X-RAY–RICH GAMMA-RAY BURSTS, PHOTOSpheres, AND VARIABILITY

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Abstract

We investigate the relationship between the quasi-thermal baryon-related photosphere in relativistic outflows and the internal shocks arising outside them, which out to a limiting radius may be able to create enough pairs to extend the optically thick region. Variable gamma-ray light curves are likely to arise outside this limiting pair-forming shock radius, while X-ray excess bursts may arise from shocks occurring below it; a possible relation to X-ray flashes is discussed. This model leads to a simple physical interpretation of the observational gamma-ray variability-luminosity relation.

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

1 INTRODUCTION

Gamma-ray burst (GRB) light curves at gamma-ray energies are often highly variable, and generally this is attributed to internal shocks occurring at some radius \( r > 10^{12} \text{cm} \) from the center of a relativistic outflow produced by a violent collapse or merger, beyond the "photospheric" radius at which the flow becomes optically thin to scattering by electrons associated with the baryons entrained.

This photosphere is a source of soft thermalized radiation, which may be observationally detectable in some GRB spectra and may also result in inverse Compton cooling of the nonthermal electrons accelerated in the shocks occurring outside it, thereby enhancing a hard giga–electron volt nonthermal component at the expense of the usual mega–electron volt synchrotron component.

At small enough radii, however, the shocks can create enough pairs to reestablish a second photosphere caused by the pairs, and it is only beyond a limiting radius that the shocks remain optically thin to pairs. The most favorable region for shocks producing highly variable gamma-ray light curves is above this radius, while shocks occurring below it would lead to a second source of less variable gamma-ray light curves (see also Rees 2002; Zhang & Meszaros 2002; Salmonson & Galama 2001). Identifying the pair-shock radius as an approximate boundary above which shocks lead to more strongly variable gamma-ray light curves and below which shocks result in smoother X-ray–rich light curves, a phenomenological jet model leads to a simple physical explanation of the quantitative form of the variability-luminosity relationship.

2 BARYONIC PHOTOSPHeres

Consider a relativistic wind outflow where the bulk Lorentz factor has a mean dimensionless entropy \( \eta = L_0/Mc^2 = 10^3 \eta_3 \) that varies \((\Delta \eta - \eta)\) on timescales \( t_e \) ranging from a minimum dynamical timescale up to the maximum burst (wind) duration \( t_w \), \( 10^{-3} s \leq t_e \leq t_w \). The flow starts from a minimum radius \( r_o = ct_{e,\text{min}} = 10^7 r_0, 7 \text{cm} \), and the Lorentz factor accelerates as \( \Gamma \propto r \) up to a coasting (or saturation) radius \( r_c \approx r_0 \Gamma_f \), beyond which it coasts as \( \Gamma = \Gamma_f \). For a simple wind, neglecting finite shell effects, \( \Gamma_f = \min[\eta, \eta_*] \), where the value \( \eta_* \) is a critical value of the dimensionless entropy given by (Meszaros & Rees 2000)

\[
\eta_* \approx \left(\frac{1}{p-1}\right) = \left(\frac{L_0 \sigma_T}{4 \pi m_p c^3 r_0}\right)^{1/4} \approx 10^3 \left(\frac{L_{52} r_0^{-1}}{10^-7}\right)^{1/4}.
\]

Here \( \eta_{p,0} \) is analogous to the definition of the compactness parameter but using the proton instead of the electron mass. The coasting \( \Gamma_f \) values follow from the criterion that the proton drag time must be longer than the expansion time for protons to start to coast. Below the baryonic photosphere, protons are naturally coupled to radiation, but in the optically thin region above the photosphere, if this occurs in the accelerating regime, the protons can still be coupled to radiation and continue to accelerate out to a radius beyond the photosphere. The comoving density in the (continuous) wind regime is \( \rho' = (L_0/4 \pi^2 m_p c^3 \eta \Gamma) \), and using the above behavior of \( \Gamma \) below and above the coasting radius, as well as the definition of the Thompson optical depth in a continuous wind \( \tau_T \approx n' \sigma_T (r/\Gamma) \), we find that the baryonic photosphere where \( \tau_T = 1 \) in the wind \( (w) \) regime, due to electrons associated with baryons, is

\[
\frac{r_{\text{ph}, w}}{r_0} = \begin{cases} \eta_*/\eta_*^{1/3} \frac{1}{r < r_0 \eta} \, , \\ \eta_4^{4/3} \frac{1}{r} \, \, \, r > r_0 \eta \, , \end{cases}
\]
where for a wind $r_c \equiv r_0 \min[\eta, \eta_s]$ is the coasting radius beyond which $\Gamma = \Gamma_s$ = constant.

The accelerating and coasting behavior is followed on average also if the outflow consists of shells of duration $t_\nu = \Delta_0/c \geq r_0$, by intervals that could similarly be of order $\sim t_c$ (or a superposition of several such frequencies), leading to an oscillatory modulation of the linear and coasting behavior. Aside from such modulation, in the optically thick regime the Lorentz factor can never exceed $\Gamma = \Gamma_s$ = constant.

For a wind made up of shells of approximate duration $t_\nu$ ejected at intervals of order $t_c$, at high $\eta$ the shells move fast enough that a photon arising in one shell never crosses more than that one shell. At lower $\eta$, however, a light ray can cross many shells before escaping, and the appropriate expression for the photosphere approximates that of the wind equation (2). The criterion for the latter to be valid is that $r > \Delta_0 r_\nu(r_0 t_c)^{-1/2}$, and the transition occurs at $\eta = \eta_T$ given by

$$\eta_T = \eta_s^{4/5} = 2.5 \times 10^2 (L_{52} r_{0.7}^{-1})^{1/5}. \quad (4)$$

Thus, one has for the baryonic photosphere

$$\frac{r_{ph}}{r_0} = \begin{cases} \eta_s^{-3/2} = \eta_s^{4/5} \eta^{-3/2}, & \eta \leq \eta_T, \\ \eta_s^{1/2} - 1/2 = \eta_s^{4/5} \eta^{1/2}, & \eta > \eta_T, \end{cases} \quad (5)$$

where the first (wind regime) occurs only in the coasting regime, while the second (shell regime) applies partly in the accelerating and partly in the coasting regimes. These regimes differ from those in Mészáros & Rees (2000) by having a break at $\eta_T = \eta_s$ instead of $\eta_s$ and by having a slope $-1/2$ above $\eta_T$ instead of $-3/2$ (due to the shell regime, neglected in our previous paper). The photosphere is in the coasting wind regime for $\eta \leq \eta_T$, in the coasting shell regime for $\eta_T \leq \eta \leq \eta_s$, and in the accelerating shell regime for $\eta \geq \eta_s^{4/3}$ (see Fig. 1). When the photosphere is above the coasting radius, the final Lorentz factor is just $\Gamma_f = \eta_s$. When the photosphere is below the coasting radius, the baryons continue to be dragged by the radiation above the photosphere until $t_{\nu} \sim r_{\nu}/c \sim \Delta_0$, where $\Gamma_f \sim \eta_s^{4/5}$, $\eta_s^{1/2}$, $\eta_{s}^{1/2}$, $\eta^{1/2}$. The final Lorentz factor is thus $\Gamma_f = \eta_s^{4/5}$, $\eta_s^{1/2}$, $\eta_s^{1/2}$, $\eta^{1/2}$ for values of $\eta_s^{4/5}$, $\eta_s^{1/2}$, $\eta_s^{1/2}$, $\eta_s^{1/2}$, $\eta_s^{1/2}$. These results in shock radii (Fig. 1) that are divided into three regimes (instead of the two in Mészáros & Rees 2000, where in the wind regime the values $\eta_T$ and $\eta_{s}^{4/5}$ were collapsed into a single $\eta_T$).

In units of the initial total luminosity $L_0$ and initial temperature at $r_0$, $\Theta_0 = kT_0/m_p c^2 \simeq 2L_{52} r_{0.7}^{-1/2}$ (i.e., $T_0 \sim 1L_{52} r_{0.7}^{-1/2}$ MeV), the lab-frame baryonic photospheric luminosity $L_{ph}$ and dimensionless temperature $\Theta_{ph}$ behave as

$$\frac{L_{ph}}{L_0} = \frac{\Theta_{ph}}{\Theta_0} = \begin{cases} (T_{ph}/r_{\nu})^{-2/3} = (\eta/\eta_{s})^{-2/3} = \eta_{s}^{-2/3} (\eta/\eta_{s})^{-2/3} \eta_T, & \eta < \eta_T, \\ (T_{ph}/r_{\nu})^{-2/3} = \eta_{s}^{-2/3}(\eta/\eta_{s}) = \eta_{s}^{-2/3} (\eta/\eta_{s}) \eta_T, & \eta < \eta_s^{4/3}, \\ 1, & \eta > \eta_{s}^{4/3}. \end{cases} \quad (6)$$

Thus, $L_{ph}/L_0 = \Theta_{ph}/\Theta_0 = \eta_{s}^{-2/3} = 2.5 \times 10^{-2} [T_{ph}/25/(1 + z)]$ keV in the observer frame for $\eta = \eta_T = \eta_s^{4/3}, ~250$, $L_{ph}/L_0 = \Theta_{ph}/\Theta_0 = \eta_{s}^{-2/3} = 10^{-4}$ for $\eta = \eta_T = \eta_s^{4/3}, ~10^3$, and $L_{ph}/L_0 = \Theta_{ph}/\Theta_0 = 1$ at $\eta = \eta_s^{4/3}, ~10^4$ (where the photosphere occurs at the coasting radius).

The internal shocks, which occur in the coasting regime at radii $r/r_0 = (\Delta_0/r_0)^{1/2}$, produce a shock photon luminosity

$$L_{sh} = \epsilon_{\gamma} \epsilon_{L_0} \sim 10^{-1} \epsilon_{\gamma} \epsilon_{s_{1/3}/1/4} L_0, \quad (7)$$

where the shock efficiency $\epsilon_{sh, -1} = 10^{-1} \epsilon_{\gamma} \epsilon_{s_{1/3}/1/4}$ is a bolometric radiative efficiency when the cooling timescale is shorter than the dynamical time. Similarly, the magnetic luminosity [if the turbulent field energy $\epsilon_B = (1/3) \epsilon_B/(1/3)$ is in equipartition with that of randomized protons and electrons] is

$$L_B = \epsilon_{\gamma} \epsilon_{\gamma} \epsilon_{L_0} \sim 10^{-1} \epsilon_{\gamma} \epsilon_{s_{1/3}/1/4} L_0 \sim L_{sh}. \quad (8)$$

Thus, $L_{sh} \lesssim L_B$ for $\eta < \eta_s^{4/3} \sim 250$, $L_{ph} \lesssim L_{sh} \sim L_B$ for $\eta_T < \eta < \eta_s^{4/3} \sim 10^3$, and $L_{ph} > L_{sh} \sim L_B$ for $\eta_{s}^{4/3} < \eta < \eta_{s}^{4/3} \sim 10^4$. This means that for $\eta > \eta_s^{4/3} \sim 10^3$ the baryonic
photospheric component dominates the nonthermal internal shock component in a bolometric sense. This will lead to inverse Compton cooling of the nonthermal electrons accelerated in the shocks, causing a weakening and softening of the nonthermal synchrotron spectrum of the shock, at the expense of a hard (\(\gtrsim\) GeV) inverse Compton component, while most of the energy will be in a thermal X-ray component.

The BATSE gamma-ray luminosity is broadband in nature and can be written as \(L_\gamma \sim (1/5)(1 - \epsilon_c)cL_{sh} \lesssim (1/5)L_{sh}\), where \(cL_{sh}\) is the inverse compton efficiency, with a peak synchrotron frequency depending on the comoving magnetic field value \(B_c\). For low values of \(\eta\), shocks occur closer in, leading to higher \(B_c\) and harder synchrotron peaks. For \(\eta \sim \eta_I\) the baryonic thermal X-ray photosphere may be responsible for the X-ray excess BATSE bursts (Preece et al. 1996). For lower \(\eta\) the photospheric thermal peak is even softer, while the shocks occur closer in and produce harder synchrotron peaks approaching the upper, less sensitive end of the BATSE band, which could lead to an apparent dominance of the soft X-ray thermal photospheric peak.

For higher values \(\eta \gtrsim \eta_s\), the thermal peak tends to blend with the synchrotron peak, resembling the canonical nonthermal GRB spectrum, while for \(\eta \lesssim \eta_s\) a hard (\(\gtrsim\) MeV) thermal component would be predicted to dominate.

### 3. SHOCKS ABOVE THE PAIR-RADIUS AND VARIABLE GAMMA-RAY LIGHT CURVES

When shells of mass \(m/2\) with Lorentz factors \(\Gamma_1\) and \(\Gamma_2\) collide, the mechanical efficiency for conversion of kinetic energy \(m^2c^2(\Gamma_1 + \Gamma_2)/2\) into internal energy is \(\epsilon_i = (\Gamma_1 + \Gamma_2 - \sqrt{2(\Gamma_1 \Gamma_2)}/\Gamma_1 + \Gamma_2)\), where as before we parameterize \(\epsilon_i = (1/4)\epsilon_{i,1/4}\). If a total of \(2N\) shells are ejected that collide and the total isotropic equivalent kinetic energy of outflow is \(E_{\text{iso}}\), the corresponding internal energy produced in the merger of two shells is \(E_{\text{int}} \sim \epsilon_i N^{-1}E_{\text{iso}}\). Of that, a fraction \(\epsilon_c\) is given to electrons, and for a high radiation efficiency in the mega-electron volt range and a high compactness parameter (i.e., high efficiency of pair formation), a fraction of order \(\epsilon_c\) of the radiated energy could be converted into pairs, and the energy in pairs in the merged shell is

\[
E_{\text{e}} = \epsilon_c \epsilon_{i,1/3} N^{-1}E_{\text{iso}} \sim 10^{50.5} \epsilon_c \epsilon_{i,1/3} N^{-1}E_{\text{iso}} \text{ ergs}. \tag{8}
\]

Assuming that \(\Gamma\) is in the range \([\Gamma_m, \Gamma_M]\), with \(\Gamma_m < \Gamma_M\), the observed radiation comes mainly from collisions involving shells at the extremes of this range and is maximized for \(\Gamma_m \ll \Gamma_M\). Such merged shells move with a center-of-momentum (CM) Lorentz factor \(\Gamma_0 \sim \sqrt{\Gamma_m \Gamma_M}\). For shells of initial lab-frame widths \(\Delta_0 \sim r_0\), for radii above the shock radius \(r_{sh} = \rho_0 r_0^2\) (which is also the “expansion radius” above which the comoving width \(\propto r\) and the comoving volume \(\propto r^3\)), the energy radiated in the shocks can be enough to create pairs that make the shocked shells optically thick to Thomson scattering, if \(\eta\) is below a certain value for which the comoving radiation compactness parameter \(l^* \sim L_{\text{com}}/mc^2 r^3 \gtrsim 1\) (Meszaros & Rees 2000). Earlier simulations involving randomly ejected shells and (baryonic) electron scattering in shocks have indicated a tendency for more variable light curves arising in more distant shocks (Panaitescu, Spada, & Meszaros 1999; Spada et al. 2000; Ramirez-Ruiz & Lloyd-Ronning 2002), as expected since closer in the scattering depth is larger. Similar results are obtained numerically when pair formation is included (e.g., Kobayashi et al. 2002). Here we pursue a simplified analytical description. For shocks at increasing radii, the pair comoving scattering depth of the shells eventually drops to unity, \(\tau_c \sim \eta_c^2 \epsilon_{i,1/4} \Gamma_c^2 \tau_{\epsilon,1/4} \sim 1\) at a characteristic limiting pair-producing shock radius

\[
r_{\pm} \sim \left(\frac{E_\gamma \Gamma_c^2}{4\pi mc^2 \epsilon_{i,1/4}}\right)^{1/2}
\sim 3 \times 10^{14} \left(\epsilon_{i,1/4} \epsilon_{\gamma,1/3} N_2^{-1} E_{54}\right)^{1/2} \left(\frac{\Gamma_M}{3 \Gamma_{m,2}}\right)^{-1/4} \text{ cm}, \tag{9}
\]

(Kobayashi et al. 2002), where \(\Gamma_m = 10^2 \Gamma_{m,2}\) and \(\Gamma_M = 10^3 \Gamma_{M,3}\). This is in the discrete shell regime, as opposed to the wind regime used by Meszaros & Rees (2000); in general, the shell regime pair density exceeds the wind regime pair density by a factor \(\left(t_{\gamma}/t_{\epsilon}\right) (\Gamma_c/\Gamma_m)^3 \gtrsim 1\), as expected since the same kinetic energy density is concentrated in shells rather than smoothed out. At this radius both the comoving scattering time and the pair formation time as well as the comoving pair annihilation time \(\left(n_c \epsilon_{i,1/4} \Gamma_c^2 \epsilon_{\gamma,1/4} \Gamma_{\gamma,1} \sim 1\right)\) become equal to the comoving expansion time \(r_{\pm}/c\Gamma_c\).

Shells with Lorentz factors \(\Gamma_M\) and \(\Gamma_m\) ejected from a starting radius \(r_0\) at time intervals \(t_{\gamma} = \Delta_0/c\) collide at a radius \(r_{sh} \sim c t_{\gamma} \Gamma_{M,3}^{-2} \gtrsim \rho_0 r_0^2\). If this shock radius is outside the limiting pair-shock radius \(r_{\pm}\), given by equation (9), pairs do not form in the shock, whereas in the opposite case an optically thick pair region does form in the merged shell, which expands until it reaches the radius \(r_{\pm}\). The shocks that occur outside the limiting pair-shock radius \(r_{\pm}\) are those for which the corresponding shells started out from \(r_0\) with a minimum time difference \(t_{\gamma} > t_{\epsilon,\pm}\), where

\[
t_{\epsilon,\pm} \sim r_{\pm}/c\Gamma_c^2 \sim 0.2 \left(\epsilon_{\gamma,1/4} \epsilon_{\epsilon,1/3} N_2^{-1} E_{54}\right)^{1/2} (\Gamma_{M,3}^{-5/4} \Gamma_{m,2}^{-5/4})^{1/4} \text{ s}. \tag{10}
\]

If the shell ejection time differences \(t_{\gamma}\) have random realizations between the minimum and maximum values \([t_{\epsilon,\pm}, t_{\epsilon,\pm}]\) over the total duration of the burst outflow \(t_b \gtrsim t_{\epsilon,\pm}\), out of the \(N\) shells ejected there will be, on average, a fraction \(\left(1 - t_{\gamma}/t_{\epsilon,\pm}\right)\) that will lead to shocks outside the pair-shock radius. For a high radiative efficiency, a fraction 0.5 of \(e_{\gamma,1/3}\) of which is taken to be in the gamma-ray range, the isotropic-equivalent gamma-ray fluence of the shocks above the limiting pair-shock radius is approximately

\[
E_\gamma \sim \frac{1}{2} \epsilon_{\gamma,1/2} \epsilon_{\epsilon,1/4} E_{54} \left(1 - \frac{t_{\epsilon,\pm}}{t_{\epsilon,\pm}}\right)^{1/2}
\sim 4 \times 10^{52} \epsilon_{\gamma,1/2} \epsilon_{\epsilon,1/4} N_2^{-1} E_{54} \left(1 - \frac{t_{\epsilon,\pm}}{t_{\epsilon,\pm}}\right) \text{ ergs}. \tag{11}
\]

Here \(t_{\epsilon,\pm}/t_{\epsilon,\pm} \sim 2 \times 10^{-2} (\epsilon_{\gamma,1/3} \epsilon_{\epsilon,1/4} N_2^{-1} E_{54})^{1/2} t_{\gamma}^{1/2} t_{\epsilon,\pm}^{-5/4} \frac{\Gamma_{M,3}^{-5/4}}{\Gamma_{m,2}^{-5/4}} \ll 1\), with \(t_{\epsilon,\pm} \lesssim t_{\epsilon,\pm} < t_b\) where \(t_b = 10 t_b, s\) is the burst duration. In this simple model \(E_\gamma\) represents the energy in the variable gamma-ray component of the burst, which arises above \(r_{\pm}\) and has variability on timescales \(\gtrsim t_{\epsilon,\pm}\). For \(t_{\epsilon,\pm} \lesssim t_b\), \(E_\gamma\) is insensitive to \(t_{\epsilon,\pm}\), but for short bursts or for \(t_{\epsilon,\pm} \gtrsim 0.1 t_b\) there is a dependence of \(E_\gamma\) on \(t_{\epsilon,\pm} \propto \Gamma_{M,3}^{-5/4}\). For small \(\Gamma_m\) the typical pair-shock radius \(r_{\pm}\) is farther out, and
the minimum variability timescale $t_{\pm}$ is longer, with a consequently smaller variable $E_{\gamma}$ (fewer bursts occur outside the more distant limiting pair-shock radius). Larger $T_m$ lead to smaller limiting pair-shock radii, shorter minimum variability timescales $t_{\pm}$, and larger isotropic equivalent $E_{\gamma}$.

4. PAIR-PRODUCING SHOCKS AND X-RAY–RICH COMPONENT

For shocks occurring at $r_{sh} \leq r_{\pm}$, i.e., below the limiting pair-shock radius where shocks can result in pair formation, the scattering optical depth of the shocked shells can become $\tau_{\pm} \geq 1$ (even when the shock is above the baryonic photosphere given by eqs. [2] and [3]). Pair formation causes the same amount of shock energy to be spread among a larger number of particles (new pairs) than in a purely baryonic outflow, and inverse Compton losses due to upscattering of its own photons (Ghisellini & Celotti 1999) become important. For a pair of shells undergoing a shock at $r < r_{\pm}$, it is expected that pair production acts as a thermostat, and for comoving compactness parameters $10 < l_{\pm} \leq 10^3$, the comoving pair temperature is $T_{\pm} \sim 3–30$ keV, with $\tau_{\pm}$ few (Svensson 1987). The scattering depth per shock due to pairs is unlikely to be much larger, because the scattering and the pair production cross sections are comparable, and unless dissipation and pair formation occur uniformly throughout the entire volume, downscattering of photons above the pair threshold rapidly leads to self-shielding (E. Ramirez-Ruiz et al. 2002, in preparation). As a specific example, we take $T_{10} \sim 10$ keV and $\tau_{10} \sim 3$ for one shock, producing a comoving spectrum peaked near $h\nu \sim 3kT_{10}/\tau_{10}^2 \sim 3T_{10}^{2/3}\tau_{10}^{2/3}$ keV. Since at any time there may be more than one shock at $r < r_{\pm}$, the photons might encounter more than one shell before escaping (e.g., Spada et al. 2000) and would also undergo adiabatic cooling between the shells by a factor $\sim r_{sh}/r_{\pm}$. These two effects combined could lower the escaping photon energy by a factor roughly estimated as $\zeta \sim 0.2$–0.5. For a CM bulk Lorentz factor $\Gamma_{cm} = (\Gamma_0M_p)^{1/2} = 300\Gamma_{1.5,25}$, the observer-frame pair-producing shock radiation peak is at

$$h\nu_{X,sh} \sim 100 T_{10}^{2/3}\tau_{10}^{2/3}\tau_{30,2}\Gamma_{2.5,25}[2/(1 + z)] \text{ keV} \quad (12)$$

The peak energy (eq. [12]) is still substantially above the blackbody value $T_{sh,bb} \sim 4(\epsilon_{1.5}\epsilon_{1.4}E_5N_2)^{-1/8}\Gamma_{2.5,25}$ keV. The BATSE distribution of peak energies (Preece et al. 2000) has $\sim 10\%$ of bursts with $h\nu_{pk} \leq 100$ keV, while the joint BATSE-BeppoSAX distribution of Kippen et al. (2002) shows that most XRFs have peak energies in the 20–100 keV range, with one exception at $3^{+8}_{-5}$ keV. A nominal value of $h\nu_{pk} \sim 20$ keV can be obtained from equation (12) with, e.g., $\Gamma_{cm} = 12(1 + z) \sim 60$.

For completeness, we note that in the extreme case where pairs are produced uniformly throughout the entire volume, thermalization and an equilibrium pair optical depth $\tau_{\pm} \propto \nu^{1/2}$ might be achieved, where $\nu$ is the comoving compactness (Guilbert, Fabian, & Rees 1983; Svensson 1987), although we expect $\tau_{\pm}$ in this case to be much smaller.

The energy in the X-ray component from shocks arising below the limiting pair-shock radius $r_{\pm}$, integrated over the burst duration $t_{bp}$, is the complement of the gamma-ray energy produced in shocks arising above $r_{\pm}$ (see eq. [11]). The X-ray isotropic equivalent fluence is

$$E_{\nu} \sim \frac{1}{2} \epsilon_k \epsilon_{\nu} \epsilon_{iso} \frac{t_{sh}}{t_{\nu}} \sim 4 \times 10^{50} \epsilon_k \epsilon_{iso} \epsilon_{\nu} \epsilon_{sh} \left(\frac{\epsilon_{1.5} \epsilon_{1.4} E_5 N_2^{-1}}{18} \right)^{3/2}$$

$$\times N_2^{-1/2} \Gamma_{cm}^{1/2} \Gamma_{2.5}^{-1/4} \Gamma_{30,2}^{-1/4} \text{ ergs} \quad (13)$$

where $t_{sh}/t_{\nu}$ is given below equation (11) and $\epsilon_k = (1/2)\epsilon_k 1/2$ is an efficiency factor to account for a fraction of order unity of the luminosity below $r_{\pm}$ that is reconverted into kinetic energy. This X-ray component could account for most of the harder XRFs (Heise et al. 2001; Kippen et al. 2002), with $h\nu_{pk} \gtrsim 20$ keV, but if more XRFs are observed with peaks as low as 3–5 keV this may require additional considerations. On the other hand, the X-ray excess GRBs discussed by Preece et al. (1996) have characteristics that, as a class, are close to those of the $r_{sh} \ll r_{\pm}$ pair-producing shocks discussed in this section.

The radiation from pair shocks with $\tau_{\pm} \sim 10$ would be subject to a moderate amount of time-smoothing $\Delta t_{\var} \sim \Delta t_{\var, \text{org}} \tau_{\pm}$, which partially degrades the original variability implied by the random ejection and shocking of shells. The smoothing would be more appreciable at the shorter timescales, where it would lead to a filling in of the narrow troughs between peaks (see also Panaitescu et al. 1999; Kobayashi et al. 2002; Ramirez-Ruiz & Lloyd-Ronning 2002). This smoothing, however, would not be expected to affect the coarser time structure of the light curve, since not many scatterings are incurred before the photons are advected with the flow.

5. VARIABILITY DEPENDENCE ON GAMMA-RAY LUMINOSITY

An observational correlation (Fenimore & Ramirez-Ruiz 2000; Reichart et al. 2001) has been reported between the isotropic equivalent luminosity $L_{\gamma}$ and a variability measure $V$ of the gamma-ray time profiles, of the form

$$L_{\gamma} \propto V^p, \quad g \approx 3.3^{+2.5}_{-1.1} \quad (14)$$

The operational definition of $V$ is related to the normalized variance, or the root mean square of the deviations from a smoothed light curve. Observations of afterglows with breaks in the light curves are believed to indicate the presence of a collimated jetlike outflow. The simplest interpretation assumes a uniform jet cross section (independent of angle out to a jet edge $\theta$), in which case the variety of break times indicates a variety of jet opening angles and the data indicate an isotropic equivalent fluence anticorrelation with jet opening angle $\theta$ (Frail et al. 2001), of the form

$$L_{\gamma,iso} \propto \theta^{-2}.$$  Alternatively, the same data can be interpreted in terms of a nonuniform (angle dependent) cross section jet with a universal jet pattern given by the same functional relation between the energy output as a function of angle $L_{\gamma}(\theta) \propto \theta^{-2}$ (Rossi et al. 2002; Zhang & Meszaros 2002; Salmonson & Galama 2002). In the latter case the data are interpreted as sampling different offsets between the observer line of sight and the jet axis. Norris (2002) and Salmonson & Galama (2002), analyzing the time-lag effects (see below) in a larger sample and including redshift and luminosity function effects, argue for a somewhat steeper angular index of $-2.5$ so

$$L_{\gamma} \propto \theta^{-p}, \quad p \sim 2–2.5.$$  (15)
In the previous sections we used an isotropic outflow, but our results continue to apply to the jet case as long as $\Gamma$ exceeds the inverse of the jet opening angle $\theta$. The model interprets the variable gamma-ray luminosity as that portion which arises from shocks above the limiting pair-shock radius $r_{\pm}$, characterized by a minimum time variability given by equation (10), $t_{\pm} \propto L_{\gamma}^{-1/3} \Gamma^{-4/3}$. This is based on equation (11) relating $E_{\gamma}\,\tau_{\gamma} \sim E_{\gamma}$ to $E_{\gamma}$ and the assumption that the average mean duration and redshift differences are overshadowed by source-intrinsic variations in $E_{\gamma}$ and $\Gamma$, so that approximately $L_{\gamma} \propto E_{\gamma}$, and assuming that $t_{\pm} \ll t_{\gamma\gamma}$, the crucial dependence of $t_{\pm}$ through $\Gamma$ rather than $\Gamma_{m}$, since it is $\Gamma_{m}$ that determines the shock radius. It is reasonable to make the Ansatz $\Gamma_{m} \propto \theta^{-2}$, and a value $q \sim 2$ (e.g., MacFadyen & Woosley 1999; Kobayashi et al. 2002) follows from momentum conservation in a “sharp boundary jet” model where the energy and $\Gamma$ are constant throughout its cross section but there is a range of opening angles (e.g., Frail et al. 2001). In a “universal jet profile” model where $L$ and $\Gamma$ vary as functions of $\theta$ (Rossi et al. 2002; Zhang & Mészáros 2002; Salmonson & Galama 2002), a value $q \sim 2$ is also expected, e.g., if the baryon loading in the jet is approximately independent of $\theta$ but the energy varies as $\theta^{-2}$ (e.g., eq. [15]). Setting $\Gamma_{m} \propto \theta^{-q}$, we have then $t_{\pm} \propto L_{\gamma}^{1/3} \Gamma_{m}^{-4/3} \propto L_{\gamma}^{1/3} \Gamma_{m}^{-4/3}$. The variability $V_{\gamma}$ of the gamma-ray light curves could be expected to scale, in an approximate way, inversely proportional to a power of the minimum variability timescale, $V_{\gamma} \propto t_{\gamma\gamma}^{-q/2}$. An approximate argument shows that such an anticorrelation exists in the GRB data, with an index $k \approx 2/3$ (e.g., Ioka & Nakamura 2001; Plaga 2001). Identifying $t_{\gamma\gamma}$ with $t_{\gamma\gamma}$, we have

$$L_{\gamma} \propto V_{\gamma}^{3/2} / [5 (q - 2)] . \quad (16)$$

If one takes $p = q$, which may be too idealized, the theoretical relation is $L_{\gamma} \propto V^{2}$, which is comparable to the lower limit fit of Fenimore & Ramirez-Ruiz (2000); the same result is obtained for $p = 5/2, q = 2, k = 1$. Using the nominal values $p = 5/2, q = 2$, and $k = 2/3$, we get $L_{\gamma} \propto V^{3}$, in good agreement with the observed best-fit relation $L_{\gamma} \propto V^{3.3}$ of Fenimore & Ramirez-Ruiz (2000).

### 6. DISCUSSION

We have discussed the properties of the quasi-thermal baryonic photospheric radiation component in GRBs. At high isotropic equivalent luminosities, this component can dominate the nonthermal shock component and appears in the hard X-ray range in the source frame. Such sources may be identified with the X-ray excess (Preece et al. 1996) class of bursts. This photospheric quasi-thermal component can inverse Compton cool the nonthermal electrons in the shocks above it, suppressing the mega–electron volt synchrotron component and enhancing an inverse Compton giga–electron volt nonthermal component. For high dimensionless entropy $\eta \sim L_{\gamma} / M c^{2}$ and low ($z \leq 1$) redshifts, the quasi-thermal component appears at hard X-rays (and in extreme cases at gamma rays), whereas for high redshifts (and/or low $\eta$) it appears at soft X-rays. We also have identified a new regime in the description of baryonic photospheres from relativistic outflows, which is valid at moderate to high $\eta$. The value of the final coasting Lorentz factor of the outflow is not automatically the value it has when the flow becomes optically thin, and has three different possible values $\eta = \eta_{\gamma}^{2/3} \eta_{\gamma}^{-1/2}$, and $\eta_{\gamma}$ as discussed below equation (5), depending on the value of the initial dimensionless entropy $\eta$.

We have quantified the location of the outermost radius at which pairs can form in internal shocks and have argued that highly variable gamma-ray light curves arise mostly from shocks above this limiting pair-shock radius. The pair-shock radius determines the approximate ratio of the fluences in a variable gamma-ray nonthermal component and in a less variable softer ($\gtrsim 20–25$ keV) X-ray component. The latter could also be responsible for X-ray excess GRBs and, for moderately low bulk Lorentz factors or moderately high redshifts $\Gamma [2 / (1 + z)] \gtrsim 60$, would be similar to most of the currently known XRF bursts (Heise et al. 2001; Kippen et al. 2002), but additional considerations may be needed to fit naturally the softest ($3–5$ keV) XRFs. Smoother X-ray components are also obtained from closer in shocks neglecting pair formation (e.g., Ramirez-Ruiz & Lloyd-Ronning 2002; Spada et al. 2000), but smoothing and softening is stronger when there is pair formation (see also Kobayashi et al. 2002). This pair X-ray component is generally softer than that of the baryonic photosphere. When present, the pair photosphere enshrouds the baryonic photosphere, but its modest opacity $n_{\pm} \lesssim 3$ is not sufficient to alter significantly the spectrum of the baryonic photosphere. One or both of these X-ray–rich components may be present, depending on the bulk Lorentz factor and isotropic equivalent total energy of the burst, and criteria are discussed for the nonthermal gamma-ray components to dominate over, or be dominated by, these X-ray components.

An individual burst is characterized in Figure 1 by an average $\eta = L_{\gamma} / M c^{2}$. The shock radius $r_{sh, 0}$ plotted in Figure 1 is for the minimum variability time $t_{\gamma} = t_{0} \sim 0.3$ ms, and a second shock radius $r_{sh, 3}$ is shown for $t_{\gamma} = 10^{3} t_{0} \sim 0.3$ s. For $t_{\gamma} \sim t_{0}$ and $\eta \lesssim \eta_{\gamma}$, the corresponding shocks occur between the baryonic $r_{ph}$ and pair-shock $r_{\pm}$ photospheres, leading to X-ray–rich bursts whose variability is partially suppressed. For $t_{\gamma} \sim 10^{3} t_{0}$ and $\eta \lesssim \eta_{\gamma}$, shocks occur at or above both $r_{ph}$ and $r_{\pm}$, leading to hard gamma rays with large variability at $\gtrsim 0.3$ s. An individual burst may have several variability timescales present, leading to both types of components simultaneously. Preponderance of one or the other leads to a short timescale but low amplitude variability X-ray–rich bursts or XRFs, or to a classical hard GRB with large amplitude variability mostly at $\gtrsim 0.3$ s. Roughly speaking, XRFs would be expected from the region $\eta < \eta_{\gamma} \sim 250$ and classical GRBs from $\eta \gtrsim \eta_{\gamma}$. A baryonic photosphere component should be present at the beginning of bursts and XRFs and in the troughs between harder peaks due to shock radiation. However, the farther beyond the coasting radius the photosphere occurs, the weaker its energy fraction is relative to the shock and/or $e^{\pm}$ component, because its energy drops as $r^{-2/3}$. Low values of $\eta$ lead to further out, weaker baryon photospheres and, at the same time, to harder, relatively stronger shocks occurring closer in to the photosphere. If long variability timescales are absent and $\eta \lesssim \eta_{\gamma}$, a soft baryon photosphere may be the most prominent component, but its total energy would be very low. A strong baryonic photosphere–dominated burst (with quasi-thermal $\gtrsim \text{MeV}$ spectrum) is possible (for shock efficiencies $\lesssim 0.1$) for $\eta \lesssim \eta_{\gamma}$, and such a component may be detectable already for $\eta \gtrsim \eta_{\gamma} \sim 10^{3}$. On the other hand, slower $\eta \lesssim \eta_{\gamma} \sim 250$ outflows are likelier to make X-ray–rich bursts through a pair-shock component.
We argue also that the relationship between variable gamma-ray radiation and the limiting pair-shock radius leads, using the phenomenologically inferred dependence between isotropic luminosity and jet angle, to a simple analytical interpretation for the observed variability-luminosity relation $L \propto V^3$ (e.g., Fenimore & Ramirez-Ruiz 2000; see also Kobayashi et al. 2002). The positive correlation between variability and harder $\nu F_\nu$ peaks discussed by Lloyd-Ronning & Ramirez–Ruiz (2002) also finds a qualitatively similar interpretation in terms of a higher variability corresponding to closer in shocks, which are more specifically in the present model shocks occurring just above the limiting pair-forming shock radius.

Several physical explanations have been proposed for the presence of a cutoff above about 1 Hz in the power density spectrum of GRB light curves (Beloborodov, Stern, & Svensson 1998) in terms of baryonic electron scattering (e.g., Panaitescu et al. 1999; Spada et al. 2000; Ramirez-Ruiz & Lloyd-Ronning 2002). Here we point out a different explanation for this, which is based on the existence of a minimum gamma-ray variability timescale (eq. [10]) above which shock radiation is free from smoothing by opacity from pair formation. This is seen in Figure 1, which shows that shocks associated with the variability timescales $t_v \gtrsim 1$ s and $\eta \gtrsim 150$ common among observed GRBs occur above the pair photosphere.

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