DISSIPATIVE PHOTOSPHERE MODELS OF GAMMA-RAY BURSTS AND X-RAY FLASHES

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ABSTRACT

We consider dissipative effects occurring in the optically thick inner parts of the relativistic outflows producing gamma-ray bursts and X-ray flashes, emphasizing in particular the Comptonization of the thermal radiation flux that is advected from the base of the outflow. Such dissipative effects—e.g., from magnetic reconnection, neutron decay, or shocks would boost the energy density of the thermal radiation. The dissipation can lead to pair production, in which case the pairs create an effective photosphere farther out than the usual baryonic one. In a slow dissipation scenario, pair creation can be suppressed, and the effects are most important when dissipation occurs below the baryonic photosphere. In both cases an increased photospheric luminosity is obtained. We suggest that the spectral peak in gamma-ray bursts is essentially due to the Comptonized thermal component from the photosphere, where the comoving optical depth in the outflow falls to unity. Typical peak photon energies range between those of classical bursts and X-ray flashes. The relationship between the observed photon peak energy and the luminosity depends on the details of the dissipation, but under plausible assumptions can resemble the observed correlations.

Subject headings: gamma rays: bursts — gamma rays: theory — X-rays: bursts

1. INTRODUCTION

Most gamma-ray burst (GRB) models invoke a relativistic outflow, probably channeled into a jet, that is energized by a central compact object. The gamma-ray and hard X-ray emission is attributed to dissipative processes in the jet. The outflow is inferred to be unsteady on timescales down to 1 ms; indeed, internal shocks are the most widely discussed dissipative process because of their ability to convert bulk kinetic energy into relativistic electrons, which then radiate (e.g., via synchrotron emission) on very short timescales. The outflow would carry baryons and also magnetic fields (which may carry as much power in Poynting flux as does the baryon kinetic energy). However, there is another inevitable ingredient of the outflow: thermal radiation. This radiation originates near the base of the outflow, where densities are high enough to guarantee (at least approximate) thermal equilibration. This thermal radiation is advected outward, as long as the jet material remains opaque, and emerges highly collimated from a "photosphere" in which the jet became optically thin.

A laminar and steady jet, viewed head-on, would give rise to emission with a thermal spectrum peaking in the hard X-ray or gamma-ray band. Moreover, the comoving energy density of this blackbody radiation could be at least comparable with that of the magnetic field. So, if dissipation generates relativistic electrons and suprathermal pairs, their energy losses due to Compton scattering of the thermal radiation would be competitive with those from synchrotron emission—perhaps even dominant. Consequently, when dissipation occurs (e.g., via internal shocks) the outcome may be a "hardened" (graybody) thermal component, along with a power-law component extending to higher photon energies. We suggest that the photon energy $E_{\rm pk}$ at which GRB spectra reach a peak may be the (probably Comptonized) ther-mal peak. We discuss how, under this hypothesis, $E_{\rm pk}$ depends on the parameters characterizing the GRB.

A key parameter in the outflow is plainly the photospheric radius, the radius at which the comoving optical depth along the jet falls to unity. In calculating this radius, we must allow for the possibility that the electrons associated with the baryons are outnumbered by electron-positron pairs (e.g., Eichler & Levinson 2000). Moreover, the number of pairs may be greatly increased by dissipative processes. The details depend on whether the photosphere lies inside or outside the saturation radius at which the bulk Lorentz factor Γ asymptotes to the dimensionless entropy $\eta = L_0/(\dot{M}c^2)$, where L_0 and $\dot{M}$ are the total energy and mass outflow rates, respectively. For a spherical flow in which the free energy emanates from a central region, $r_0 \sim \alpha r_g = \alpha 2GM/c^2$, comparable to the Schwarzschild radius r_g of a central object of mass M (where $\alpha \geq 1$), the bulk Lorentz factor grows as $\Gamma(r) \sim r/r_0$ outside of r_0 up to a saturation radius $r_s \sim r_0\eta$, where it saturates at a value $\Gamma \sim \eta$. This simple behavior applies for a spherical outflow (or a conical one with jet opening half-angle $\theta_j < \Gamma^{-1}$) in which there are no internal shocks. We focus on this as an illustrative case (bearing in mind that the effective value of r_0 may be increased by dissipation in the inner jet). Moreover, extensions to the cases of convergent or divergent jets are straightforward. Inside the saturation radius, the observer-frame photospheric luminosity L_γ is approximately the total luminosity of the outflow L_0, since the increasing Doppler boost just cancels the adiabatic decay of the comoving characteristic photon energy. On the other hand, if the photosphere of an adiabatic flow occurs outside the saturation radius $r > r_s$, the Lorentz factor no longer grows, and $L_\gamma(r) = L_0(r/r_s)^{-2/3} < L_0$, the greater part of the energy being in kinetic form, $L_k \sim L_0$ (e.g., Mészáros & Rees 2000). If this photospheric luminosity were the bulk of the observed radiation, the radiative efficiency would be low in the latter case.

However, the above scenario can change substantially due to dissipative effects such as magnetic reconnection (e.g., Thompson 1994; Giannios & Spruit 2005), neutron decay (e.g., Beloborodov 2003), or internal shocks. If the dissipation occurs below the photosphere, the adiabatic decrease of the radiative luminosity

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beyond the saturation radius can be compensated by reconversion of some fraction $\epsilon_d \leq 1$ of the kinetic energy into radiation, which would reenergize the photospheric component. Moreover, dissipation outside the nominal photosphere may lead to sufficient pair formation to create a second photosphere, lying outside the original nominal photosphere that would have been obtained in the absence of dissipation.

Thus, if there were subphotospheric dissipation, the observable photospheric luminosity would be boosted by the energy recovered from the kinetic energy, which becomes available for conversion into radiation or pairs. Moreover, the dissipated energy would go mainly into Comptonization of the thermal radiation advected out from the central engine. Above r_s, the photospheric luminosity can be boosted to a value $L_\gamma = \epsilon_d L_0 > L_0 (r/r_s)^{-2/3}$, depending on the dissipation efficiency. We suggest that the peak energy of the photon spectrum of gamma-ray bursts should be identified with the peak of this Comptonized spectrum.

2. PHOTOSPHERES, DISSIPATION, AND PAIRS

In the dissipation regions of the flow, all suprathermal or relativistic electrons and pairs lose energy by Compton scattering of the thermal radiation (which is roughly isotropic in the comoving frame). Synchrotron losses might dominate for high-γ electrons, but for those with modest γ, synchrotron emission is inhibited by self-absorption; these lose their energy primarily by Compton scattering even if the magnetic energy density exceeds that of the thermal radiation. They will cool down and thermalize in a time that is short compared to the dynamic time.

Relativistic electrons moving through blackbody radiation Compton-boost each scattered photon by γ^2, producing a power law rather than just boosting each photon by a small amount. However, if the slope of the injected power law is steeper than -2, most of the energy will be at the low-energy end, and all the energy of electrons with, say, $\gamma \lesssim 3$ would go into what would look like a broad thermal peak. They would emit no synchrotron radiation (because of self-absorption), and they would not boost any of the thermal photons by more than a factor $\epsilon_c \sim 10$.

If the primary dissipation were mainly by strong shocks, most of the energy might be channeled initially into very high gamma electrons, which would produce photons with a power-law spectrum extending to very high energies; production of photons above 1 MeV in the comoving frame would only require $\gamma \sim 10$–30 for Compton scattering, and little more than 10^3 for synchrotron emission. However, pair production can change this situation, leading again to a situation in which energy is ultimately dissipated via thermal Comptonization. If the compactness parameter is more than unity (which, as shown below, is often the case for the usual parameters considered), most of the energy in photons with >1 MeV energies in the comoving frame are converted into pairs with very modest γ. These pairs then lose their energy (as described above) by Compton cooling, resulting in a quasi-thermal spectrum, whose characteristic peak is a factor $\epsilon_c \lesssim 10$ above the original thermal peak, i.e., in the tens to hundreds of keV. These pairs effectively establish a new photosphere outside the one that would have been present in their absence, and the dissipation (or shocks) responsible for these pairs will effectively be a subphotospheric dissipation, which has different characteristics from the more familiar shocks that occur well outside the photosphere (e.g., Ghisellini & Celotti 1999; Kobayashi et al. 2002; Pe'er & Waxman 2004).

For a GRB outflow of radiative luminosity L and bulk Lorentz factor Γ in the observer frame, at a radius r, the comoving scattering opacity due to $e^\pm$ pairs in the high-comoving compactness regime is

$$\tau'_\pm \sim l'^{1/2} \sim \left(L\sigma_T/4\pi m_e c^3 \Gamma^3 r\right)^{1/2}, \quad (1)$$

where l' is the comoving frame compactness parameter (e.g., Pe'er & Waxman 2004). Here we have approximated $L(>1 \text{ MeV}) \sim \epsilon_d L_0$, where L_0 is the total luminosity in the observer frame, and we have taken other efficiency factors to be of order unity. The functional dependence of equation (1) can be obtained by considering in the comoving frame (primed quantities, as opposed to unprimed quantities in the observer frame) the balance between the rate at which pairs annihilate and the rate at which pairs are formed. The latter is the rate at which photons capable of pair-producing are introduced into the flow, i.e., the photon density above $m_e c^2$ divided by the comoving dynamic time, $n'^2_\pm \sigma_T c \sim (L/4\pi r^2 m_e c^3 \Gamma^2)(c\Gamma/r)$, from which follows the pair optical depth $\tau'_\pm \sim n'_\pm \sigma_T (r/\Gamma)$. The pair photosphere $r_{\text{ph},\pm}$ is the radius where $\tau'_\pm \sim 1$, or

$$r_{\text{ph},\pm} \sim (\epsilon_d/2\alpha)\left(m_p/m_e\right)(L_0/L_E)\Gamma^{-3}r_0$$
$$\sim 2\times10^{14}L_{51}\epsilon_{d,-1}\alpha^{-1}\Gamma_2^{-3} \text{ cm}, \quad (2)$$

where $L_E = 4\pi GM m_p c/\sigma_T \simeq 1.25\times10^{39}m_1$ ergs s^{-1} is the Eddington luminosity; $r_0 = \alpha r_g$, where $\alpha \geq 1$, and $r_g = 2GM/c^2 \simeq 3\times10^6 m_1$ cm is the Schwarzschild radius for a central object (e.g., black hole) of mass $M \sim 10 m_1$ solar masses; and $\eta = L/(\dot{M}c^2)$ is the dimensionless entropy of the relativistic outflow.

On the other hand, the scattering opacity due to the ordinary electrons associated with baryons in the flow would give rise to a "baryonic photosphere" at

$$r_{\text{ph},b} \sim (1/2\alpha)(L_0/L_E)\eta^{-1}\Gamma^{-2}r_0$$
$$\sim 1.2\times10^{12}L_{51}\alpha^{-1}\eta_2^{-1}\Gamma_2^{-2} \text{ cm}, \quad (3)$$

with the same notation as above.

For an outflow that starts at $r_0 = \alpha r_g = 3\times10^6\alpha m_1$ cm, where $\alpha \geq 1$ and the initial Lorentz factor $\Gamma_0 \sim 1$, under adiabatic conditions energy-momentum conservation leads to a Lorentz factor that grows linearly as $\Gamma(r) \propto r/r_0$ until it reaches a saturation radius, $r_s \simeq r_0\eta \simeq 3\times10^8\alpha m_1\eta_2$ cm, beyond which the Lorentz factor saturates to $\Gamma \simeq \eta = $ constant. One can then define two critical limiting Lorentz factors,

$$\eta_b = \left(\frac{1}{2\alpha}\frac{L}{L_E}\right)^{1/4} = 7.9\times10^2(L_{51}/\alpha m_1)^{1/4}, \quad (4)$$

$$\eta_\pm = \left(\epsilon_d\frac{m_p}{m_e}\right)^{1/4}\eta_b = 2.9\times10^3\left(L_{51}\epsilon_{d,-1}/\alpha m_1\right)^{1/4}, \quad (5)$$

which characterize the behavior of the baryon and pair photospheres below and above the saturation radius. The pair photosphere behaves as

$$\frac{r_{\text{ph},\pm}}{r_0} = \begin{cases} \eta_\pm & \text{for } r < r_s \\ \eta_\pm(\eta/\eta_\pm)^{-3} & \text{for } r > r_s, \end{cases} \quad (6)$$
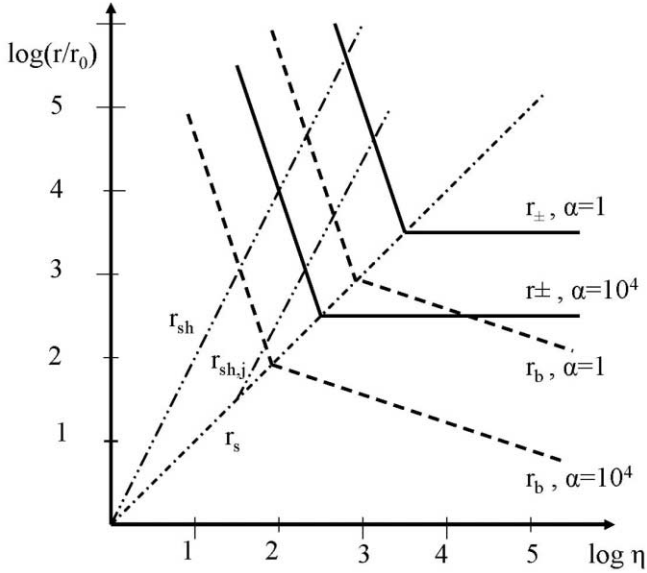
FIG. 1.—Radii of the pair (*solid lines*) and baryon (*dashed lines*) photospheres as a function of η for $\epsilon_d = 10^{-1}$, $L_0 = 10^{51}$ ergs s^{-1}, $\alpha = 1$ ($r_0 = 3 \times 10^6 m_1$ cm), and $\alpha = 10^4$ ($r_0 = 3 \times 10^{10} m_1$ cm). Also shown are the saturation radius $r_s = r_0 \eta$ and the spherical minimum shock radius $r_{sh} = r_0 \eta^2$. Instabilities at the nozzle θ of a jet could lead to shocks at a lower minimum radius, $r_{sh,j}$, while magnetic dissipation could in principle occur both above and below r_s.

FIG. 2.—Kinetic luminosity and photospheric radiation luminosity as a function of radius. Beyond the saturation radius the luminosity decays as $L_\gamma \propto r^{-2/3}$, but beyond $\epsilon_d^{3/2} r_s$ the fraction ϵ_d of the kinetic energy reconverted into radiative (pair) form becomes significant. Also shown is the value of the observer temperature. Comptonization at the pair photosphere (see text) could boost this by an additional factor $\lesssim 10$.

and the pair photosphere occurs at $r < r_s$ for $\eta > \eta_\pm$. The baryon photosphere behaves as

$$r_{ph,b}/r_0 = \begin{cases} \eta_b (\eta/\eta_b)^{-1/3} & \text{for } r < r_s \\ \eta_b (\eta/\eta_b)^{-3} & \text{for } r > r_s, \end{cases} \quad (7)$$

and the baryon photosphere occurs at $r < r_s$ for $\eta > \eta_b$. This is shown schematically in Figure 1 for values of $\alpha = 1$ and 10^4, i.e., initial radii $r_0 = 3 \times 10^6 \alpha_0 m_1$ and $3 \times 10^{10} \alpha_4 m_1$ cm.

The pair photosphere (eq. [2]) will be above the baryon photosphere provided that

$$\epsilon_d > (m_e/m_p), \quad (8)$$

where ϵ_d characterizes the dissipation efficiency producing photons above energies $m_e c^2$ in the comoving frame.

3. CHARACTERISTIC PHOTON LUMINOSITIES AND TEMPERATURES

When the conditions of equation (8) are satisfied, the real (outermost) photosphere is not the baryon photosphere but the pair photosphere. The pair photosphere will have a luminosity $L_{\gamma\pm} = \epsilon_d L_0 \leq L_0$. At $r < r_s$, the observed radiation is insensitive to the actual details of the photosphere: the decrease with r in the comoving-frame luminosity is compensated by the observer-frame boost given by the increasing Lorentz factor Γ; moreover, there is less scope for dissipation (except in the case when Poynting flux far exceeds the radiative flux in the jet).

On the other hand, for a photosphere at $r > r_s$ the luminosity decays as $L_\gamma = L_0 (r/r_s)^{-2/3}$ in the adiabatic regime. However, if dissipation occurs above r_s, this leads to an effective luminosity

$$L_\gamma \sim \epsilon_d L_0. \quad (9)$$

This luminosity is achieved at the baryon photosphere if dissipation occurs below this radius, even in the absence of sig-
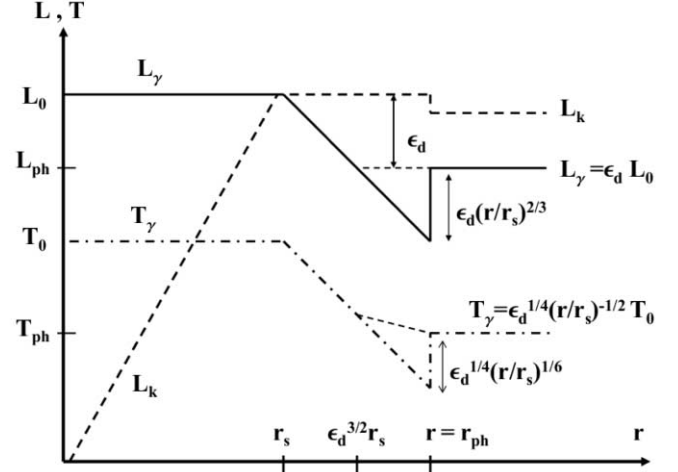
nificant pair formation, or above the baryon photosphere if dissipation above the baryon photosphere leads to a pair photosphere radius $r_\pm$ such that $\epsilon_d (r_\pm/r_s)^{2/3} \geq 1$ (see Fig. 2). In such cases the effective photosphere luminosity exceeds what would have emerged from a nondissipative outflow by $\epsilon_d (r/r_s)^{2/3}$.

The characteristic initial temperature of the fireball outflow is

$$T_0 = (L_0/4\pi r_0^2 ca)^{1/4} = 1.2 L_{51}^{1/4} (\alpha m_1)^{-1/2} \text{ MeV}, \quad (10)$$

which for a larger $\alpha = 10^4$ (i.e., for a larger initial radius $r_0 = 3 \times 10^{10} \alpha_4 m_1$ cm) would be $T_0 = 12.1 L_{51}^{1/4} (\alpha_4 m_1)^{-1/2}$ keV.

For $r < r_s$ the observer-frame effective photospheric temperature (even in the presence of dissipation) remains as $T_\gamma = T_0$, being boosted by the growing Lorentz factor back to its initial value. For $r > r_s$, adiabatic effects (in the absence of dissipation) cause the temperature to fall off as $T_\gamma = T_0 (r_s/r)^{-2/3}$. Since, however, dissipation leads to a luminosity $\epsilon_d L_0$ that can exceed the adiabatic value, this results in a temperature T_γ that drops $\propto r^{-1/2}$. If dissipation is maintained all the way to the (baryonic or pair-dominated) photosphere, the temperature is

$$T_{\gamma,d} = \epsilon_c \epsilon_d^{1/4} (r/r_s)^{-1/2} T_0, \quad (11)$$

where a factor $\epsilon_c \gtrsim 1$ accounts for possible departures from a blackbody. This temperature is larger by a factor $\epsilon_c \epsilon_d^{1/4} (r/r_s)^{1/6}$ than the adiabatic photosphere temperature T_γ (Fig. 2).

Thus, one consequence of dissipation is that, even for $\alpha = 1$, e.g., with $\epsilon_d = 10^{-1}$ and $r_d/r_s = 10^2$, the characteristic temperatures can be $kT_\gamma \sim 60$ keV, i.e., peak photon energies in the X-ray flash range, while for $\alpha = 10^4$ this energy can easily be as low as a few keV. If dissipation is important only for some range of radii, starting at $r_s \epsilon_d^{3/2}$ but ceasing, say, at some radius r_c below the photosphere, r_{ph}, then the adiabatic decay $L_\gamma \propto T_\gamma \propto r^{-2/3}$ resumes above r_c until r_{ph}, so the photospheric luminosity would be lower than implied by equations (9) and (11).

We should note that, in the present context, r_0 is essentially the radius beyond which the Lorentz factor starts to grow as $\Gamma \propto r/r_0$. Thus any dissipation in the inner "cauldron" or along the inner jet effectively pushes out r_0. This could come about

because of entrainment, or because of the oblique shocks that occur when the jet is initially poorly collimated (as exemplified in numerical calculations of collapsar models, such as in Zhang & Woosley [2004], where in effect $r_0 \gtrsim 10^8$ cm). The initial reference temperature T_0 is correspondingly lower. (Another effect that could change the reference temperature is if the inner jet is Poynting dominated, so that only a small fraction of the flux is in the radiation. In this case, pairs can be even more dominant inside r_s. There could be modifications to the outflow dynamics if the field were tangled and did not obey the straightforward Bernoulli equation for a relativistic gas [cf. Heinz & Begelman 2000].)

Dissipation need not necessarily lead to pair formation. For example, in a "slow heating" scenario (such as that in Ghisellini & Celotti 1999), the accelerated particles, and the photons associated with them, could all have energies substantially below ~0.5 MeV. Dissipation would then not enhance the photospheric radius, but, even so, as indicated above, the characteristic photon energies and photospheric luminosity could be substantially boosted over what their adiabatic values would have been.

An important feature of the model is that millisecond variations, either at the photosphere or due to internal shocks farther out, may still be traced back to irregularities in the jet boundary at r_0, since the characteristic timescale for a nozzle of opening half-angle θ_j is $t_{var} \sim r_0\theta_j/c$, rather than r_0/c itself, which can be less than 1 ms even if r_0 is of the order 10^8 cm. If internal shocks are to develop, they must be induced by unsteady conditions near the base of the jet (resulting in changes in η and the saturation Lorentz factor). While for the usual minimum variability timescale $t_{var} \sim r_0/c$ shocks would develop above the line marked r_{sh} in Figure 1, for $t_{var} \sim r_0\theta/c$ the shocks can form at radii $r_{sh,j}$, a factor θ_j smaller than that for the spherical case; see Figure 1. Dissipation at such or smaller radii is also possible, e.g., in the case of oblique shocks induced by irregularities in the walls of the jet, during the collimation of an initially poorly collimated jet, or in the case of dissipation due to magnetic reconnection.

Note also that any variability at r_0 would alter the conditions at the photosphere (and the value of the photospheric radius). Moreover, the photospheric changes can be rapid. Obviously this is true if the photosphere lies below the saturation radius; however, this condition is not necessary, and provided that the photospheric radius is within $\eta^2 r_0$, there is no smearing of variability on any timescale down to r_0/c. We would therefore, generically, expect an internal shock to be slightly preceded by a change in the luminosity of the thermal component (and in E_{pk}). Indeed, one is led to conjecture that rapid variations in the photosphere could be at least as important as the associated internal shocks in causing rapid variability in GRBs. In contrast to shocks, variations in the photospheric emission could as readily account for a short dip as for a short peak. Detailed evidence of spectral softening during both the rise and fall of individual subpulses (cf. Ryde 2004) could clarify the relative contributions of these effects.

4. SPECTRUM FORMATION

When dissipation occurs, one expects the photospheric spectrum to be "gray" rather than an accurate blackbody, because there would not (except near the base of the jet) be processes capable of producing the new photons appropriate to a blackbody with the enhanced energy density. All photons emerging from the photosphere, however, will have undergone multiple scatterings. In the case of shock dissipation, a power-law relativistic electron energy distribution may be formed, which would upscatter the thermal photons into a power-law photon distri-
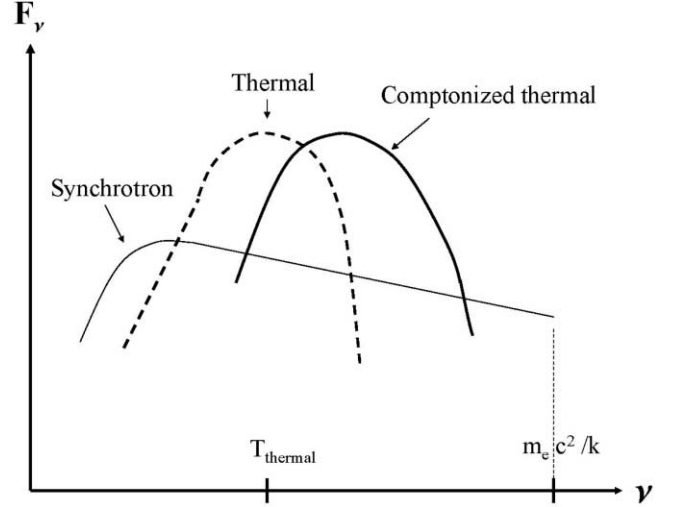

Fig. 3.—Schematic comoving frame spectrum showing the photospheric (thermal) spectrum and its Comptonized component, as well as a shock synchrotron component (assumed to arise farther out). This is the generic spectrum characterizing a slow dissipation model (see text). Shocks with pair formation could lead to an additional component at higher energies.

bution whose index is similar to that of synchrotron radiation, with pair formation being possible at comoving energies $\gtrsim m_e c^2$. At the pair photosphere the comoving inverse Compton cooling time is $t'_{IC} \sim 4 \times 10^{-3} L_{51}\epsilon^2_{d,-1}\alpha^{-2}\gamma^{-1}_{e,3}$ s, while the dynamic time is $t'_{dyn} \sim 3 \times 10^1 L_{51}\epsilon_{d,-1}\alpha^{-1}\eta^{-4}_2$ s. The interplay between the electron and photon distributions requires a detailed analysis and is discussed in A. Pe'er et al. (2005, in preparation). For a slow heating scenario, such as that of Ghisellini & Celotti (1999) but with the added feature of dissipation (e.g., from magnetic reconnection or from multiple shocks and/or magnetohydrodynamic [MHD] turbulence behind them), one expects the electrons to be heated to more modest values, say $\gamma_e \lesssim$ few, but the electrons might get reheated every Compton cooling time. In this case pair formation is at best modest (A. Pe'er et al. 2005, in preparation), so the effects outside the baryonic photosphere are not significant. If slow dissipation occurs at or below the baryonic photosphere, where pair formation is suppressed, the inverse Compton (IC) cooling time is $t'_{IC} \sim 10^{-4} L_{51}\alpha^{-1}\gamma^{-1}_{e,0.5}\eta^{-4}_2$ s, while the dynamic time is $t'_{dyn} \sim 2 \times 10^{-1} L_{51}\alpha^{-1}\eta^{-4}_2$ s. The dissipative baryonic photosphere thermal peak is at $3kT_{\gamma,b} \sim 20L^{-1/2}_{51}\epsilon^{1/4}_{d,-1}\alpha^{1/2}\eta^2_2$ keV, which (depending on γ_e) may get upscattered by factors $\sim 1-10$. The schematic shape of the spectrum is given in Figure 3 (cf. Pe'er et al. 2005; A. Pe'er et al. 2005, in preparation), showing the original quasi-thermal Wien component, the upscattered photospheric component resulting from subphotospheric dissipation and Comptonization, and a possible additional synchrotron component from shocks outside the photosphere. The peak frequency scales with the amount of dissipation according to a power law that depends on how many new photons are produced. (The photon production depends on the radial dependence of the dissipation and on the detailed dissipation mechanism.)

The dependence of the spectral peak energy on the burst parameters, as observed in a given energy range by a given instrument, depends on the specific mechanism responsible for the spectrum in that energy range. In the Burst and Transient Source Experiment (BATSE) and *BeppoSAX* energy range (roughly 20 keV to 0.5 MeV), there is a quantitative relationship observed between the spectral peak energy E_{pk} and the isotropic-equivalent luminosity of the burst in that energy range, L_{iso} (which requires a knowledge of the redshift of the burst). This

relationship (Amati et al. 2002) is $E_{pk} \propto L_{iso}^{1/2}$ for a score of bursts with redshifts. Our generic assumptions naturally yield a peak in the relevant range, but cannot predict any correlation with L without a more specific model that relates L to the other significant parameters, in particular r_0 and η. Without going into details, we may consider several possibilities.

If one seeks to explain this relationship by interpreting the peak energy as the synchrotron peak in a simple standard internal shock scenario, one expects the dependence (e.g., Zhang & Mészáros 2002)

$$E_{pk} \propto \Gamma^{-2} t_{var}^{-1} L^{1/2}, \qquad (12)$$

where $L_{\gamma, iso} \sim L$. Here Γ and t_{var} are the Lorentz factor of the outflow and its typical variability timescale, respectively. If the latter two quantities are approximately the same for all bursts, this would reproduce the Amati et al. (2002) relation. However, it is not obvious why there should be a constancy of Γ and t_{var} across bursts, even if approximate.

If the spectral peak is of a quasi-thermal origin determined by the photosphere (possibly shifted up by Comptonization, e.g., from pair dissipative effects such as discussed above), and if there are enough photons to guarantee an approximate blackbody distribution, the peak photon energy in the observer frame is, using equation (2),

$$E_{pk} \propto \Gamma k T_{pk}' \propto \Gamma \left(L/\Gamma r^2\right)^{1/4} \propto \Gamma^2 L^{-1/4} \propto L^{(8\beta-1)/4}, \quad (13)$$

which depends mainly on the Lorentz factor Γ. If the latter in turn depends on L, e.g., as $\Gamma \propto L^{\beta}$, one obtains the last part of equation (13). For instance, taking the observed Frail et al. (2001) relation $L_{\gamma, iso} \propto \theta^{-2}$ between the jet opening half-angle θ inferred from the light-curve break, and using the causality relation $\theta \sim \Gamma^{-1}$, equation (13) becomes $E_{pk} \propto L^{3/4}$.

If dissipation occurs mainly very close to the central engine, this could result in a larger radius r_0, where r_0 is defined as the radius beyond which $\Gamma \propto r/r_0$. Assuming that the "drag" or dissipation at the base increases r_0 according to, e.g., $r_0 \propto L^{-\beta'}$ for a photosphere occurring inside the saturation radius $r_0 < r_{ph,\pm} < r_s$, the growth of the Lorentz factor $\Gamma \propto r/r_0$ cancels out the adiabatic drop $T' \propto r^{-1}$ of the comoving temperature, and one has

$$E_{pk} \propto r_0^{-1/2} L^{1/4} \propto L^{(2\beta'+1)/4}. \qquad (14)$$

Hence, for $\beta' = (0.5, 1)$ one has $E_{pk} \propto (L^{1/2}, L^{3/4})$.

In the extreme "photon starved" case (likely to apply if the dissipation is concentrated not far inside the photosphere, where the photon number N_γ is constant), one would have $E_{pk} \propto L/N_\gamma \propto L$.

Thus, a variety of E_{pk} versus L dependences might in principle be expected, depending on the uncertain physical conditions just below the photosphere, some of which approximate the reported $L^{1/2}$ behavior. Ghirlanda et al. (2004) have recently claimed an empirical correlation between E_{pk} and a different quantity, the angle-corrected total energy $E_{tot} = E_{iso}(1 - \cos\theta_j) \sim E_{iso}(\theta_j^2/2)$, where $E_{iso} \simeq L_{iso} t_\gamma$, and t_γ is the burst duration. They find a tighter correlation between E_{pk} and E_{tot} than between E_{pk} and E_{iso} for bursts with observed redshifts and breaks. Furthermore, in contrast to the Amati et al. (2002) $E_{pk} \propto E_{iso}^{1/2}$ dependence, they deduce from the data a steeper slope, $E_{pk} \propto E_{tot}^{0.7}$. Taking a standard burst duration and jet opening angle, this is of

the form discussed in equations (13) and (14). A critique of the methods for obtaining both types of correlations from the data is given by Friedman & Bloom (2005). We should note that such correlations are generally derived assuming that the efficiency of gamma-ray production is the same for all bursts, independently of the luminosity or the total energy. If, however, the efficiency were lower for the weaker (and therefore softer) bursts, then the correlation would have a flatter E_{pk} versus E_{tot} slope than that currently derived from the data. This is because, for a given gamma-ray isotropic luminosity, the momentum outflow per unit solid angle would be higher than they assume. This means that the standard jet-break argument would imply a narrower beam than that inferred under the constant efficiency assumption, and therefore a lower L_{tot} (for a given E_{pk}) than the values currently derived.

5. DISCUSSION

We have considered dissipative effects below the photosphere of GRB or X-ray flash (XRF) outflows, such as, e.g., those due to magnetic reconnection or shocks. Such dissipation can lead to copious pair formation dominating the photospheric opacity. Alternatively, if dissipation occurs not too far above an initial photosphere, it can result in a second effective photosphere situated outside the initial one.

Subphotospheric dissipation can increase the radiative efficiency of the outflow, significantly boosting the quasi-thermal photospheric component so that it may well dominate the much discussed synchrotron component from nonthermal shocks outside the photosphere. The hypothesis that GRB emission is dominated by a Comptonized thermal component offers a natural explanation for the thermal GRB spectra discussed most recently, e.g., by Ryde (2004). It can also naturally explain the steeper than synchrotron lower energy spectral indices (Preece et al. 2000; Lloyd et al. 2000) noticed in some bursts.

The quasi-thermal peak of the photospheric spectrum is controlled by the total luminosity L_0 and by the reference injection radius r_0, above which the Lorentz factor starts to grow linearly. Dissipation near the central object or along the inner jet can result in an increase in r_0, thus lowering the reference temperature of the outflow that characterizes the quasi-thermal photospheric component. The characteristic variability timescales $r_0 \sin\theta/c$ for jets with the observationally inferred opening half-angles θ are in the millisecond range. The spectral peak of the dissipation-enhanced photospheric component, upscattered in energy by factors of ~10 due to electrons accelerated in the dissipation process, results in typical photon energies ranging between those of classical bursts and X-ray flashes.

The relationship between the observed photon peak energy and the luminosity can have a variety of functional forms, which depends on a number of as yet poorly determined parameters. However, plausible assumptions can lead to relationships of the type $E_{pk} \propto L_{iso}^{1/2}$ (Amati et al. 2002), or $E_{pk} \propto E_{tot}^{0.7}$ (Ghirlanda et al. 2004). Although more physics and a more specific model are needed before we can explain the correlations, the idea that E_{pk} is essentially a thermal peak seems better able to account for a "standardized" value in a given class of objects, because there is not a steep Γ-dependence (and indeed to the first order the Γ factor cancels out, because adiabatic cooling in the comoving frame is compensated by the Doppler boosting).

In summary, our main result is that a spectral peak at photon energies in the range of tens to hundreds of keV, typical of XRFs and GRBs, can naturally arise from an outflowing jet, in which dissipation below a baryonic or pair-dominated photosphere enhances the radiative efficiency and gives rise to a Comptonized

thermal spectrum. On this hypothesis, the recently discovered correlations between L and E_{pk} should be an important diagnostic of how the key jet parameters—physics near the "sonic point," baryon contamination, etc.—depend on L.

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