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COMPTON-DRAGGED GAMMA-RAY BURSTS ASSOCIATED WITH SUPERNOVAE

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ABSTRACT

It is proposed that the gamma-ray photons that characterize the prompt emission of gamma-ray bursts are produced through the Compton-drag process, which is caused by the interaction of a relativistic fireball with a very dense soft photon bath. If gamma-ray bursts are indeed associated with supernovae, then the exploding star can provide enough soft photons for radiative drag to be effective. This model accounts for the basic properties of gamma-ray bursts, i.e., the overall energetics, the peak frequency of the spectrum, and the fast variability, with an efficiency that can exceed 50%. In this scenario, there is no need for particle acceleration in relativistic collisionless shocks. Furthermore, although the Poynting flux may be important in accelerating the outflow, no magnetic field is required in the gamma-ray production. The drag also naturally limits the relativistic expansion of the fireball to $\Gamma \lesssim 10^4$.

Subject headings: gamma rays: bursts — radiation mechanisms: nonthermal — supernovae: general

1. INTRODUCTION

In the leading scenario for gamma-ray bursts (GRBs) and afterglows, the gamma-ray event is produced by internal shocks in a hyperrelativistic inhomogeneous wind (Rees & Mészáros 1994), while the afterglow is produced when the fireball drives a shock wave in the external interstellar medium (Mészáros & Rees 1997). Even if there is a large consensus that both gamma rays and afterglow photons are produced by the synchrotron process, recently some doubts have been cast on the synchrotron interpretation for the burst itself (Liang 1997; Ghisellini & Celotti 1999; Ghisellini, Celotti, & Lazzati 1999).

The nature of the progenitor is still a matter of active debate because the sudden release of a huge amount of energy in a compact region (generating a fireball) does not leave behind any traces of the way in which this energy has been produced. For this reason, the study of the interactions of the fireball with the surrounding medium seems to be the most powerful means of unveiling the GRB progenitor. At least two models are in competition: (1) the merging of a binary system composed of two compact objects (Eichler et al. 1989) and (2) the hypernova-collapsar model (Woosley 1993; Paczyński 1998), i.e., the core collapse of a very massive star to form a black hole.

After the discovery and multiwavelength observation of many afterglows, circumstantial evidence has accumulated for GRB explosions located in dense regions and associated with supernova-like phenomena. In fact, (1) host galaxies have been detected in many cases (Sahu et al. 1997; see Wheeler 1999 for a review), and some of them show starburst activity (Djorgovski et al. 1998; Hogg & Fruchter 1999); (2) large hydrogen column densities have sometimes been detected in X-ray afterglows (Owens et al. 1998); (3) nondetections of several X-ray afterglows in the optical band can be due to dust absorption (Paczyński 1998); (4) a possible iron line feature has been detected in the X-ray afterglow of GRB 970508 (Piro et al. 1999; Lazzati, Campana, & Ghisellini 1999a; Lazzati, Ghisellini, & Celotti 1999b), and (5) the rapid decay with time of several afterglows can be explained by the presence of a preexplosion wind from a very massive star (Chevalier & Li 1999). More recently, the possible presence of supernova (SN) emission in the late afterglow light curves of GRB 970228 (Galama et al. 1999; Reichert 1999) and GRB 980326 (Bloom et al. 1999) has added support in favor of the association of some GRBs with the final evolutionary stages of massive stars.

Although, in these models, the available energy is larger than in the case of compact binary mergers, the very small efficiency of internal shocks (see, e.g., Spada, Panaitescu, & Mészáros 1999) seems to be inconsistent with the fact that more energy can be released during the burst proper than during the afterglow (Paczyński 1999; see also Kumar & Piran 1999).

In this Letter, we show that if GRBs are associated with supernovae, Compton drag inside the relativistic wind can produce both the expected energetics and the peak energy of the spectrum of a classical long-duration GRB. In this new scenario, the efficiency is not limited by internal shock interactions, and the successful modeling of afterglows with external shocks is left unaffected.

The Compton-drag effect has already been invoked for GRBs by Zdziarski, Svensson, & Paczyński (1991) and Shemi (1994). Cosmic background radiation (at high redshift), central regions of globular clusters, and active galactic nuclei were identified as plausible sources of soft photons, but none of these scenarios were able to account for all the main properties of GRBs. However, the growing evidence for the association of GRB explosions with star-forming regions and supernovae opens up new perspectives for this scenario.

2. COMPTON DRAG IN A RELATIVISTIC WIND

We consider a relativistic ($\Gamma \gg 1$) wind of plasma propagating in a bath of photons with typical energy $\epsilon_{\rm seed}$. A fraction $\sim$min $(1, \tau_{\rm T})$ of the photons are scattered by the inverse Compton effect to energies $\epsilon \sim \Gamma^2 \epsilon_{\rm seed}$, where $\tau_{\rm T}$ is the Thomson

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opacity of the wind. Due to relativistic aberration, the scattered photons propagate in a narrow cone forming an angle $1/\Gamma$ with the velocity vector of wind propagation. By this process, a net amount of energy $E_{\rm CD}$ is converted from the kinetic energy of the wind to a radiation field propagating in the direction of the wind itself, where $E_{\rm CD} \sim \min(1, \tau_{\rm T}) V u_{\rm rad}(\Gamma^2 - 1)$. V is the volume filled by the soft photon field of energy density $u_{\rm rad}$ that is swept up by the wind.

Let us assume that the GRB fireball, instead of being made by a number of individual shells (see, e.g., Lazzati et al. 1999a), is an unsteady (in both velocity and density) relativistic wind expanding from a central point. After an initial acceleration phase, the density of the outflowing wind decreases with radius as $n(r) \propto r^{-2}$, giving a scattering probability $\sim\min[1, (r/r_0)^{-2}]$, where r_0 is the radius at which the scattering probability equals unity. After the first scattering, the photons propagates in the same direction of the flow, and the probability of a second scattering is reduced by a factor of $\sim\Gamma^2$.

If such a wind flows in a radiation field with an energy density $u_{\rm rad}(r)$, the total energy transferred to the photons when the fireball reaches a distance R is given by[2]

$$E_{\rm CD}(R) = 4\pi\Gamma^2 \left[\int_0^{r_0} u_{\rm rad}(r) r^2\, dr + \int_{r_0}^{R} r_0^2 u_{\rm rad}(r) dr\right], \quad (1)$$

where, for simplicity, we assume that a constant Γ has been reached (see also § 3). The transparency radius r_0 depends on the baryon loading of the fireball, which is parameterized by $\eta_b \equiv E/(Mc^2)$, where E/M is the ratio between the total energy and the rest mass of the fireball. Then r_0 is given by[3]

$$r_0 = 5.9 \times 10^{13} E_{52}^{1/2} \eta_{\{b,2\}}^{-1/2} \text{ cm}. \quad (2)$$

2.1. *A Simple Scenario*

We initially consider a simple scenario that can illustrate the basic features of the Compton-drag effect. Let us assume that the GRB is triggered at a time Δt (on the order of a few hours) after the explosion of a supernova (Woosley, MacFadyen, & Heger 1999; Cheng & Dai 1999). By this time, the SN ejecta, moving with velocity $\beta_{\rm SN} c$, have reached a distance $R_{\rm SN} = v_{\rm SN}\Delta t \sim 5.4 \times 10^{13} \beta_{\{\rm SN,-1\}}(\Delta t/5 \text{ hr})$ cm. Let us also imagine that the SN explosion is asymmetric, e.g., with no ejecta in the polar directions. Despite this asymmetry, the ejecta uniformly fill with radiation the entire volume within $R_{\rm SN}$. If $R_{\rm SN} < r_0$, the energy extracted by Compton drag is

$$E_{\rm CD} = \begin{cases} \dfrac{4\pi R_{\rm SN}^3}{3}\Gamma^2 u_{\rm rad}, & R_{\rm SN} \le r_0, \\[2ex] \dfrac{4\pi r_0^3}{3}\Gamma^2 u_{\rm rad}\left(3\dfrac{R_{\rm SN}}{r_0} - 2\right), & R_{\rm SN} > r_0. \end{cases} \quad (3)$$

According to Woosley et al. (1994), the average luminosity of a Type II SN[4] during Δt is on the order of $L_{\rm SN} \sim 10^{44}$ ergs s^{-1}, with a blackbody emission at a temperature $T_{\rm SN} \sim 10^5$ K. It follows that in this case, $u_{\rm rad} = aT_{\rm SN}^4 \sim 7.6 \times 10^5 T_{\{\rm SN,5\}}^4$ ergs cm^{-3} (consistent with $R_{\rm SN}$ assumed above). The efficiency ξ of Compton drag in extracting the fireball energy is very large; from the first condition of equation (3), we obtain

$$\xi \equiv \frac{E_{\rm CD}}{E} \sim 0.6 E_{52}^{-1} \beta_{\{\rm SN,-1\}}^3 \left(\frac{\Delta t}{5 \text{ hr}}\right)^3 T_{\{\rm SN,5\}}^4 \Gamma_2^2, \quad (4)$$

Note here that a high efficiency can be reached even for $\Gamma \sim 100$. Note also that the drag itself can limit the maximum speed of the expansion, even in a wind with very small baryon loading (as discussed in § 3). Each seed photon is boosted by $\sim 2\Gamma^2$ in frequency, yielding a spectrum peaking at $h\nu \sim 2\Gamma^2(3kT_{\rm SN}) \sim 0.5\Gamma_2^2 T_{\rm SN,5}$ MeV.

2.2. *A More Realistic Scenario*

The previous scenario requires that the GRB explode a few hours after a supernova. However, there is a plausible alternative, independent of whether the massive ($>30\ M_\odot$) star (assumed to be the progenitor of the GRB) ends up with a supernova explosion or not, that can produce a GRB even if the relativistic flow and the core collapse of the progenitor star are simultaneous or separated by a relatively small time interval (Woosley et al. 1999; MacFadyen, Woosley, & Heger 1999).

In fact, there is a somewhat general consensus (e.g., MacFadyen & Woosley 1999; Aloy et al. 1999; but see also Khokhlov et al. 1999) that a relativistic wind can flow in a relatively baryon-free funnel created by a bow shock following the collapse of the iron core of the star. Even if the details of this class of model are still controversial, the formation of the funnel seems to be a general outcome. Let us estimate its luminosity and, more precisely, the amount of energy in radiation crossing the funnel walls at a time t_f after its creation. With respect to the total luminosity of the star, assuming it radiates at its Eddington limit $\sim L_{\rm Edd}$, there would be a reduction by the geometrical factor equal to the ratio of the funnel to star surfaces, which is on the order of the funnel opening angle ϑ. However, immediately after its creation, the funnel luminosity is much larger than $\vartheta L_{\rm Edd}$ because of two effects that we will discuss in turn.

First, the walls of the funnel contain an enhanced amount of radiation with respect to the surface layers of the star: once "trapped" in the interior of the star, the radiation can escape through the funnel walls, thus enhancing the luminosity inside the funnel for a short time. Photons produced at a distance s from the wall surface cross it at a time $t_f \sim \tau_s s/c = \sigma n s^2/c$, where σ is the relevant cross section. The comparison of this value with the Kelvin time $t_{\rm K} \sim \sigma n R_\star^2/c$ that is needed for radiation to reach the star surface yields $s/R_\star \sim (t_f/t_{\rm K})^{1/2}$. After the time t_f, the radiation produced in the layer of width ds crosses the funnel surface carrying the energy $dE_f \sim \vartheta\tau_\star L_{\rm Edd}\, ds/c$, with the corresponding luminosity being

$$L_f \sim \frac{\vartheta}{2} L_{\rm Edd}\left(\frac{\tau_\star R_\star}{c t_f}\right)^{1/2}. \quad (5)$$

For $t_f = 100$ s and a 10 $M_\odot$ star with $R_\star \sim 10^{13}$ cm ($\tau_\star \approx 10^8$), this effect can enhance the funnel luminosity by $\sim 10^4$.

Let us now consider a second plausible enhancing factor. If the funnel has been produced by the propagation of a bow shock in the star, the matter in front of the advancing front is compressed, with a pressure increase of $\mathcal{M}^2$, where $\mathcal{M}$ is the

[2] All the calculations are made in spherical symmetry. In case of beaming, all the quoted numbers should be considered as equivalent isotropic values.

[3] Here and in the following, we adopt the notation $Q = 10^x Q_x$, using cgs units.

[4] This luminosity decreases by a factor of ~100 for Type Ibc supernovae, while the typical frequency increases by a factor of 10.

Mach number of the shock in the star. This (optically thick) gas then flows along the sides of the funnel and relaxes adiabatically to the pressure of the external matter (its original pressure). The result is that the funnel is surrounded by a sheath (cocoon) with density lower than that of the unshocked stellar material by a factor $\mathcal{M}^{3/2}$ (a polytropic index of 4/3 has been used in the adiabatic cooling). The diffusion of photons through this rarefied gas into the funnel is then even faster, resulting in a further increase of the luminosity by $\mathcal{M}^{3/4} \sim 200$, where a shock speed $\beta_{\rm sf} c = 0.1c$ (MacFadyen & Woosley 1999) and a sound speed $\beta_s c = 10^{-4} c$ have been assumed.

By taking into account both effects, the funnel luminosity corresponds to

$$L_f \sim L_{\rm Edd} \frac{\vartheta}{2} \left(\frac{\tau_\star R_\star}{c t_f}\right)^{1/2} \left(\frac{\beta_{\rm sf}}{\beta_s}\right)^{3/4} \sim 10^{45} \vartheta_{-1} \frac{M_\star}{10\ M_\odot}\ \text{ergs s}^{-1}, \tag{6}$$

which leads to an energy loss for Compton drag $L_{\rm CD} \sim \Gamma^2 L_f \sim 10^{49} \vartheta_{-1} \Gamma_2^2 (M_\star/10\ M_\odot)$, compared with the observed luminosity $\langle L_{\rm GRB}\rangle \sim 10^{49} \pi \vartheta_{-1}^2$ ergs s^{-1}. Here the average luminosity is considered over the entire burst duration: for single pulses, we should take into account an extra factor Γ^2 in equation (6) because of the Doppler contraction of the observed time.

The typical radiation temperature associated with this luminosity, assuming a blackbody spectrum, is enhanced with respect to the temperature of the star surface by $[L_f/(\vartheta L_{\rm Edd})]^{1/4} \sim (\tau_\star R_\star)^{1/8} (c t_f)^{-1/8} (\beta_{\rm sf}/\beta_s)^{3/16}$. Adopting the numerical values used above, the enhancement is on the order of 50, corresponding to a funnel temperature $T_f \sim 2 \times 10^5$ K (for a surface temperature of the star of ∼5000 K). This value is similar to the one estimated in the simple scenario of the previous subsection and thus leads to similar Compton frequencies.

3. PROPERTIES OF THE OBSERVED BURSTS

If the wind is homogeneous, the spectrum of the scattered photons resembles that of the incident photons, i.e., a broad blackbody continuum peaked at a temperature $T_{\rm drag} \sim 2\Gamma^2 T$. While the observed characteristic photon energy would therefore be $\epsilon \sim 0.5 \Gamma_2^2 T_5 (1+z)^{-1}$ MeV, in good agreement with the observed distribution of peak energies of BATSE GRBs (assuming again $\Gamma = 100$; see below), the spectrum would not reproduce the observed, smoothly broken power-law shape (Band et al. 1993). However, the assumptions of a perfectly homogeneous wind and of an isothermal radiation field are very crude, and one might reasonably expect that different regions of the wind are characterized by different values of Γ and different soft field temperatures. If we assume, e.g., that the temperature of the soft photon field varies with radius according to a power law $T(r) \propto r^{-\delta}$, the time-integrated spectrum will have a high-energy power-law tail $F(\nu) \propto \nu^{-(3-3\delta)/\delta}$. In addition, the bulk Lorentz factor of the flow is likely to vary on a timescale much shorter than the integration time required to obtain a spectrum with the BATSE data (∼1 s), and hence the analyzed spectra are the superposition of drag spectra by many different Lorentz factors. A third effect adding power to the high-energy tails of the spectrum is the reflection of upscattered photons in the pre-SN wind. These photons are scattered again by the fireball and can reach energies of $\sim 0.5\Gamma$–$50\Gamma_2$ MeV. The computation of the actual spectrum resulting from all these effects depends on many assumptions and is beyond the scope of this work.

The effects described above, which can increase the funnel luminosity over the Eddington limit, take place in nonstationary conditions. At the wind onset, it is likely that the temperature gradient in the walls of the funnel is large, but this is soon erased because of the high luminosity of the walls. This causes both the total flux and the characteristic frequency of the soft photons to decrease, and hence a hard-to-soft trend is expected. Moreover, it has been shown by Liang & Kargatis (1996) that the peak frequency of the spectrum in a single pulse at time t is strongly related to the flux of the pulse integrated from the beginning of the pulse to the time t. In our scenario, this behavior can be easily accounted for if we consider a shell slowed down by the drag itself: the Lorentz factor (and hence the peak frequency of the spectrum) at a time t is related to the energy lost by the shell, i.e., to the integral of the flux from the beginning of the pulse to the time t.

The observed minimum variability timescale is related to the typical size of the region containing the dense seed photon field, which corresponds to either $R_\star$ or $R_{\rm SN}$ depending on which of the two scenarios described above applies. Thus, the relevant light-crossing time—divided by the time compression factor—is

$$t_{\rm var} \sim \frac{R}{c\Gamma^2} \sim 3 \times 10^{-2} R_{13} \Gamma_2^{-2}\ \text{s}. \tag{7}$$

Longer timescales are instead expected if the relativistic wind is smooth and continuous.

Another interesting feature of this scenario is the possibility that the bulk Lorentz factor of the wind is self-consistently limited by the drag itself. The pressure of the soft photons starts braking the fireball in competition with the pressure of internal photons. The limiting Lorentz factor is hence reached when the internal pressure $p'_{\rm fb} \propto (T_0/\Gamma)^4$ is balanced by the pressure of the external photons as observed in the fireball comoving frame $p' \propto \Gamma^2 T_{\rm SN}^4 (1+\tau_{\rm T})^{-1}$, where $\tau_{\rm T}$ is the scattering optical depth of the wind. This gives

$$\Gamma_{\rm lim} \sim 2 \times 10^4 T_{\{\rm SN,\,5\}}^{-1/2} E_{52}^{1/4} R_{\{0,\,7\}}^{-5/8} \eta_{\{b,\,5\}}^{-1/8}, \tag{8}$$

where R_0 is the radius at which the fireball is released. Equation (8) reduces to $\Gamma_{\rm lim} \sim 10^4 (T_{0,\,11}/T_{\rm SN,\,5})^{2/3}$ if the fireball becomes transparent before reaching the coasting phase. With such high Γ, the Compton drag would be maximally efficient, causing the fireball to decelerate immediately until its Γ reaches the value given by $L_{\rm CD} = L_f \Gamma^2$, implying

$$\Gamma = \left(\frac{L_{\rm kin}}{L_f}\right)^{1/2} \sim 300 \left(\frac{L_{\{\rm kin,\,50\}}}{L_{\{f,\,45\}}}\right)^{1/2}. \tag{10}$$

In general, these limits are smaller than the maximum Γ set by the baryon load only, but they are still in agreement with the values recently inferred for GRB 990123 (Sari & Piran 1999). In addition, it is likely that the external parts of the relativistic wind, which are in closer connection with the funnel walls, are dragged more efficiently than the central ones since, at the beginning, the soft photons coming from the walls can penetrate only a small fraction of the funnel before being upscattered by relativistic electrons. This may result in a polar-structured wind, with higher Lorentz factors along the symmetry axis, gradually decreasing as the polar angle increases.

4. DISCUSSION

A crucial requirement for our model is the association of GRBs with the final evolutionary stages of very massive stars since these provide the large amount of seed photons emitted at distances $\sim 10^{13}$ cm from the central trigger, and these seed photons are needed for the Compton drag to be efficient. The efficiency of the conversion of the bulk kinetic energy of the flow into gamma-ray photons is very good, solving the observational challenge of gamma-ray emission being more energetic than the afterglow emission (Paczyński 1999). Furthermore, in this scenario, there is no requirement for either the efficient acceleration in collisionless shocks or the presence/generation of an intense (equipartition) magnetic field, although the Poynting flux may still be important in accelerating the outflow (it being more efficient than neutrino reconversion into pairs).

We have investigated the main properties of a GRB produced by Compton drag in a relativistic wind in a very general case. A moderately beamed burst ($\vartheta \lesssim 10°$; Woosley et al. 1999) can thus be produced, and, without any fine-tuning of the parameters, the basic features of classic GRBs are accounted for.

In particular, the peak energy of the burst emission simply reflects the temperature of the supernova seed photons, upscattered by the square of the bulk Lorentz factor. The simplest hypothesis predicts a quasi-thermal spectrum; however, it is easy to imagine an effective multitemperature distribution that would depend on unconstrained quantities such as the variation of the spectrum of the SN photons with radius and the degree of inhomogeneity of the wind.

Although, in this scenario, there is no requirement for the internal shocks to be set up, they can of course occur, thus contributing a small fraction of the observed gamma-ray flux. On the other hand, the wind is expected to escape from the funnel of the star with still highly relativistic motion, so that an external shock can be driven in the interstellar medium and produce an afterglow, similar to the scenario already studied by several authors. It is likely that this afterglow would develop in a nonuniform density medium because of the presence of the massive star wind occurring before the SN explosion (Chevalier & Li 1999).

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