Compton dragged gamma-ray bursts: the spectrum

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Abstract
We calculate the spectrum resulting from the interaction of a fireball with ambient soft photons. These photons are assumed to be produced by the walls of a funnel in a massive star. By parametrizing the radial dependence of the funnel temperature we calculate the deceleration of the fireball self-consistently, taking into account the absorption of high energy γ-rays arising from interactions with the softer ambient photons. The resulting spectrum is peaked at energies that are in agreement with observations, has a ν2 slope in the X-ray band and a steep power law high energy tail.

Key words: radiation mechanisms: non-thermal – gamma-rays: bursts – X-rays: general.

1 Introduction
We have recently proposed (Lazzati et al. 2000, hereafter Paper I) that the gamma-ray-burst (GRB) phenomenon originates from the interaction of a relativistic fireball with a dense photon environment, leading to Compton drag. On one hand, this is an inevitable effect if the progenitors of GRBs are massive stars that are about to explode or have just exploded as supernovae; on the other hand, this mechanism greatly alleviates the efficiency problem faced by the standard internal shock scenario (Lazzati, Ghisellini & Celotti 1999; Kumar 1999; Panaitescu, Spada & Meszaros 1999). In Paper I we have discussed the basic Compton-drag scenario, showing the way in which this process can convert bulk motion energy directly into radiation with a remarkably high efficiency and, on the basis of simple estimates, the way in which the resulting spectrum should peak, in a ν2 representation, around ∼1 MeV, as is observed.

Here we quantitatively and self-consistently estimate the predicted spectrum, assuming that the fireball propagates in a funnel inside a massive star, and show that, independent of the details of the model, it satisfactorily resembles what is observed. As the funnel walls emit a blackbody spectrum, and the scattered photons are boosted by the square of the Lorentz factor (Γ) of the fireball, the local spectrum has a blackbody shape at a temperature enhanced by Γ2. However, the observed spectrum, which is a convolution of all the locally emitted spectra, is not one of a blackbody as a consequence of four main effects:

(i) the funnel walls would not be at a uniform temperature, but there should be a gradient between the internal and external parts;
(ii) if the Compton-drag process is efficient, the fireball decelerates;
(iii) the very high energy emission produced in the internal regions can interact with the ambient photons producing electron–positron pairs;
(iv) the fireball may become optically thin to scattering outside the funnel where the ambient photons are characterized by the same temperature, but where their energy density is progressively diluted with distance.

2 Basic Assumptions
We postulate that the fireball propagates with a bulk Lorentz factor Γ inside a funnel cavity, the walls of which emit blackbody radiation at a temperature T, of conical shape with a semi-aperture angle ψ. The calculation starts at the distance z0, which is assumed to be the end of the acceleration phase and, for consistency, we verify that the power emitted at z < z0 is negligible.

We assume that the fireball is and remains cold in the comoving frame. At z0, in fact, the internal energy has been already used to accelerate the fireball, and thus the protons are subrelativistic. On the other hand, leptons might still be hot and/or being re-heated at z0 when the bulk scattering process starts to be efficient. However, in a few (Compton) cooling time-scales the leptons would reach the (subrelativistic) Compton temperature. It is thus reasonable to also treat the leptonic component as cold when estimating both the dynamics and the resulting spectrum.

The initial Lorentz factor (at z0) is indicated as Γ0 and the fireball energy is therefore Ef = Γ0Mc2, where Mc is the rest mass of the fireball. For simplicity, the dependence of the temperature on z, between z0 and the radius of the star zc, has been parametrized by a power law:

\[ T(z) = T_0 \left( \frac{z}{z_0} \right)^{-b} = T_0 \left( \frac{z}{zc} \right)^{-b}, \]

where T0 is the temperature at the top of the funnel.
Inside the funnel, we approximate the local radiation energy density of the ambient photons as $U(z) = aT^4(z)$. Beyond $z_e$, and in the region where the fireball remains optically thick (i.e. for $z < z_T$, see below), $U(z)$ is characterized by the same temperature, but decreases. As the relevant quantity is the amount of radiation that is indeed scattered by the fireball, we parametrize the dependence on $z$ of the product $U(z) \times (\text{the scattering rate})$ as $(z/z_e)^{-\delta}$.

We consider $g$ a free parameter. Inside the funnel $g = 0$, whereas outside the funnel a value of $g > 2$ can account for a decrease in the scattering rate caused by the changing of the typical scattering angle (photons come preferentially at smaller angles as $z$ increases). As the scattering rate is $\propto (1 - \beta \cos \theta)$, where $\theta$ is the angle between the photon and the electron directions, far from the star surface $(1 - \beta \cos \theta) \propto (z/z_e)^{-2}$, corresponding to $g \approx 4$. Furthermore, some of the radiation produced by the massive star could be reflected and re-isotropized by the scattering material, of unknown radial density profile, which is likely to surround the massive star progenitor. In particular, if this forms a wind with a $z^{-2}$ profile, the energy density of the re-isotropized radiation scales as $z^{-3}$ and dominates the seed photon distribution at large distances. In this case $U(z) \times (\text{the scattering rate})$ can have a complex profile: being flat in the vicinity of the surface of the star, then decreasing as $z^{-2}$ and as $z^{-4}$ for increasing $z$ and becoming flatter when the component associated with the re-isotropized photons dominates. It is also possible, that as a result of intermittent stellar activity, the stellar wind is not continuous. In this case a single shell may dominate the scattering, producing a homogeneous and isotropic scattered radiation field that dominates the total radiation energy density beyond some critical distance.

The distance $z_T$ at which the fireball becomes optically thin to scattering is

$$z_T = \left( \frac{\sigma_T E_i}{\pi \psi_m c^2 x_{\psi}} \right)^{1/2} \sim 3.7 \times 10^{14} \psi^{-1} a_{E(1/1)}^{1/2} \Gamma_{0.2}^{1/2} \text{cm},$$

where the conventional representation $Q = Q_0 10^6$ and cgs units are adopted. It then becomes likely that the fireball becomes transparent at $z > z_e$ (because the radius of red supergiants is $z_e \leq 10^{13}$ cm).

As long as the fireball is opaque to scattering, the interaction with photons boosts their energy by a factor of $\sim 2\Gamma^2$. Therefore, the (local) total energy emitted by the fireball through the Compton-drag process (over a distance $dz$) is

$$dE(z) = 2\pi \psi^2 z^2 aT_0^2 \left( \frac{z}{z_0} \right)^{-g} \Gamma^2 dz \quad \text{for} \quad z < z_e,$$

and

$$dE(z) = 2\pi \psi^2 z^2 aT_0^2 \left( \frac{z}{z_e} \right)^{-g} \Gamma^2 dz \quad \text{for} \quad z > z_e.$$  

The factor 2 in front of the right-hand side of these equations takes into account that the preferred scattering angle is $\sim 90^\circ$, corresponding to an average energy boost of $2\Gamma^2$.

Let us now consider the spectral shape. For this it is convenient to use dimensionless photon energies and temperatures, defined as $x = h\nu/(mc^2)$ and $\Theta = kT/(mc^2)$, respectively.

The resulting Compton spectrum has a blackbody shape, of effective temperature $T_{eff} = 2\Gamma^2 T$ (or $\Theta = 2\Gamma^2 \Theta$), i.e. the local spectral distribution produced within $dz$ is given by:

$$dE(z, x) = \frac{\pi^2 \psi^2 x^2}{\Gamma^2} \frac{\left(m_e c\right)^3}{h^3} \left( \frac{1}{\exp(x/\Theta) - 1} \right) dz$$

for $z < z_e$  

and

$$dE(z, x) = \frac{\pi^2 \psi^2 x^2}{\Gamma^2} \frac{\left(m_e c\right)^3}{h^3} \left( \frac{x^3}{\exp(x/\Theta) - 1} \right) dz$$

for $z > z_e,$

where $\Theta_{e, s} = 2T^2 \Theta_s$. Equations (5) and (6) are correctly normalized, i.e. the integrated energies correspond to the energies expressed in (3) and (4).

### 3 THE FIREBALL DYNAMICS

As long as the fireball remains optically thick for scattering, and that this occurs in the Thomson regime, the dynamics (deceleration) of the fireball as a result of radiative drag obeys

$$M_{\text{f}} c^2 \frac{d\Gamma}{dz} = -2\pi \psi^2 z^2 aT^4 \Gamma^2.$$  

Assuming the temperature profile of equation (1) we obtain

$$\Gamma = \frac{\Gamma_0}{\left[ 1 + 2\pi \psi^2 aT_0^4 \left( \frac{z}{z_0} \right)^{2-g} \right]} \quad \text{for} \quad z_0 < z < z_e,$$

and thus the deceleration radius $z_d$, defined as the distance at which $\Gamma$ is halved, corresponds to

$$z_d = z_0 \left[ 1 + \frac{E_0(3 - 4b)}{2\pi \psi^2 aT_0^4 \left( \frac{z}{z_0} \right)^{2-g}} \right]^{1/(3-4b)}.$$  

Beyond $z_d$, the Lorentz factor decreases with distance as a power law, the slope of which is determined by the temperature profile.

Outside the star radius ($z > z_e$) the Lorentz factor follows

$$\Gamma = \frac{\Gamma_0}{\left[ 1 + \frac{E_0(3 - 4b)}{2\pi \psi^2 aT_0^4 \Gamma_{0.2} \left( \frac{z}{z_0} \right)^{2-g}} \right]} \quad \text{for} \quad z_e < z < z_T.$$  

Note that Klein–Nishina effects are important for incoming photon energies such that $\Gamma > 1$, i.e. when $\Theta > 1/(3\Gamma)$. For simplicity we neglect Klein–Nishina effects when calculating $\Gamma(z)$, but we assume that no scattering events occur when $\Theta > 1/(3\Gamma)$ in the calculation of the spectrum. This simplification is justified as long as most of the fireball energy is lost in the Thomson scattering regime (see Fig. 2, which shows that $\Gamma$ starts to decrease at distances where the temperature is small enough to ensure scattering entirely in the Thomson regime).

When the fireball becomes optically thin, the number of scattered photons is correspondingly reduced, and the process becomes less efficient. As shown by equation (2), this is likely to happen at some distance from the star surface, where the photon density is also reduced, thus further decreasing the efficiency of the process. In the numerical calculations we have, however, included the optically thin scattering regime, and one can see its contribution in Fig. 1 (dotted line).
Let us therefore estimate the photon–photon optical depth, $\tau_{\gamma\gamma}$, by integrating the product of the photon–photon cross-section $\sigma_{\gamma\gamma}(x, x_T)$ and the photon density above the threshold $n_g(z)$ over the gamma-ray path, i.e. from the site of creation $z_1$ to infinity, and over the photon energies:

$$\tau_{\gamma\gamma}(z_1, x) = \int_{x_T}^{\infty} \int_{z_1}^{\infty} \sigma_{\gamma\gamma}(x', x)n_g(z, x') \, dz \, dx'$$  \hspace{1cm} (11)

As $\sigma_{\gamma\gamma}(x', x)$ is peaked at the threshold energy, equation (11) can be simplified (Svensson 1984, 1987) to

$$\tau_{\gamma\gamma}(z_1, x) = \frac{\sigma_T}{5n_e\epsilon^2} \int_{x_T}^{\infty} x_T U(z, x_T) \, dz,$$

where $U(z, x_T) = n_g(x_T)GQ\hat{z}x_T$ is the photon energy density at threshold, at the location $z$, i.e.

$$U(z_1, x_T) = \frac{8\pi h}{c^3} \left( \frac{m_e^2 \epsilon}{h} \right)^4 \hat{x} \frac{1}{\exp[x_T/(\Theta z_1)] - 1}.$$  \hspace{1cm} (13)

The radiation flux produced at the location $z_1$ is then decreased by the factor $\exp[-\tau_{\gamma\gamma}(z_1, x_T)]$ while crossing the funnel.

The absorbed radiation will be reprocessed by the pairs and redistributed in energy. Each electron and positron will have an energy $\gamma \sim x/2$ at birth, and will cool as a consequence of the Compton-drag process. The positrons will then annihilate in collisions with the electrons in the fireball, producing a Doppler blueshifted annihilation line at $x \sim \Gamma$. We have neglected these reprocessing mechanisms, because, as can be seen in Fig. 1, the amount of energy absorbed in $\gamma-\gamma$ collisions is small, amounting to a few per cent at most.

5 THE SPECTRUM

The observed total spectrum can be computed by integrating equations (3) and (4) over $z$, taking into account photon–photon absorption. The contribution produced within the star is given by

$$E(x) = \pi^2 \sigma_T \epsilon^2 m_e^2 \left( \frac{m_e \epsilon}{h} \right)^3 \int_{0}^{x_T} \frac{z^2 \hat{x}}{\epsilon^2 \exp[x/(\Theta z)] - 1} \, dz$$

for $z_0 < z < z_*$.  \hspace{1cm} (14)

whereas beyond $z_*$ the number of target photons able to interact with high-energy gamma-rays to produce pairs is negligible, and thus, ignoring photon–photon absorption, we obtain

$$E(x) = \pi^2 \sigma_T \epsilon^2 m_e^2 \left( \frac{m_e \epsilon}{h} \right)^3 \int_{0}^{x_T} \frac{z^2 (z/z_*)^{-\gamma} \hat{x}}{\epsilon^2 \exp[x/(\Theta z)] - 1} \, dz$$

for $z_* < z < z_T$.  \hspace{1cm} (15)

In Fig. 1 we show three examples of the predicted spectrum corresponding to different values of the initial bulk Lorentz factors. To illustrate the main features of the model, and the importance of photon–photon absorption, this is calculated both with and without the photon–photon absorption term. Together with the total spectrum, the separate contributions for $z < z_*$ and for $z_T < z < z_*$ are reported. In Fig. 2 we show the corresponding $\Gamma$ profiles. The effect of the star surface temperature (and of the entire funnel, as the parameter $b$ is assumed to be the same for all cases) can be clearly seen in Fig. 3. Note the $\nu^{-1/2}$ power-law shape in the X-ray band for the high-temperature case. The extension of this power law branch depends on the value of $g$. In the case shown ($g = 2$) the radiation energy density outside the
funnel remains sufficiently large to cause the deceleration of the fireball, and this is responsible for the power law tail between 10 and 100 keV. For larger values of $g$ the extension of this power law decreases as $(z/z_0)^{-3b}$. If $\Gamma$ remains constant (= $\Gamma_0$), the spectrum $E(x) \propto x^2$, whereas for $\Gamma$ decreasing as $\Gamma \propto (z/z_0)^{-3-g}$

\[E(x) \propto x^{-(3-3b)/b} \quad \text{for} \quad z > z_d, \] (17)

which, for $b = 0.5$, gives $E(x) \propto x^{-3}$.

(ii) $z_d < z < z_e$: here, $\Gamma$ decreases as $(z/z_0)^{-(3-4b)}$ and thus

\[E(x) \propto x^{-(3-(1-b)/(6-7b))} \quad \text{for} \quad z_d < z < z_e, \] (18)

which, for $b = 0.5$, results in $E(x) \propto x^{-3/5}$.

(iii) $z_e < z < z_f$: at these distances the ambient radiation energy density decreases as $(z/z_0)^{-3}$. If $\Gamma$ remains constant (= $\Gamma_0$), the spectrum $E(x) \propto x^2$, whereas for $\Gamma$ decreasing as $\Gamma \propto (z/z_0)^{-3-g}$

\[E(x) \propto x^{-1/2} \quad \text{for} \quad z_e < z < z_f, \] (19)

which is independent of $g$.

In conclusion, in the case of efficient Compton drag, and independent of the particular choice of parameters, the predicted spectrum is always characterized (in order of decreasing energy) by: a steep high energy tail; a first break flagging the deceleration of the fireball; a second break corresponding to radiation produced at the top of the funnel – above which the temperature of the ambient photons remains constant; a third break, below which the spectrum is $\propto x^{-1/2}$, corresponding to the deceleration of the fireball because of the isothermal photon bath; and finally a fourth break, below which the spectrum $F(x) \propto x^2$. One obtains such a hard spectrum, instead of the familiar slope $F(x) \propto x$ corresponding to scattering of isotropically distributed electrons and seed photons, because only the photons scattered along the forward direction are observed.\(^1\)

6 DISCUSSION

If the fireball propagates in a dense photon environment the Compton-drag effect must necessarily be taken into account, and it may even be the dominant emission mechanism able to decelerate the fireball without the need of internal shocks and without invoking the build-up of large magnetic fields.

In this letter we have shown that the predicted spectrum, rather than being simply a black body spectrum boosted in energy, has a complex shape with power-law segments corresponding to the decrease in temperature of the funnel, deceleration of the fireball and dilution of the radiation energy density as the fireball propagates outside the funnel while remaining optically thick.

The general features of the predicted spectrum qualitatively agree with observations, as they can explain the steep power-law high-energy tail, the peak of the emission and a hard tail in the X-ray band. The latter feature is particularly interesting because other models made different predictions. In the standard internal shock synchrotron model, in fact, the spectrum cannot be harder than $E^{1/3}$ in the thin part, and it is very unlikely that self-absorption can take place in the X-ray band (Granot, Piran & Sari 2000). This would in fact imply a huge density of relativistic particles, making the inverse Compton effect largely dominate the total radiation output. This radiation would be emitted at higher and yet unobserved frequencies, and would then worsen the already severe efficiency problem.

In the quasi-thermal Comptonization model, on the other hand,\(^1\) this can be seen by integrating equation (7.23) of Rybicki & Lightman (1979) in the angle range $0 < \theta_h < 1/\Gamma$.

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\(\Gamma\) 2000 RAS, MNRAS 316, L45–L49 

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the typical predicted spectral shape in the X-ray band is \( \sim \rho^0 \) down to the typical frequencies of the seed soft photons, i.e. the IR–optical band (Ghisellini & Celotti 1999; Meszaros & Rees 2000).

The existing observations of a significant fraction of burst spectra harder than \( \nu^{1/3} \) (Crider et al. 1997; Preece et al. 1999a,b) are therefore already a challenge to existing models, and may suggest a Compton-drag origin of this portion of the spectrum. However, the situation is not already clear-cut because, to receive enough photons to study the spectral shape, integration times are much longer than the dynamical time-scales of the system, with the spectrum rapidly evolving in time. More sensitive instruments, such as the Burst Alert Telescope (BAT, a coded mask detector more sensitive than BATSE) onboard the foreseen Swift satellite will probably overcome this limitation.

We must also stress that the Compton-drag scenario is not alternative to the more conventional internal shock one. Indeed, the front of the fireball will decelerate first, plausibly causing subsequent non-decelerated parts to shock even if the central engine is working in a continuous way. This would produce additional radiation, either by the synchrotron and inverse Compton processes or by quasi-thermal Comptonization, depending on the details of the particle acceleration mechanism (see Ghisellini & Celotti 1999). We then expect spectral evolution: as the latter radiation mechanisms produce a steeper low-energy tail, a hard-to-soft transition (i.e. from \( \nu^2 \) to \( \nu^{1/3} \) or \( \nu^0 \)) would occur.

In this paper, we have considered the illustrative case of a single fireball moving out through an extended stellar envelope, along a funnel that is empty of matter but pervaded by thermal radiation from the funnel walls. The fireball itself (for typical parameters) remains optically thick until it expands beyond the stellar surface. A burst with complex time-structure could be modelled by a series of fireballs or expanding shells. However, in this more general case, the later shells would suffer less drag, as not enough time may have elapsed to replenish the entire funnel cavity with seed photons. Indeed one expects the spikes to be more powerful the longer the time interval is between them, as more seed photons could pervade the cavity. This, besides causing internal shocks with the first shell, which has been efficiently decelerated by Compton drag, will also result in a distribution of \( \Gamma \)-factors: they will become greater on axis, where few seed photons can efficiently Compton drag the shells, and smaller towards the border of the funnel, where seed photons can be replenished by the funnel walls.

We plan to investigate these possibilities and the consequences on the associated predicted afterglows in future work.

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