QUASI-THERMAL COMPTONIZATION AND GAMMA-RAY BURSTS

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ABSTRACT

Quasi-thermal Comptonization in internal shocks formed between relativistic shells can account for the high-energy emission of gamma-ray bursts. This is in fact the dominant cooling mechanism if the typical energy of the emitting particles is achieved either through the balance between heating and cooling or as a result of electron-positron pair production. Both processes yield subrelativistic or mildly relativistic energies. In this case, the synchrotron spectrum is self-absorbed, providing the soft-seed photons for the Comptonization process, whose spectrum is flat \( F(\nu) \sim \text{const} \), ending in either an exponential cutoff or a Wien peak, depending on the scattering optical depth of the emitting particles. The self-consistent particle energy and optical depth are estimated and found in agreement with the observed spectra.

Subject headings: gamma rays: bursts — radiation mechanisms: nonthermal — radiation mechanisms: thermal

X-rays: general

1. INTRODUCTION

After the observational breakthrough of BeppoSAX (Costa et al. 1997; van Paradijs et al. 1997), we are now starting to disclose the physics of gamma-ray bursts (GRBs). The huge energy and power release required by their cosmological distances supports the fireball scenario (Cavallo & Rees 1978; Rees & Mészáros 1992; Mészáros & Rees 1993), whose evolution and behavior are unfortunately largely independent of its origin.

We still do not know in detail how the GRB event is related to the afterglow emission, but in the most accepted scheme of the formation of and the emission from internal/external shocks (Rees & Mészáros 1992, 1994; Sari & Piran 1997), the former is due to collisions of pairs of relativistic shells (internal shocks), while the latter is generated by the collisionless shocks produced by shells interacting with the interstellar medium (external shocks). The short spikes (\( t_s \sim 10 \text{ ms} \)) observed in the high-energy light curves suggest that shell-shell collisions occur at distances \( R \sim 10^{15} - 10^{17} \text{ cm} \) from the central source within a plasma moving with a bulk Lorentz factor of \( \Gamma \geq 100 \). The fireball starts to be decelerated by the interstellar medium farther out, at a distance that depends on the assumed density of this material.

Up until now, the main radiation mechanism that is assumed to give rise to both the burst event and the afterglow is synchrotron emission (Rees & Mészáros 1994; Sari, Narayan, & Piran 1996; Sari & Piran 1997; Panaitescu & Mészáros 1998). In fact, if electrons, protons, and the magnetic field share the available energy \( E = 10^{50} E_{50} \text{ ergs} \), the electrons reach typical random Lorentz factors of \( \gamma \sim m_e/m_p \), while the assumption of a Poynting flux \( L_B = R^2 \Gamma^2 B^2 c/2 = 10^{50} L_{50} \text{ ergs s}^{-1} \) implies a comoving magnetic field on the order of \( B \sim L_{B50}^{1/2}/(R \Gamma) \sim 10^{12} L_{B50}^{1/2}/(R_1 \Gamma_1) \text{ G} \). For these values of \( \gamma \) and \( B \), the typical observed synchrotron frequency is \( \nu_s \sim 0.5 L_{B50}^{1/2}/[R_1(1+z)] \text{ MeV} \), independent of the bulk Lorentz factor \( \Gamma \) and in excellent agreement with the observed values of the peak of GRB spectra in a \( \nu F(\nu) \) representation.

Therefore, in this model, the assumption of energy equipartition plays a key role. And, at a closer look, this implicitly requires a number of constraints to be satisfied in the emission regions. The main one concerns the acceleration of the electrons, which must be impulsive (i.e., on timescales much shorter than the cooling ones). Further requirements will be discussed in § 2.

One can envisage an alternative scenario, which we describe in § 3, where the key role is instead played by the balance between the cooling and heating processes. This would be favored if the emitting region occupies the entire shell volume rather than the narrow region associated with a planar shock, as in the “equipartition” model.

An immediate prediction following this hypothesis is that the typical energy of the emitting electrons is mildly relativistic and that the main radiation process is quasi-thermal Comptonization. This model also implies the existence of a characteristic observed frequency of a few MeV, which is controlled by the feedback introduced by the effect of electron-positron pair production.

Earlier attempts to explain burst radiation with multiple Compton scatterings have been done by Liang et al. (1997), who considered emission by a very dense population of non-thermal relativistic electrons in a weak magnetic field, in the context of bursts located in an extended Galactic halo. Liang (1997) later extended this model for bursts at gigaparsec distances, assuming an emitting region on the order of \( \sim 10^{15} \text{ cm} \), a magnetic field value on the order of 0.1 G, and thermal electrons at a temperature of a few keV.

Instead, we first discuss the constraints that the “synchrotron” scenario must satisfy and then, using the very same parameters (except for the electron energy), propose that quasi-thermal Comptonization can quite naturally dominate the cooling. Some consequences of our scenario are presented in § 4, while our findings are discussed in § 5.

For simplicity, we will consider spherical shells of comoving...
2. CONSTR ANTS ON THE EQUIPARTITION SCENARIO

Shell-shell collision produces a shock that has to accelerate particles to an energy \( \gamma m_e c^2 \) in a timescale \( t_{\text{cool}} \) much shorter than the cooling time \( t_c \); otherwise, the final random energy of the particle would be controlled by its cooling rate and not by the equipartition condition. Once accelerated to the equipartition Lorentz factor \( \gamma \approx m_e/m_p \), and radiatively cooled down, a particle will not be accelerated again during the shell-shell interaction, in order for the total random energy not to exceed the available energy in the relative bulk motion of the two shells.

The inverse Compton power must be, at most, of the same order as the synchrotron one; otherwise, the luminosity we observe in the hard X-rays would be largely dominated by the luminosity emitted beyond the GeV band, with obvious problems for the total energetics. This results in the requirement that \( \gamma t_{\gamma} \gamma^2 \leq 1 \), where \( \gamma \) is the optical depth of the emitting relativistic electrons (not to be confused with the total scattering optical depth \( \tau \) of the shell, which is of order unity at \( R \approx 10^{13} \) cm). This condition is valid if the Compton scattering process is entirely in the Thomson regime, which is appropriate as long as \( \gamma \leq 760B_{-10} \). Higher energy electrons would scatter in the Thomson-limit synchrotron photons of lower frequencies: for \( \gamma \sim 10^3 \), the reduction in the Compton luminosity due to Klein-Nishina effects is on the order of \( \sim 38^g \).

The condition \( \gamma t_{\gamma} \gamma^2 \leq 1 \) translates into demanding that the width \( dR' \) of the emitting region of the shock (assumed to be plane parallel) satisfies \( dR'/\Delta R' \leq \tau_{\gamma}^{-1} \gamma^{-2} \sim 10^{-6} \). This in turn controls the cooling timescale, which must be on the order of \( \Delta R'/c \), and we obtain

\[
t_{\text{cool}} = \frac{\Delta R'}{c \gamma^2} \sim \frac{\Delta R'}{\tau_{\gamma}} \frac{\sigma_T B^2 (1 + \gamma t_{\gamma} \gamma^2)}{6 \pi m_e c^2} \approx 86 \frac{\Delta R_{11} B_{12}^2}{\tau_{\gamma}},
\]

indicating that either the magnetic field is stronger than \( 10^4 \) G or the total shell optical depth \( \tau_{\gamma} \sim 0.1 \), to allow the electrons to reach \( \gamma \sim 10^3 \). Let us consider also that if the density of electrons is increased by electron-positron pair production, the mean energy per lepton is less than \( m_e/m_p \) by the factor equal to the ratio of leptons to proton densities \( n_{\gamma}^l/n_{\gamma}^p \).

We consider these constraints—i.e., impulsive acceleration, limited Compton power, and the absence of copious pair production—quite demanding (at least among the limits imposed by the considerations of radiation processes). The emitting region cannot be very compact and cannot have a width larger than \( dR' \). Note that in some models, the shocked region is instead considered to be complex and extended, as in the shock structure resulting from Rayleigh-Taylor instability, which can occupy the entire shell volume (Pilla & Loeb 1998). In this case, there are many more electrons emitting at a given time, with a consequent building up of the radiation energy density and an increased Compton luminosity. In addition, the larger cooling rate can limit the typical electron energies to values much below the equipartition one. In the next section, we will therefore examine the possibility that the above conditions are not satisfied.

3. QUASI-THERMAL COMPTONIZATION

Let us assume that the heating process for a typical electron lasts for the duration time of the shell-shell interaction, \( \Delta R'/c \). The maximum amount of energy given to a single lepton is on the order of \( m_e c^2 n_{\gamma}^l/m_p' \) (not to violate the total energetics), which, when released over the above timescale, corresponds to a total (average) heating rate \( E_{\text{heat}} = n_{\gamma}^l m_p c^2/\Delta R' \). The typical electron energy is given by balancing \( E_{\text{heat}} \) and \( E_{\text{cool}} \) as

\[
E_{\text{heat}} = \frac{(4/3) \sigma_T c U_{\gamma} \beta^2 (1 + U'_{\rho}/U')}{n_{\gamma}^l m_p},
\]

where \( U'_{\rho}/U' \) is the magnetic-to-radiation energy density ratio and \( n_{\gamma}^l \) includes a possible contribution from \( e^\pm \) pairs:

\[
\gamma^2 \beta^2 \equiv \frac{3 \pi R^2 n_{\gamma}^l m_p c^3}{\Delta R' \sigma_T U_{\gamma} \beta^2 (1 + U'_{\rho}/U')} \equiv \frac{n_{\gamma}^l m_p}{n_{\gamma}^m m_e} \frac{1}{1 + U'_{\rho}/U'} \frac{3 \pi}{T},
\]

where \( U' = \left[ L_{\gamma} / (R m_e c^3) / (\Delta R') / (R_{\gamma}) \right] \) is the compactness parameter of the region emitting a (comoving) luminosity \( L' \). Electron-positron pairs can be important for values of the compactness greater than unity (Svensson 1982, 1984, 1987), and they can even dominate the particle number density. A typical value in this situation is \( U' = 270 (U'_{\rho}/R_{\gamma}) (\Delta R') / (R_{\gamma}) \), and therefore equation (2) yields typical electron (and positron) energies moderately relativistic at most. As detailed below, the small energy of the emitting particles implies the following:

1. The synchrotron emission is self-absorbed.
2. The main radiation mechanism is multiple Compton scattering, and the self-absorbed synchrotron emission is the source of soft-seed photons.

Even though the particle distribution may not have time to thermalize, it will be characterized by a mean energy, and it will possibly be peaked at this value. It is then convenient to introduce an “effective temperature” \( T = \lambda T/(m_e c^2) \).

The synchrotron luminosity can then be estimated by assuming that the spectrum is described by the Rayleigh-Jeans part of a blackbody spectrum, up to the self-absorption frequency \( \nu_{\gamma}^c \):

\[
L'_{s} \sim \frac{8 \pi}{3} m_e R^2 T \Theta(\nu_{\gamma}^c)^3 \sim 7.6 \times 10^{41} \Theta R_{11}^2 (\nu_{\gamma}^c)^3 \text{ ergs s}^{-1},
\]

where \( \nu_{\gamma}^c \) can be derived again by approximating the particle distribution as a Maxwellian of temperature \( \Theta' \). An approximate prescription, which holds for \( 0.1 \leq \Theta' \leq 3 \), has been derived by interpolating the analytic approximations to the cyclotron-synchrotron emission reported by Mahadevan, Narayan, & Yi (1996). This gives, for \( B_{\parallel} \sim 1 \) and \( \gamma \sim 1, \nu_{\gamma}^c \sim 2.75 \times 10^{13} (\Theta')^{-1/3} \text{ Hz} \).

A generalized Comptonization parameter \( \gamma \), which is approximately valid in the transrelativistic and quasi-transparent conditions also, can be defined as \( \gamma = 4 \sigma_T \Theta (1 + t_{\gamma}) (1 + 4 \Theta' \Theta) \). The ratio of the Compton to the synchrotron powers can then be approximated by \( \epsilon' \), and thus, in order to emit a Compton moving luminosity \( L'_{s} = 10^{46} L'_{s} \text{ ergs s}^{-1} \), the \( \gamma \)-parameter must be on the order of \( \gamma = \ln (L'_{s} / L_{s}) \sim 11.5 \ln (L'_{s} / L_{s}) \). With this value of \( \gamma \) and with \( t_{\gamma} \) of order unity, the Comptonized high-energy spectrum has a \( F(\nu) \propto \)}
solution giving \( \omega \) and \( \phi \) respectively, the resulting observed spectrum would extend between the energies \( h\nu / T_l(1+z) \) and \( \sim 2kT_lT_l(1+z) \), with \( F(\nu) \sim \text{const} \).

The spectrum emitted by a single shell will evolve very rapidly: after the observed acceleration time \( \Delta T_l(1+z)/T_l(c) \), particles cool on a similar timescale, while the Comptonization spectrum steepens and the emitted power decreases. (Eventually, the same particles can be reheated by a collision with another shell.) The time-integrated emission (even for a short exposure time of, say, a second) will result from all the cooling and reheating histories of many shell-shell interactions. Any Wien hump and/or feature in the spectrum of individual shells will be smoothed out. The hard power-law continuum, if typical of all shells, would instead be preserved even when integrating over the exposure time.

3.1. The Role of Electron-Positron Pairs

As anticipated, \( e^+e^- \) pair production can play a crucial role: this process would surely be efficient for intrinsic compactnesses \( l' > 1 \), and, on the one hand, it would increase the optical depth and, on the other, would act as a thermostat by maintaining the temperature in a narrow range. For the temperatures of interest here, photon-photon collisions are the main pair-production process. Note that our definition of \( l' \) corresponds to the optical depth for \( \gamma\gamma \rightarrow e^+e^- \) within the shell width. Additional pairs will be produced outside the shell region, increasing the lepton content of the surrounding medium. Detailed time-dependent studies of the optical depth and temperature evolution for a rapidly varying source have not yet been pursued. Results concerning a steady source in pair equilibrium indicate that for \( l' \) between 10 and \( 10^3 \), the maximum equilibrium temperature is on the order of 30–300 keV (Svensson 1982, 1984), if the source is pair dominated (i.e., the density of pairs outnumbers the density of protons). Indeed, in this situation, we expect to be close to pair equilibrium, since this would be reached in about a dynamical timescale. Note that the quoted numbers refer to a perfect Maxwellian particle distribution. However, pairs can be created, even if the temperature is subrelativistic, by the photons and particles in the high-energy tails of the real distribution: for a particle density decreasing slower than a Maxwellian one, more photons are created above the threshold for photon-photoplus pair production, and thus pairs become important for values of \( l' \) lower than in the completely thermal case.

We conclude that an “effective” temperature of \( KT' \sim 50 \) keV (\( T' \sim 0.1 \)), and \( \tau_e \sim 4 \) dominated by pairs, can be a consistent solution giving \( \sim 11 \). We stress here that the assumption of a soft-seed photon distribution of luminosity \( L \sim 10^{41} \) ergs s\(^{-1} \) implies that any self-consistent solution must give \( \sim 10 \), in order to produce a Compton luminosity matching the observed one.

4. SOME CONSEQUENCES

If the high-energy spectrum is due to quasi-thermal Comptonization, it will be sensitive to the amount of available seed photons. Let us assume that the complex light curve of a GRB can be explained by the internal shock scenario, in which the emission is produced by the collisions of many shells. The first colliding shells will give rise to a certain synchrotron self-absorbed radiation, while subsequent shells, besides producing their own synchrotron radiation, will be illuminated by photons coming from the previous shell-shell collisions. The amount of seed photons is then bound to increase, in turn increasing the cooling rate of the electrons and positrons, which will therefore reach a lower temperature and produce a softer spectrum. This can qualitatively explain the hard-to-soft behavior of GRB emission and the (weak) correlation between duration and hardness (shorter bursts have harder spectra; Fishman & Meegan 1995 and references therein). More quantitative details will require the knowledge of the exact time-dependent feedback introduced by pair emission, which is difficult to assess.

If the intrinsic effective temperature is on the order of 50 keV, then the observed Comptonized spectrum extends to \( \sim 100\nu_1\Gamma_l/(1+z) \) MeV. Note that there are no severe constraints on such high values of the energy at which the \( vF(\nu) \) spectrum of the burst peaks (Cohen, Piran, & Narayan 1998), and these values may be reached during the very first parts of the burst light curve (e.g., the first second). Time-integrated spectra may instead be well fitted as emission from particles of lower temperatures.

One would also expect that (again, especially during the initial first phase) a Wien peak at the electron temperature would be formed for bursts with \( \tau_e \) larger than 3–5. This value of the optical depth would be easily reached for particularly strong bursts, where more copious pair production might occur.

In Comptonization models, photons of higher energies undergo more scatterings: if variability is caused by a change in the seed photon population, this may cause the high-energy flux to lag the flux at softer frequencies. In this case, the relevant timescale is the average time between two scatterings, which, in the observer frame, is on the order of \( \Delta T_l(1+z)/(\Gamma_l c) \sim 0.03\Delta T_1(1+z)/(\Gamma_l c) \) and within the possibility of current detectors. However, there can be no lag if variability is caused by a sudden increase (decrease) of the number of emitting electrons or their effective temperature. Constraints on the Comptonization scenario may come from detailed studies on how to reproduce very fast variability, since multiple scatterings will tend to smooth out any very short change.

If the progenitors of GRBs are hypernovae (Paczynski 1998), the density in the vicinity of the central source is dominated by the pre–hypernova wind. This can lead to optical depths around unity at distances \( R \sim 10^{12}–10^{13} \), which is just where shell–shell collisions are assumed to take place. There is then the possibility that the GRB events are due to shocks with this material, rather than shocks between the shells. The implications of this scenario for the emission models are very interesting and will be discussed elsewhere (Ghisellini et al. 1998). Here we would like to stress that (i) in the case of shocks between shells and the pre–hypernova wind, the large densities involved suggest that inverse Compton emission is favored with respect to the synchrotron process, and (ii) if the (still unshocked) interstellar material has total optical depth \( \tau_e \) around or greater than unity, photons will be downscattered, introducing a break in the emergent spectrum at the observed energy \( 511[\Gamma_l(1+z)] \) keV (Gulilbert, Fabian, & Rees 1983; Pozdznyakov et al. 1983). Furthermore, the interstellar matter will act as a “mirror,” sending back the scattered photons and thus increasing the amount of Compton cooling in the emitting region (see Ghisellini & Madau 1996 for an application of this mirror model to blazars).

We then conclude that the high-density environment of the hypernova poses problems for the nonthermal equipartition scenario and favors quasi-thermal Comptonization as the main radiation process. Afterglow emission would start at the de-
celeration radius \( R_2 \sim 10^{16}[E_{50}/(n\Omega^2)]^{1/3} \) cm, after \( \sim 40(E_{50}/n)\Omega^{1/3}/\Gamma_{50} \) s from the start of the burst (see, e.g., Wijers, Rees, & Mészáros 1997), where \( n \) is the number density of the interstellar material. At \( R_2 \), the reduced densities of particles and photons diminish the importance of the Compton emission even for ultrarelativistic electron energies, but the requirement \( \tau_\gamma \gamma < 1 \) can still limit \( \gamma \), especially in the case of dense star-forming environments.

5. DISCUSSION

In the equipartition scenario, the rough equality between the energy density of protons, leptons, and the magnetic field leads to a remarkably good agreement with the observed characteristics of the spectra. In order for this to be achieved, the acceleration of electrons has to be impulsive, it has to take place in a very limited volume of the interacting shell, and \( e^\pm \) pair density has to be small enough not to lower the mean lepton energy significantly.

At the other extreme, when the particle acceleration occupies the entire shell volume and/or lasts for a shell-light crossing time, the mean lepton energy is controlled by the balance between the heating and the cooling rate. This leads to mildly relativistic or subrelativistic lepton energies. Therefore, the synchrotron emission is inhibited by self-absorption and provides soft photons for the dominant inverse Compton scattering process. The ratio of the observed multiple Compton scattering luminosity to the self-absorbed synchrotron luminosity is on the order of \( 10^5 \) and determines the required Comptonization \( \gamma \)-parameter, i.e., the product of the particle optical depth and temperature. As long as the compactness of the emission region is greater than unity, relativistic temperatures cannot be achieved because, in this case, electrons and positrons are copiously produced, increasing the optical depth and decreasing the temperature. If, on the other hand, the temperature is low and the optical depth is large, photons are trapped inside the shell, and part of the radiation luminosity is used to expand it. As a result, there is a narrow range of optical depths and temperatures that accounts for the observed spectra. This may be why GRBs preferentially emit at \( \sim 1 \) MeV. A detailed analysis is needed to determine the exact shape and time behavior of the predicted spectra: studying the time evolution of hot compact sources, relaxing the assumption of pair balance, would be particularly relevant in this respect.

We conclude that there are at least two possible regimes that yield the observed spectrum and peak frequency of the GRBs, depending on the nature of the dissipation/acceleration mechanism. There is even the possibility that they coexist: one can in fact imagine that an initially planar (and “narrow”) shocked region can soon be subjected to instabilities, thus starting dissipating energy over a larger volume. In this situation, both the nonthermal “equipartition” and quasi-thermal “heating/cooling balance” regimes would be at work.

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