \frac{1}{2} \). IC scattering by electrons accumulated at \( \gamma_{\gamma} \) results in a somewhat flatter (\( \alpha = \frac{1}{2} \)) spectral slope. The thick dash-dotted curve shows a scenario with weak magnetic field (\( \epsilon_B = 10^{-4} \)) and \( \epsilon_p = 1 \). In this scenario, energetic electrons upscatter photons to high energies; however,
pair production suppresses emission above giga–electron volts. The number of pairs created in this scenario is found numerically to be $n_{\perp}/n_{\parallel} \simeq 4$, in agreement with the analytical predictions of § 4.

The numerical results obtained for $\tau_{ve} = 1$ and 10 are presented in Figures 4 and 5. These results further show the similarity between the spectra obtained by the two different dissipation mechanisms, as well as the weak dependence of the spectral shape on the value of the magnetic field or on $v_{pl}$. Multiple scattering creates the flat energy per decade spectrum above the thermal peak at $3\Gamma m_{e}c^2/(1 + z) \approx 30–100$ keV, for a typical $\Gamma = 100–300$. This energy is very similar, albeit somewhat lower than the observed spectral break energy, $\sim 300$ keV (Band et al. 1993; Preece et al. 1998a, 2000). The flat spectrum extends up to $\Gamma(\gamma/\bar{\gamma})m_{e}c^2/(1 + z) \approx 10$ MeV.

In the internal shock scenario, the energetic photons produce pairs, resulting in a cutoff above $\sim 100$ MeV. The number density of pairs in the scenario $\tau_{ve} = 1$ is found numerically to be $n_{\perp}/n_{\parallel} \simeq 4$ ($v_{pl} = 0$), $n_{\perp}/n_{\parallel} \simeq 5$ ($v_{pl} = 1$). Similar, but somewhat lower values ($n_{\perp}/n_{\parallel} = 1.5, 3.5$, respectively) are obtained for $\tau_{ve} = 10$. We thus find that, indeed, pairs do not outnumber the baryon-related electrons by a large factor, due to pair annihilation. The pair annihilation phenomenon causes the small peak observed at $\Gamma m_{e}c^2/(1 + z) = 25$ MeV (see Pe’er & Waxman 2004), for
further discussion). In summary, for intermediate values of the optical depth $\tau_{\gamma e} \approx 0.1$–10, multiple IC scattering produces an approximately flat energy per decade spectrum in the range $\approx \text{keV}–\text{sub-GeV}$ for the two dissipation mechanism considered regardless of the values of the magnetic field or of $\epsilon_{\text{pl}}$.

Figure 6 shows the asymptotic case of high optical depth, $\tau_{\gamma e} = 100, 1000$. For optical depths larger than ~100, multiple Compton scattering leads to the formation of a Wien peak at sub-MeV energies, regardless of the details of the acceleration mechanism, or the values of any of the free parameters (see also Pe’er et al. 2005). As shown in § 4, for high values of the optical depth the number density of pairs is strongly suppressed by pair annihilation. This calculation is confirmed by the numerical results, $n_{\perp}/n_{\text{el}} \approx 1$.

### 6. SUMMARY AND DISCUSSION

In this work, we have considered the effect of a photospheric component on the observed spectrum after kinetic energy dissipation that occurs below or near the thermal photosphere of the outflows in GRBs or XRFs, with particular emphasis on the resulting spectral shape. Two dissipation mechanisms were considered: the slow heating mechanism, and the internal shock scenario. We showed that the resulting spectra are largely independent of the details of the dissipation mechanism or of the values of any of the free parameters. A strong dependence was found only on the value of the optical depth, $\tau_{\gamma e}$. For $\tau_{\gamma e}$ smaller than a few tens, an approximately flat energy per decade spectrum is obtained above a break energy at few tens to hundreds of kilo–electron volts and up to sub-GeVs, while steep slopes are obtained below the break energy (see Figs. 4–5). For higher values of the optical depth, a Wien peak is formed at sub-MeV energies, with steep slope at lower energies, and sharp cutoff above this energy. The production of pairs does not change this result, since pair annihilation limits the ratio of the number density of pairs to baryon–related electrons to a factor of not larger than a few.

Our theoretical and numerical results emphasize the important role of IC scattering in the formation of the spectra in scenarios involving an important thermal component. As a consequence, the observed spectral slopes in the two model considered are significantly different from the synchrotron model prediction. The (nearly) thermal component observed for low to intermediate values of the optical depth, and the Wien peak formed for high value of the optical depth, can account for the increasing evidence for steep spectral slopes observed below ~100 keV in the early stages of GRBs (Preece et al. 2002; Ghirlanda et al. 2003; Ryde 2004). As shown in Pe’er et al. (2005) and as confirmed here (Figs. 1 and 6), for a scattering optical depth larger than ~100, dissipation at a subphotospheric radius can naturally lead to a clustering of the peak energy at ~100 keV–1 MeV, as observed in many bursts (Brainerd 1998; Preece et al. 1998a, 2000).

For low values of the optical depth, $\tau_{\gamma e} \approx 0.1$–1, the advected thermal component produces an observed peak at a few tens of kilo–electron volts. This peak energy is consistent with the peak energy observed in XRFs at ~25 keV (Heise et al. 2001) and X-ray–rich GRBs (XRRs), which show peak energy clustering at ~50 keV (Sakamoto et al. 2005). Thermal peak energies in this range of few tens of kilo–electron volts are obtained for a wide range of values of the free model parameters, $10^{-4} \leq L_2 \leq 1$, $1 \leq \Gamma_2 \leq 3$, and $1 \leq \alpha \leq 10^3$, which result in an optical depth ~1 (see eqs. [1] and [2]). It is difficult to obtain in our model a thermal peak energy at energies below a few keV combined with an optical depth larger than ~$10^{-3}$. Therefore, for a peak below a few kilo–electron volts, electrons maintain their energy (see § 4).

We conclude that if XRF peak energies are due to advected thermal component, no XRFs with peak energy below a few kilo–electron volts are expected.

The inclusion of an advected thermal component in the prompt emission calculations can help resolve the question of the accelerated electron distribution, which is of high theoretical importance. As pointed out by Baring & Braby (2004), fitting the data of many GRBs by synchrotron and Compton emission only requires the acceleration of the majority of electrons to a power-law distribution, leaving far too few thermal particles than to be expected as a seed distribution for the power-law population. Here we overcome this problem by assuming that the thermal component seen in many bursts (e.g., Ryde 2004), and presumably existing in all of them at early times, is advected from the core. As a result, a variety of spectra can be obtained under different assumptions on the values of the free parameters, without the requirement that most of the electrons are accelerated to a power-law distribution (in the presented figures of the power-law acceleration model, we allow the parameter $\epsilon_{\text{pl}}$ to vary between 0 and 1).

The spectra presented in this paper (§ 5) describe the emission resulting from a single dissipation phase (e.g., single collision between two shells) for a particular choice of model parameters. Observed spectra are expected to be combination of spectra produced by many such dissipation processes, which are characterized by different parameters, e.g., different locations of collisions within the expanding wind. A detailed comparison with observations therefore requires a detailed model describing the distribution of single dissipation parameters within the fireball wind (see, e.g., Daigne & Mochkovitch 1998; Panaitescu et al. 1999; Guetta et al. 2001). The construction and investigation of such detailed models would involve additional assumptions and parameters, which is beyond the scope of this manuscript.

Nonetheless, the similarities found between the resulting spectra, produced under reasonable assumptions about the unknown values of the free parameters, suggests that the key results would remain valid for a spectrum that is obtained from a series of dissipation processes. These results include an observed break energy at few tens of kilo–electron volts to sub-MeVs, accompanied by slopes steeper than the optically thin synchrotron spectrum at lower energies, which are obtained for a single dissipation event for any value of the optical depth, and flat energy per decade spectra obtained up to sub-GeV energy for optical depths of $10^{-2}$ to a few tens. For a superposition of dissipation events at various radii, flatter slopes can be obtained.

A superposition of dissipation events is also required in our model in order to explain both peak clustering at ~300 keV and flat energy per decade spectra at higher energies, as observed in many GRBs. A high optical depth $\tau_{\gamma e} \sim 100$ is required in order to obtain a Wien peak at sub-MeV energies, while lower values of $\tau_{\gamma e}$ are required in order to obtain a flat energy per decade spectra at higher energies. As detailed in Pe’er & Waxman (2004) and Pe’er et al. (2005), for the case of high optical depths the photons are expected to lose ~30% of their energy to the bulk motion of the plasma during an adiabatic expansion phase before escaping; thus, the observed Wien peak energy is expected at ~300 keV. We conclude that given such a superposition, a clustering of the peak energy, steep slopes below the peak, and flat energy per decade spectra above the break are a natural consequence of a model containing a significant thermal emission component. These results show good agreement with a large number of observations of GRB prompt emission spectra (Band et al. 1993; Preece et al. 2000; Ryde 2004).
The fluxes predicted by the model are within the detection capability of the Swift\(^4\) satellite in the kilo–electron volt range and the Gamma-Ray Large Area Space Telescope\(^5\) satellite in the sub-GeV range. Observations of a cutoff at high energies will therefore provide information about the optical depth during the dissipation phase, hence constraining the value of \(\Gamma\), one of the least constrained free parameters of the model.

The results may be applicable to a variety of compact objects, such as GRBs, XRFs, and blazars. At least 60 flat-spectrum radio quasars having flat \((\alpha \approx 0)\) spectral indexes in the Energetic Gamma-Ray Experiment Telescope range were reported (e.g., Mukherjee et al. 1997). Similar spectra were obtained for several quasars using the BeppoSAX satellite (Tavecchio et al. 2000), while clustering of blazars peak energies at 1–5 MeV was reported by McNaron-Brown et al. (1995). We thus conclude that similar mechanisms to the ones presented here may explain some of the observable effects in radio quasars as well.

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APPENDIX

ENERGY LOSS RATE OF HIGH-ENERGY PHOTONS DUE TO COMPTON SCATTERING WITH MILDLY RELATIVISTIC ELECTRONS

The rate of scattering by a single electron having Lorentz factor \(\gamma\) passing through space filled with a unit density, isotropically distributed, monoenergetic photons with energy \(\alpha_1 m_e c^2\) was first derived by Jones (1968) and summarized in Pe’er & Waxman (2005):

\[
\frac{d^2N(\gamma, \alpha_1)}{dt d\alpha} = \frac{\pi r_0^2 c \alpha}{2 \gamma^4 \beta \alpha_1^2} [F(\zeta_+) - F(\zeta_-)].
\] (A1)

Here \(\alpha\) is the energy of the outgoing photon in units of \(m_e c^2\), \(r_0\) is the classical electron radius, \(\beta = (1 - 1/\gamma^2)^{1/2}, \zeta_{\pm}\) are the upper and lower integration limits, and \(F(\zeta)\) is given by the sum of 12 functions, \(F(\zeta) = \sum_{i=1}^{12} f_i(\zeta)\), where the functions \(f_i(\zeta)\) are summarized in, e.g., equation (20) of Pe’er & Waxman (2005).

The limits \(\zeta_{\pm}(\alpha; \gamma; \alpha_1)\) depend on the energy of the outgoing photon, \(\alpha\), for given \(\alpha_1\) and \(\gamma\) (see Pe’er & Waxman 2005). For \(\alpha_1 \gg 1\) and \(\gamma \approx 1\), this dependence is degenerated and \(\zeta_{\pm} = 1 \pm \beta\). In this approximation, in calculating the difference of the 12 terms, \(\Delta f_i \equiv f_i(\zeta_+) - f_i(\zeta_-)\), only four out of the 12 terms are of order \(\alpha_1^{-1}\) or \(\alpha^{-1}\), while the other eight are of lower orders in \(\alpha\), \(\alpha_1\). Thus, to leading order in \(\alpha, \alpha_1\),

\[
\Delta f_5 \approx f_5(\zeta_+) - f_5(\zeta_-) \approx \frac{2\beta}{(1 - \beta^2)\alpha_1},
\]

\[
\Delta f_9 \approx \frac{2\beta(1 + \beta^2)}{\alpha},
\]

\[
\Delta f_{10} \approx \frac{\beta^3}{\alpha},
\]

\[
\Delta f_{11} \approx \frac{2\beta(1 + \beta^2)}{\alpha^2} \left[ \alpha_1 - \alpha + \left(1 + \frac{\alpha_1}{\alpha}\right) \left(1 + \frac{\beta^2}{2}\right) \right].
\] (A2)

The (normalized) energy loss rate of photon being scattered by unit density, isotropically distributed, monoenergetic electrons at Lorentz factor \(\gamma\) is therefore given by

\[
\frac{d\alpha(\gamma, \alpha_1)}{dt} = \int_{\alpha_{\text{min}}}^{\alpha_1} \frac{d^2N(\gamma, \alpha_1)}{dt d\alpha} d\alpha' \sim c \sigma_{\text{T}} \left(1 - \frac{\beta^2}{2} + \frac{3\beta^2}{16}\right) = c \sigma_{\text{T}} \left(1 - \frac{5\beta^2}{16}\right) + O(\alpha_1^{-1}).
\] (A3)

In order to check the validity of this approximation for values of \(\alpha_1\) not much larger than 1, a numerical integration of equation (A1) was carried out, using the exact formula. The results of the energy loss rate are presented in Figure 7. From the figure we find that the approximation \(d\varepsilon/dt \approx c \sigma_{\text{T}} m_e c^2 / 2\), where \(\varepsilon = \alpha m_e c^2\) for the energy loss rate of photon due to Compton scattering with unit density, isotropically distributed, monoenergetic electrons with \(\gamma \approx 1\) is valid for \(\varepsilon \geq f m_e c^2\), with \(f \approx 3\).
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Fig. 7.— Energy loss rate of photon with initial energy $\alpha_1 n_e c^2$ due to Compton scattering with unit density, isotropically distributed, monoenergetic electrons with velocity $\beta = 0.1$. [See the electronic edition of the Journal for a color version of this figure.]