OPTICAL AND $\gamma$-RAY EMISSIONS FROM INTERNAL FORWARD–REVERSE SHOCKS: APPLICATION TO GRB 080319B?

Y. W. Yu\textsuperscript{1,2}, X. Y. Wang\textsuperscript{1}, and Z. G. Dai\textsuperscript{1}

\textsuperscript{1}Department of Astronomy, Nanjing University, Nanjing 210093, China; yuyw@nju.edu.cn, xywang@nju.edu.cn, dzg@nju.edu.cn
\textsuperscript{2}Institute of Astrophysics, Huazhong Normal University, Wuhan 430079, China

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ABSTRACT

In the popular internal shock model for the prompt emission of gamma-ray bursts (GRBs), collisions between a series of relativistic shells generate lots of paired forward and reverse shocks. We show that the synchrotron emission produced by the forward and reverse shocks respectively could peak at two quite different energy bands if the Lorentz factors of these two types of shocks are significantly different from each other (e.g., one shock is relativistic and the other is Newtonian). We then investigate whether this scenario is applicable to the case of GRB 080319B and find that a bimodal distribution of the shell Lorentz factors, peaking at $\sim 400$ and $\sim 10^5$, is required. In addition, this scenario predicts an accompanying inverse-Compton (IC) GeV emission with a luminosity comparable to (not much higher than) that of the synchrotron MeV emission, which can be tested with future Fermi observations.

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

1. INTRODUCTION

Since the pioneering works by Rees & Mészáros (1994) and Paczyński & Xu (1994), it has been widely argued that the prompt emission of gamma-ray bursts (GRBs) arises from internal shocks in a relativistic fireball that consists of a series of shells with different Lorentz factors and that the observed $\gamma$-ray emission is usually attributed to synchrotron or inverse-Compton (IC) emission from power-law electrons in the shocks (Mészáros & Rees 1997). Especially, in the case of IC $\gamma$-ray emission, we can also expect a bright synchrotron low-energy (e.g., optical) emission as seed photons for IC scattering. Based on this synchrotron and synchrotron self-Compton (SSC) emission scenario, Mészáros & Rees (1999) accounted for the bright optical flash of GRB 990123.

In this paper we put forward an alternative model, in which prompt optical and $\gamma$-ray emissions may be produced only by synchrotron radiation in internal shocks. It is natural to expect that internal shocks should consist of paired forward and reverse shocks, which are produced simultaneously by collisions between the shells in the fireball. In the case where the Lorentz factors of these two types of shocks are quite different, two groups of electrons accelerated by a forward and a reverse shock respectively during one collision are expected to reach different characteristic energies. Therefore, by combining the two synchrotron components contributed by these two shocks, we can get a bimodal photon spectrum, which might be able to account for some GRBs that have two spectral components. In particular, under some peculiar conditions, an optical peak can be found to fit the prompt optical emission of some GRBs such as the naked-eye GRB 080319B.

For GRB 080319B, with its prompt $\gamma$-ray fluence ($6.13 \pm 0.13) \times 10^{-4}$ erg cm$^{-2}$ (20 keV–7 MeV; Racusin et al. 2008b) and redshift $z = 0.937$ (Vreeswijk et al. 2008), its isotropic equivalent $\gamma$-ray energy release is estimated to be $E_{\gamma,\text{iso}} = 1.4 \times 10^{54}$ erg (20 keV–7 MeV), which is among the highest ever measured. More surprisingly, this GRB was found to be associated by an extraordinarily bright optical flash peaked at a visual magnitude of 5.3 (Racusin et al. 2008b; Bloom et al. 2009), which is visible even for the unaided eye. Compared with the extrapolation of the $\gamma$-ray spectrum to the optical band, the observed flux density of the optical flash ($\sim 20$ Jy) is about 10,000 times higher. In addition, thanks to the high time resolution, a fluctuating structure can be seen clearly in the optical light curve.

For this rare multiwavelength prompt emission, any scenario invoking a single synchrotron component from only one emitting region is obviously not viable. Moreover, different from some other bright optical flashes such as those of GRBs 990123 and 061007, an external shock origin for the optical flash of GRB 080319B is disfavored due to the short duration of the optical pulse and the lack of an increasing optical pulse duration throughout the whole prompt phase, as argued by Kumar & Panaitescu (2009). Instead, an internal dissipation origin may be plausible. As argued by Mészáros & Rees (1999) for GRB 990123, Kumar & Panaitescu (2009) and Racusin et al. (2008) explained the prompt emission of GRB 080319B in the synchrotron and SSC emission scenario. However, as a direct consequence, they predicted a remarkably strong GeV emission, whose luminosity is about 10 times higher than the observed MeV one. In contrast, for several bright GRBs detected by EGRET on Compton Gamma Ray Observatory (CGRO), the GeV fluence is not higher than that in the MeV Burst and Transient Source Experiment (BATSE) energy band (e.g., Sommers et al. 1994; Hurley et al. 1994).

However, in our internal shock-produced two-component synchrotron emission scenario, the problem of GeV emission excess can be avoided, whereas the prompt optical and MeV $\gamma$-ray fluxes of GRB 080319B can be still explained if an unusually high variability of Lorentz factors exists in the fireball. Moreover, since internal shocks can take place many times, we can naturally understand the fluctuating structure of the optical and $\gamma$-ray light curves. However, the emissions in the two energy bands are not strictly correlated with each other, since the durations of the shocks determined by the widths of the shells could be different and the specific shape of the light curves is dependent on the specific structure of the shells. For simplicity, these complications will not be considered in our simple model.

Compared with the internal shock model suggested by Mészáros & Rees (1999) for GRB 990123, our work differs from theirs in two aspects. First, as discussed above, we suggest that
both optical and γ-ray emissions arise from synchrotron emission, differing from the synchrotron plus SSC scenario proposed by them. Second, in their works, only nonrelativistic to semirelativistic internal shocks are considered, which are produced by collisions between shells with similar Lorentz factors, while in our paper some relativistic internal (reverse) shocks would be involved in order to fit the observations of GRB 080319B.

This paper is organized as follows: in Section 2, we describe the dynamics and synchrotron emission of internal forward–reverse shocks and the model parameters are expressed as functions of some observational quantities. In Section 3, we constrain the model parameters by the observations of GRB 080319B and then some implications from these results are discussed. In addition, the contribution to the prompt emission by IC scattering of the electrons is also considered with the inferred model parameters. Finally, conclusions and discussion are given in Section 4.

2. TWO-COMPONENT SYNCHROTRON EMISSION FROM INTERNAL SHOCKS

2.1. The Dynamics and Electron Energy Distributions

In the internal shock model, the central engine of GRBs is assumed to eject a fireball consisting of a series of shells with different Lorentz factors γ_{shell} during the prompt phase. Considering two shells that are ejected subsequently, for example, if the posterior shell (denoted by 4) moves more rapidly than the prior one (denoted by 1), a collision takes place at radius \( R_{\text{c}} \approx 2\gamma_{i}c\delta t \phi_{i}^{-1} \) (Yu & Dai 2009) to produce an emission pulse, where \( \delta t \) is the observed variability time and \( \phi_{i} = 1 + z \) is introduced due to the cosmological dilution of time. Because of the collision, a pair of shocks (i.e., internal shocks) could arise: a forward shock propagating into shell 1 and a reverse shock propagating into shell 4. The shocked regions in shells 1 and 4 are denoted by 2 and 3, respectively.

According to the jump conditions between the two sides of a shock (Blandford & McKee 1976), we can calculate the comoving internal energy densities of the two shocked regions by \( \rho_{2} = (\gamma_{21} - 1)(4\gamma_{21} + 3)n_{1}m_{p}c^{2} \) and \( \rho_{3} = (\gamma_{23} - 1)(4\gamma_{23} + 3)n_{3}m_{p}c^{2} \), where \( \gamma_{21} \) or \( \gamma_{23} \) is the Lorentz factor of region 2 or 3 relative to the unshocked region 1 or 4. The comoving proton number density \( n_{i} \) of unshocked region \( i \) can be calculated by \( n_{i} = L_{k,i}/(4\pi R_{k,i}^{2}\gamma_{i}^{2}m_{p}c^{2}) \) for an isotropic kinetic-energy luminosity \( L_{k,i} \) and a bulk Lorentz factor \( \gamma_{i} > 1 \) of the unshocked shell. The mechanical equilibrium between the two shocked regions requires \( \epsilon_{2} = \epsilon_{3} \), which yields

\[
\frac{(\gamma_{21} - 1)(4\gamma_{21} + 3)}{(\gamma_{23} - 1)(4\gamma_{23} + 3)} = \frac{n_{4}}{n_{1}} = \left( \frac{\gamma_{4}}{\gamma_{1}} \right)^{2}, \tag{1}
\]

where \( L_{k,1} = L_{k,4} = L_{k} \) is supposed. By assuming \( \gamma_{4} \gg \gamma_{1} \), the above equation leads to (Yu & Dai 2009)

\[
\begin{align*}
\gamma_{21} - 1 &= \frac{1}{2} \left( \frac{\gamma_{1}}{\gamma} + \frac{\gamma}{\gamma_{1}} \right) - 1 \approx \xi \lesssim 1, \\
\gamma_{23} &= \frac{1}{2} \left( \frac{\gamma_{4}}{\gamma} + \frac{\gamma_{4}}{\gamma} \right) \approx \frac{\gamma_{4}}{2\gamma} \tag{2}
\end{align*}
\]

where \( \gamma = \gamma_{1}(1 + \sqrt{2\xi}) \) is the Lorentz factor of the shocked regions. This indicates that the reverse shock is relativistic and the forward shock is Newtonian. Moreover, following Dai & Lu (2002), the total number of the electrons swept up by the forward and reverse shocks during a period of \( \delta t \) can be expressed by

\[
N_{2} = 2\sqrt{\frac{2L_{k}\delta t}{(\phi_{1}\gamma_{1}m_{p}c^{2})}} \quad \text{and} \quad N_{3} = L_{k}\delta t/(\phi_{4}\gamma_{4}m_{p}c^{2}), \tag{3}
\]

Both forward and reverse shocks can accelerate particles to high energies and amplify magnetic fields. As usual, we assume that the energies of the hot electrons and magnetic fields are fractions \( \epsilon_{e} \) and \( \epsilon_{B} \) of the total internal energy, respectively. Thus, the strength of the magnetic fields is given

\[
B_{i} = (8\pi\epsilon_{B}e\epsilon_{e})^{1/2} = \frac{1}{\sqrt{2\gamma_{21}^{3}}/\gamma_{21}^{2}y_{i}^{6}} \tag{4}
\]

For the shock-accelerated electrons, a power-law energy distribution, \( dN/d\gamma_{e} \propto \gamma^{-p} \) for \( \gamma_{e} \gg \gamma_{e,m}, \) is assumed. The characteristic random Lorentz factors of these hot electrons in regions 2 and 3 are determined respectively by

\[
\begin{align*}
\gamma_{e,m,2} &= \epsilon_{e}\epsilon_{m}/m_{p} \left( \frac{\gamma_{21} - 1}{\epsilon_{e}} \right) \approx \epsilon_{e}\epsilon_{m}/m_{p}, \tag{5}
\gamma_{e,m,3} &= \epsilon_{e}\epsilon_{m}/m_{p} \left( \frac{\gamma_{23} - 1}{\epsilon_{e}} \right) \approx \epsilon_{e}\epsilon_{m}/m_{p} \tag{5}
\end{align*}
\]

where \( \epsilon_{e} = \epsilon_{e}(p - 2)/(p - 1) \). That \( \gamma_{e,m,3} > \gamma_{e,m,2} \), due to \( \gamma_{2}/\gamma_{4} \gg 2\xi \), indicates that the characteristic energy of the reverse-shocked electrons is much higher than that of the forward-shocked electrons. Therefore, the resulting synchrotron photons emitted by these two types of electrons are expected to peak at two different energy bands and thus two distinct spectral components would be observed as in GRB 080319B. To be specific, the reverse shock is responsible for emission in a relatively high-energy band, while the forward shock contributes to a relatively low-energy component.

In both shocked regions, the hot electrons with energies above \( \gamma_{e,c,i}m_{p}c^{2} \) lose most of their energies during a cooling time \( t_{c,i} \), where the cooling Lorentz factor is determined by \( \gamma_{e,c,i} = \delta t_{i}m_{p}c^{2}/(\gamma_{i}\sigma T_{B}^{0}r_{i,c}) \). The parameter \( r_{i} \), defined as the ratio of the total luminosity to the synchrotron one, is introduced by considering the cooling effect due to the IC emission besides the synchrotron cooling. As pointed out by Ghisellini et al. (2000), the theoretical synchrotron spectrum arising from these electrons, calculated by using the standard assumption that the magnetic field maintains a steady value throughout the shocked region, leads to a spectral slope \( F_{\nu} \propto \nu^{-1/2} \) below \( \sim 100 \text{ keV} \), which is in contradiction to the much harder spectra observed. In order to overcome this problem, Pe‘er & Zhang (2006) suggested that the magnetic field created by a shock could decay on a length scale \( \lambda_{R,i} \) much shorter than the comoving width \( \Delta_{i} \) of the shocked region, i.e., \( \lambda_{R,i} = \Delta_{i}/f_{R,i} \), \( f_{R,i} > 1 \). In other words, the shocked region can be roughly divided into a magnetized part immediately after the shock front and a further unmagnetized part. Under this assumption, the cooling time of the electrons should be determined by the time during which the electrons traverse the magnetized region, i.e., \( t_{c,i} = \delta t_{R,i} \). Although the size of the magnetized region is reduced significantly by the field-decay effect (as found in Section 3.3), this region could be still wide enough for electrons to lose a great part of their energy when they traverse it. In this case, the cooling Lorentz factor \( \gamma_{e,c,i} \) of electrons is not
much higher than $\gamma_{c,m,i}$ so that the radiation efficiency of the electrons is not reduced drastically compared to the case without any magnetic field decay.

2.2. Two-component Synchrotron Emission

With the electron distributions and the magnetic fields described above, we can give the resultant synchrotron spectra using the method developed by Sari et al. (1998). The reference peak energies of the synchrotron spectra of the forward and reverse shocks are taken to be in optical and $\gamma$-ray bands, respectively. Then, the model parameters can be expressed as functions of some observational quantities.

For the electrons in the magnetized reverse-shocked region, two break frequencies of the synchrotron spectrum are given by

$$
\nu_{m,3} = \frac{q_e}{2\pi m_e c \phi_c} \gamma_{c,m,3}^2 B_3 \eta
$$

where $\gamma_{c,m,3}$ are functions of the magnetic field decay. The quantities in the left-hand sides of Equations (6)–(7) are functions of some observational quantities.

The peak flux density of the spectrum at $\nu_p = \min[\nu_{m,3}, \nu_c]$ reads

$$
F_{\nu, p, 3} = \frac{\phi_c}{4\pi d_L^2} \frac{m_e c^2 \sigma_T}{3q_e} f_{B,3} B_3 \eta
$$

where the parameter $f_{B,3}$ is introduced because, at any moment, only a fraction $1/f_{B,3}$ of the total reverse-shocked electrons located at the magnetized region and other electrons in the un-magnetized region do not contribute to the synchrotron emission. The quantities in the left-hand sides of Equations (6)–(8) can be inferred from an observed prompt $\gamma$-ray spectrum, while the right-hand sides are functions of the model parameters. We can therefore solve these equations to find the values of some model parameters,

$$
L_k = 2.5 \times 10^{33} \text{ erg s}^{-1} \gamma_{3,0} \epsilon_{v, -1}^{-1}
\times \left(\frac{d_L^2}{L_{28}} F_{\nu, p, 3, -25} \nu_c^{1/2} \nu_{3,0}^{1/2} \nu_{m,3,20}^{1/2}\right)
\equiv \frac{\nu_{3,0}}{\nu_c} L_{k, 28}
$$

$$
\gamma_4 = 4 \times 10^4 y_{2,5}^{1/4} \nu_{3,0}^{1/4} \nu_{c, -1}^{1/4} \nu_{v, 3,20}^{1/4} \nu_{m,3,20}^{1/4}
\times \left(\delta_{L,28}^{-1/2} d_{L,28}^{-1/2} F_{\nu, p, 3, -25}^{-1/2} \nu_c^{1/2} \nu_{3,0}^{1/2} \nu_{v, 3,20}^{1/2} \nu_{m,3,20}^{1/2}\right)
\quad \equiv \left(\frac{\nu_{3,0}}{\nu_c} L_{k, 28}\right)^{1/2}
$$

and $f_{B,3}$ are strongly dependent on $\gamma$ that can be constrained by optical observations. The fact that $L_k$ is independent of the parameter $f_{B,3}$ indicates that the hypothesis of magnetic field decay does not increase the energy requirement of the model. This is because all the reverse-shocked electrons, when they traverse the tiny magnetized region at different times, have released a great part of their energy to $\gamma$-rays (as indicated by $\nu_{m,3} \sim \nu_{c,3}$) via synchrotron emission. The observed $\gamma$-ray emission is mainly contributed by the emission from this tiny region behind the shock front, no matter whether the other part of the shocked region is magnetized or not.

In order to study the properties of the forward shock, we now focus on a possible low-energy emission, whose spectral information is however not as rich as the $\gamma$-ray component. We calculate the peak frequency of the synchrotron spectrum produced by the forward shock as

$$
\nu_{m,2} = \frac{q_e}{2\pi m_e c \phi_c} \gamma_{c,m,2}^2 B_2 \eta
$$

where $\gamma_{c,m,2}$ is introduced because, at any moment, only a fraction $1/f_{B,2}$ of the total forward-shocked electrons located at the magnetized region and other electrons in the un-magnetized region do not contribute to the synchrotron emission. The quantities in the left-hand sides of Equations (6)–(8) can be inferred from an observed prompt $\gamma$-ray spectrum, while the right-hand sides are functions of the model parameters. We can therefore solve these equations to find the values of some model parameters,

$$
F_{\nu, o, 2} = \frac{\phi_c}{4\pi d_L^2} \frac{m_e c^2 \sigma_T}{3q_e} f_{B,2} B_2 \eta
$$

where $f_{B,2}$ is introduced because, at any moment, only a fraction $1/f_{B,2}$ of the total forward-shocked electrons located at the magnetized region and other electrons in the un-magnetized region do not contribute to the synchrotron emission. The quantities in the left-hand sides of Equations (6)–(8) can be inferred from an observed prompt $\gamma$-ray spectrum, while the right-hand sides are functions of the model parameters. We can therefore solve these equations to find the values of some model parameters,

$$
\gamma_4 = 4 \times 10^4 y_{2,5}^{1/4} \nu_{3,0}^{1/4} \nu_{c, -1}^{1/4} \nu_{v, 3,20}^{1/4} \nu_{m,3,20}^{1/4}
\times \left(\delta_{L,28}^{-1/2} d_{L,28}^{-1/2} F_{\nu, p, 3, -25}^{-1/2} \nu_c^{1/2} \nu_{3,0}^{1/2} \nu_{v, 3,20}^{1/2} \nu_{m,3,20}^{1/2}\right)
\equiv \left(\frac{\nu_{3,0}}{\nu_c} L_{k, 28}\right)^{1/2}
$$

where and hereafter the convention $Q = 10^Q$, $Q_0$ is adopted in cgs units. The quantities in the brackets are basically determined by the observational data and the values of $L_k$, $\gamma_4$, and $f_{B,3}$ are modulated by the remaining free parameters. In particular, $\gamma_4$ is a free parameter which can be constrained by optical observations. The fact that $L_k$ is independent of the parameter $f_{B,2}$ indicates that the hypothesis of magnetic field decay does not increase the energy requirement of the model. This is because all the reverse-shocked electrons, when they traverse the tiny magnetized region at different times, have released a great part of their energy to $\gamma$-rays (as indicated by $\nu_{m,3} \sim \nu_{c,3}$) via synchrotron emission. The observed $\gamma$-ray emission is mainly contributed by the emission from this tiny region behind the shock front, no matter whether the other part of the shocked region is magnetized or not.

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$$
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where $\gamma_{c,m,2}$ is introduced because, at any moment, only a fraction $1/f_{B,2}$ of the total forward-shocked electrons located at the magnetized region and other electrons in the un-magnetized region do not contribute to the synchrotron emission. The quantities in the left-hand sides of Equations (6)–(8) can be inferred from an observed prompt $\gamma$-ray spectrum, while the right-hand sides are functions of the model parameters. We can therefore solve these equations to find the values of some model parameters,

$$
F_{\nu, o, 2} = \frac{\phi_c}{4\pi d_L^2} \frac{m_e c^2 \sigma_T}{3q_e} f_{B,2} B_2 \eta
$$

where $f_{B,2}$ is introduced because, at any moment, only a fraction $1/f_{B,2}$ of the total forward-shocked electrons located at the magnetized region and other electrons in the un-magnetized region do not contribute to the synchrotron emission. The quantities in the left-hand sides of Equations (6)–(8) can be inferred from an observed prompt $\gamma$-ray spectrum, while the right-hand sides are functions of the model parameters. We can therefore solve these equations to find the values of some model parameters,
where the kinetic-energy luminosity $L_{k} = 2.5 \times 10^{52} \text{ erg s}^{-1}$ ($\gamma_{3}/\epsilon_{k}) L_{k} = 5.2$, obtained in Equation (9) has been substituted. The value of $\epsilon$ is mainly determined by the observational data and insensitive to the remaining free parameters $\epsilon_{B}$ and $\epsilon_{e}$. For common GRBs, the value of $\tau_{\gamma o}$ is thought to be very high (e.g., $\sim 10^3$) and thus the value of $\gamma$ could be lower than 100, which leads to hundreds for $\gamma_{4}$. However, for GRB 080319B, $\tau_{\gamma o}$ is deemed to be not larger than unity to ensure a bright optical flash and we suggest $\tau_{\gamma o} = 0.1$ as a reference value hereafter.

3. APPLICATION TO GRB 080319B

GRB 080319B triggered the Swift Burst Alert Telescope (15–350 keV) at $T_{0} = 06:12:49$ UT on 2008 March 19 (Racusin et al. 2008a) and was simultaneously detected by the Konus $\gamma$-ray detector onboard the Wind satellite (20 keV–15 MeV; Golenetskii et al. 2008). The time-averaged Konus–Wind $\gamma$-ray spectrum can be fitted well by the Band function (Band et al. 1993) with a low-energy slope of $0.855^{+0.014}_{-0.0014}$ below the peak of $E_{p} = 675 \pm 22$ keV and a high-energy slope of $-3.59^{+0.32}_{-0.62}$ above the peak (Racusin et al. 2008b). The burst had a peak flux of $F_{p} = (2.26 \pm 0.21) \times 10^{-5}$ erg cm$^{-2}$ s$^{-1}$ and thus the peak isotropic equivalent luminosity was $L_{p,iso} = (1.01 \pm 0.09) \times 10^{53}$ erg s$^{-1}$. Using the values of $F_{p}$ and $E_{p}$, we roughly estimate the peak flux density of the $\gamma$-ray spectrum, $F_{p}/E_{p} \approx 14$ mJy. Compared with the extrapolation of the $\gamma$-ray spectrum to the optical band, the observed flux density of the optical flash ($\sim 20$ Jy) is about 10,000 times higher.

3.1. The Model Parameters

Adopting $z = 0.937$ ($d_L = 1.9 \times 10^{28}$ cm), $\Delta t \sim 3$ s, $F_{V,o} \sim 20$ Jy (at $v_o = 5 \times 10^{14}$ Hz), $F_{V,3} \sim 14$ mJy, $h_{V,m} = 675$ keV and denoting $V_{3} \equiv x \times V_{m,3} \equiv 10^{3}x_{0} \times V_{m,3}$ for GRB 080319B, we derive the model parameters,

$$L_{k} \approx 2 \times 10^{54} \text{ erg s}^{-1} \times \gamma_{3/0}^{1/2} \epsilon_{3/0}^{-1},$$

$$\gamma \approx 400 \gamma_{3/0}^{1/2} \epsilon_{3/0}^{-1/2},$$

$$\gamma_{4} \approx 9 \times 10^{3} \times \gamma_{3/0}^{3/2} \epsilon_{3/0}^{-3/2},$$

$$f_{B,2} \approx 700 \times \gamma_{3/0}^{3/2} \epsilon_{3/0}^{-1/2},$$

$$f_{B,3} \approx 7 \times 10^{3} \times \gamma_{3/0}^{3/2} \epsilon_{3/0}^{-1/2},$$

where a fiducial value of unity is assumed for $y_{3}$, which will be proved in Section 3.2. Besides a constraint by the maximum allowed equipartition value ($\epsilon_{e} \lesssim 0.3$ and $\epsilon_{B} \lesssim 0.3$), the remaining free parameters $\epsilon_{B}$ and $\epsilon_{e}$ satisfy $\epsilon_{e} < 0.09 \gamma_{3/0}^{2/3}$ given $\gamma > \gamma_{3/0}$. The upper limit of $\epsilon_{e}$ is insensitive to the value of $\epsilon_{B}$ (strictly, with a decrease of $\epsilon_{B}$, the upper limit of $\epsilon_{e}$ increases slightly). Taking $\epsilon_{e} < 0.09 \gamma_{3/0}^{2/3}$ as a conservative estimate, we find

1. $L_{k} \gtrsim 2 \times 10^{54} \text{ erg s}^{-1}$. This is a natural result due to the high observed $\gamma$-ray luminosity ($\sim 10^{53}$ erg s$^{-1}$) of GRB 080319B. The MeV $\gamma$-ray radiation of the reverse shock can be estimated by $\eta \approx 0.05 \gamma_{3/0}^{-2} \epsilon_{3/0}^{-1}$. We next calculate the total isotropic-equivalent energy release of GRB 080319B.

$$E_{k,iso} = 2E_{k,iso}/\eta \approx 5 \times 10^{55} \times \gamma_{3/0}^{-1/2} \epsilon_{3/0}^{-1} \text{ erg},$$

where a factor of 2 is introduced by considering a similar amount of energy carried by the forward shocks. Using a very small jet angle $\theta_{j} = 0\,\,^\circ$ that is found by Racusin et al. (2008b), we get the beaming-corrected energy release of GRB 080319B.

$$E_{k,jet} = E_{k,iso} \gamma_{3/0}^{2}/2 \sim (3 \times 10^{50}) \times \gamma_{3/0}^{-1} \epsilon_{3/0}^{-1} \text{ erg},$$

which is a typical value for common GRBs.

2. $\gamma \sim 400$. This is a typical value for the Lorentz factor of merged GRB ejecta after internal shocks. Thus, we can obtain the internal shock radius for GRB 080319B,

$$R_{i} \approx 2\gamma_{3}^{2} \times \delta t \phi_{c}^{-1} \approx 10^{16} \delta t_{60} \phi_{c,3}^{-1} \gamma_{12,6}^{2} \text{ cm},$$

where a wind-like circumburst medium ($\rho \approx A r^{-2}$ with $A = 5 \times 10^{11} \text{ g cm}^{-3} A_{s}$) is assumed. Moreover, as claimed by Racusin et al. (2008b), a tenuous wind with an upper limit of $A_{s} \sim 0.03$ is required by the afterglow data of GRB 080319B. In this case, the deceleration radius is larger. Therefore, we conclude that the deceleration of the shells can be ignored at the times where internal shocks among the shells occur.

3. $\gamma_{4} \gtrsim 10^{3}$. This high Lorentz factor is allowable for acceleration of an initial fireball with very low baryon contamination (Piran 1999).

4. $f_{B,3} \gtrsim 10 f_{B,2}$. Our constraint on $f_{B,3}$, which is much larger than unity and much less than $\Delta/\lambda_{s}$ ($\lambda_{s}$ is the plasma skin depth), is consistent with that found by Pe’er & Zhang (2006) for other GRBs. However, the physical underpinning of these values of $f_{B}$ is unknown. So the difference in $f_{B}$ for different shocked regions lacks a reasonable physical explanation and it might be related to different shock strengths of the forward and reverse shocks.

From the above discussion, we can see that most of the inferred model parameters are reasonable and acceptable even in common GRBs, except for an unusually high variability of Lorentz factors denoted by $\gamma_{4}/\gamma_{1} \sim 300$. Although the $\gamma$-ray emission of GRB 080319B seems to be not unusual, the relatively high value of $\Delta t \sim 3$ s (vs. $\sim 10$ ms for common GRBs) implies an unusually large internal shock radius, which ensures the synchrotron self-absorption frequency below the optical band and reduces the magnetic field strength. Therefore, in order to produce sufficiently energetic $\gamma$-ray photons, it is necessary to invoke some highly relativistic internal shocks that require high variability of Lorentz factors in the fireball. Since the GRB central engine is far from being thoroughly understood, it is difficult to demonstrate whether the central engine can produce such a drastically varying fireball or not, but some possible origins can be still imagined. In the collapsar model, for example, a shell passing through the envelope of a progenitor star should sweep up and clear away the envelope material, leaving a channel behind the shell. A following shell will pass through this clear channel and thus have a very low baryon contamination. This might lead to a highly relativistic ($\gamma_{shell} \sim 10^{5}$) shell. However, due to lateral diffusion of the channel wall, this channel will be possibly contaminated by baryons again some time
result, relatively slow and rapid shells are generated alternately. In reality, this process is unlikely to be so regular because a slow/rapid shell could be followed by another slow/rapid shell. Therefore, the temporally correlated optical and $\gamma$-ray emissions from this process could be polluted by the emission due to the collisions between slow–slow or rapid–rapid shells. This along with other effects (e.g., different shock-crossing times of shells and so on, see Section 4) may lead to the fact that the prompt optical and $\gamma$-ray emissions are not correlated finely, as observed in GRB 080319B.

To summarize, if some unknown physical processes of the central engine can give rise to a fireball in which the Lorentz factors and densities vary drastically in a variability timescale of a few seconds, our internal shock-produced two-component synchrotron emission scenario may account for some fundamental features of the optical and MeV $\gamma$-ray emissions of GRB 080319B.

3.2. Inverse Compton Emission

We denote $\tau_i = \sigma_T N_i / (4 \pi R^2 f_{B,i})$ and $\tau_j^p = \sigma_T N_j / (4 \pi R^2)$ as the Thomson optical depth of the magnetized and unmagnetized regions, respectively. The SSC emission from the two magnetized regions is considered first. Following Sari & Esin (2001) for $p = 2.5$, we estimate the two break energies of the SSC spectrum contributed by the forward shock by

$$h_{\nu, m, 2}^{\text{SSC}} = \gamma_{\nu, m, 2}^2 \hbar \nu_{m, 2} = 1.5 \text{ keV} \left( \frac{3^{16/3}}{3^{1/2} / 3/8 \gamma_{E, B, -1}^{1e, -1}} \right)^{29/3},$$

$$h_{\nu, r, 2}^{\text{SSC}} = \gamma_{\nu, r, 2}^2 \hbar \nu_{r, 2} = 70 \text{ GeV} \left( \frac{3^{16/3}}{3^{1/2} / 3/8 \gamma_{E, B, -1}^{1e, -1}} \right)^{29/3}. $$

The peak flux at $\nu_{r, 2}^{\text{SSC}}$ can be estimated by

$$\nu_{r, 2}^{\text{SSC}} \approx \nu_{r, 2}^{\text{SSC}} \gamma_{\nu, r, 2}^2 F_{\nu, r, 2} \left( \frac{\nu_{r, 2}^{\text{SSC}}}{\nu_{m, 2}^{\text{SSC}}} \right)^{(p-1)/2} \sim 2 \times 10^{-8} \text{ erg cm}^{-2} \text{ s}^{-1}. $$

For the reverse shock, as the SSC peak enters the Klein–Nishina regime, the real peak energy can be determined by (Gupta & Zhang 2007; Fragile et al. 2004)

$$h_{\nu, K, 3}^{\text{SSC}} = \frac{\nu_{K, 3}^2 \mu_{e, d}^2}{h \nu_{m, 3}} = 60 \text{ GeV} \left( \frac{3^{1/2} / 3/8 \gamma_{E, B, -1}^{1e, -1}}{3^{1/2} / 3/8 \gamma_{E, B, -1}^{1e, -1}} \right)^{29/3},$$

at which a negligible flux is found from

$$\nu_{K, 3}^{\text{SSC}} \approx \nu_{K, 3}^{\text{SSC}} \frac{F_{\nu, K, 3} \nu_{\nu, K, 3}^{\text{SSC}}}{2 \gamma_{\nu, r, 3}^2 \nu_{r, 3}^{\text{SSC}}} \left( \frac{2 \gamma_{\nu, r, 3}^2 \nu_{r, 3}^{\text{SSC}}}{\nu_{K, 3}^{\text{SSC}}} \right)^{1/3} \sim 8 \times 10^{-11} \text{ erg cm}^{-2} \text{ s}^{-1}. $$

According to the definition of parameter $y_i$, we obtain

$$y_2 = 1 + \frac{L_2^{\text{SSC}}}{L_2^{\text{opt}}} \lesssim 1 + \frac{1}{\left[ F_{\nu, l, 2}^{\text{SSC}} / F_{\nu, l, 2}^{\text{opt}} \right]} \approx 1.2, $$

$$y_3 = 1 + \frac{L_3^{\text{SSC}}}{L_3^{\text{opt}}} \lesssim 1 + \frac{\left[ F_{\nu, l, 3}^{\text{SSC}} / F_{\nu, l, 3}^{\text{opt}} \right]}{F_{\nu, l, 3}^{\text{MeV}}} \approx 1, $$

where $[F_{\nu, l}^{\text{opt}}] \approx 10^{-7} \text{ erg cm}^{-2} \text{ s}^{-1}$ and $[F_{\nu, l}^{\text{MeV}}] \approx 2.3 \times 10^{-5} \text{ erg cm}^{-2} \text{ s}^{-1}$. This indicates that the SSC emission is relatively unimportant and the estimations in Equations (19)–(32) under the assumption, $y_2 \sim y_3 \sim 1$, are self-consistent.

Although a part of the electron energy can be released via synchrotron and SSC emissions when the electrons traverse the magnetized regions, about half energy of the reverse-shocked electrons and almost all energy of the forward-shocked electrons are still held by the electrons when they enter into the unmagnetized regions. This remaining energy can be no longer released via synchrotron radiation because of the lack of magnetic fields, but can via external inverse Compton (EIC) scattering due to the existence of the radiation fields. Determined by this EIC cooling, the cooling Lorentz factor of the electrons in both the unmagnetized regions reads

$$\gamma_{e, c} = \frac{m_e c \phi_{e, c}}{\frac{3}{2} \sigma_T Y_{\gamma} \delta t},$$

where the radiation energy density contributed by the synchrotron radiation from the two magnetized regions can be calculated by $u_{\nu} = (y_2 + y_3 - 2)(B^2/8\pi)$, then we have

$$\gamma_{e, c} = \frac{6 \pi m_e c \phi_{e, c}}{Y_{\gamma} \sigma_T B^2 \gamma \delta t} \approx 15 \gamma_0^{3/2} \gamma_{e, c}^{-1},$$

where $\gamma_{e, c} < (y_{e, m, 2}, y_{e, m, 3})$. For an EIC spectrum produced by upscattering the seed photons from magnetized region $j$ by the electrons in unmagnetized region $i$, we calculate its two characteristic break frequencies by

$$\nu_{L, i}^{(j)} = 2 \gamma_{e, L, i} \nu_{L, j}^{(j)},$$

$$\nu_{H, i}^{(j)} = 2 \gamma_{e, H, i} \nu_{H, j}^{(j)}.$$  

and estimate its peak flux at the peak frequency $\nu_{H, i}^{(j)}$ roughly by

$$[F_{\nu, i}^{(j)}]_{H, i} \approx \nu_{H, i}^{(j)} \nu_{H, i}^{(j)} F_{\nu, H, j}^{(j)} (\nu_{H, i}^{(j)} / \nu_{L, i}^{(j)})^{(p-1)/2},$$

where the subscript $L$ represents the low break frequency of the seed photons and the low break Lorentz factor of the target electrons, while $H$ represents the high ones. Considering the two synchrotron components for seed photons and the two population unmagnetized electrons, four EIC components are expected:

1. For $i = 3$ and $j = 2$, we have

$$h_{\nu, l, 3}^{(2)} = 2 \gamma_{e, l, 2} \hbar \nu_{m, 2} = 0.5 \text{ keV} \left[ \frac{2^{1/2} / 3^{1/2} / 3^{1/2} / 3^{1/2}}{3^{1/2} / 3^{1/2} / 3^{1/2} / 3^{1/2}} \right]^{29/3} \gamma_{e, c}^{-1},$$

$$h_{\nu, H, 3}^{(2)} = 2 \gamma_{e, m, 2} \hbar \nu_{m, 2} \approx 6 \text{ TeV} \left[ \frac{2^{1/2} / 3^{1/2} / 3^{1/2} / 3^{1/2}}{3^{1/2} / 3^{1/2} / 3^{1/2} / 3^{1/2}} \right]^{29/3} \gamma_{e, c}^{-1}. $$

Although the peak frequency $\gamma_{e, H, 3}^{(2)}$ is just around the Klein–Nishina break determined by $h_{\nu, K, 3}^{(2)} = \gamma^2 m_e^2 c^4 / (h \nu_{m, 2}) \approx 6 \text{ TeV}$, we may still derive a relatively low flux from Equation (39),

$$[F_{\nu, i}^{(j)}]_{H, i} \approx 8 \times 10^{-8} \text{ erg cm}^{-2} \text{ s}^{-1}. $$
2. For $i = 3$ and $j = 3$, we have

$$h_{L,3}^{(3)} = 2y_{e,c} T_2 x_3 h_{e,c}$$

$$= 0.3 \text{ GeV} x_0^{5/16} y_{-0.7}^{3/8} x_3^{3/8} \epsilon_{e,c}^{-1} \epsilon_{-1}^{1/2}.$$  (43)

$$h_{H,3}^{(3)} = 2y_{e,m}^2 x_3$$

$$= 0.6 \text{ BeV} x_0^{-3/16} y_{-0.7}^{3/8} x_3^{3/8} \epsilon_{e,c}^{-1} \epsilon_{-1}.$$  (44)

The peak frequency $\nu_{H,3}^{(3)}$ is however higher than the Klein–Nishina break frequency, $h_{K,N,3}^{(3)} = \gamma^2 m_e c^4 / (h v_{m,3}) \approx 60 \text{ GeV}$, and thus the real peak flux can be calculated at $\nu_{K,N,3}^{(3)}$ to be

$$[vF]\nu_{K,N,3}^{(3)} \sim 2 \times 10^{-7} \text{ erg cm}^{-2} \text{ s}^{-1}.$$  (45)

Summarizing the above two cases, the EIC process of the reverse-shocked electrons produces a TeV and a GeV emission component, both of which are much weaker than the observed synchrotron MeV emission with a flux of $\sim 2.3 \times 10^{-5} \text{ erg cm}^{-2} \text{ s}^{-1}$.

3. For $i = 2$ and $j = 2$, we have

$$h_{L,2}^{(2)} = 2y_{e,c} T_2 x_2 h_{e,c}$$

$$= 0.5 \text{ keV} x_0^{-1/2} y_{-0.7}^{1/2} x_2^{1/2} \epsilon_{e,c}^{-1} \epsilon_{-1}^{3}.$$  (46)

$$h_{H,2}^{(2)} = 2y_{e,m}^2 x_2$$

$$= 10 \text{ MeV} x_0^{35/32} y_{-0.7}^{35/16} x_2^{35/16} \epsilon_{e,c}^{45/16}.$$  (47)

and the peak flux at $\nu_{H,2}^{(2)}$

$$[vF]\nu_{H,2}^{(2)} \sim 7 \times 10^{-7} \text{ erg cm}^{-2} \text{ s}^{-1}.$$  (48)

Due to the stronger synchrotron MeV emission, this relatively weaker MeV emission component is likely to be covered up.

4. For $i = 2$ and $j = 3$, we have

$$h_{L,3}^{(3)} = 2y_{e,c} T_2 x_3 h_{e,c}$$

$$= 0.3 \text{ GeV} x_0^{5/16} y_{-0.7}^{3/8} x_3^{3/8} \epsilon_{e,c}^{-1} \epsilon_{-1}^{1/2}.$$  (49)

$$h_{H,3}^{(3)} = 2y_{e,m}^2 x_3$$

$$= 0.9 \text{ GeV} \epsilon_{e,c}^2 \epsilon_{-1}.$$  (50)

Distinct from the emission components discussed in the previous three cases, the flux at $\nu_{H,2,3}^{(3)}$

$$[vF]\nu_{H,2,3}^{(3)} \sim 2 \times 10^{-7} \text{ erg cm}^{-2} \text{ s}^{-1},$$  (51)

indicates a strong GeV emission that is as strong as the synchrotron MeV emission. This means that the energy of the forward-shocked electrons would be mainly released by upscattering the synchrotron MeV $\gamma$-ray photons from the reverse shock, while the reverse-shocked electrons lose a great part of their energy via synchrotron cooling directly. Due to the strong emission within the energy regime from sub-GeV to GeV, electron–positron pairs might be produced by collisions between the sub-GeV and GeV photons. According to the results obtained above, we can give approximately an upper limit luminosity of $L_{\text{lim}} \sim 10^{53} \text{ erg s}^{-1}$ for the sub-GeV ($\epsilon_{\gamma} \sim 0.1 \text{ GeV}$) emission. Then, the optical depth due to pair production interactions can be roughly estimated as

$$\tau_{\gamma\gamma} \lesssim \frac{3}{16} \frac{L_{\text{lim}}}{4\pi R^2 c^2 \epsilon_{\gamma}} \frac{R}{\gamma} \sim 0.1 L_{\text{lim},33} \gamma_{2,6}^{-2} R_{16}^{-1},$$  (52)

which indicates that the pair production effect is not significant. For simplicity, the further contribution from the secondary electrons is ignored here.

To summarize, the contributions by the SSC and EIC emission to the observed optical and MeV $\gamma$-ray emissions are insignificant and the two synchrotron components are dominant in the observed bands for GRB 080319B. In contrast, some higher energy emission components would be produced by the EIC process. Although most of these components are weak, the flux of the strong GeV emission can reach as high as $\sim 10^{-5} \text{ erg cm}^{-2} \text{ s}^{-1}$, which is comparable to the observed MeV one.

4. SUMMARY AND CONCLUSIONS

In the popular internal shock model for the prompt emission of GRBs, paired forward and reverse shocks are produced by collisions between some relativistic shells with different Lorentz factors. In this paper, we have considered this model in an unusual situation where a bimodal distribution of the shell Lorentz factors exists. As a result, the Lorentz factors of the forward and reverse shocks are quite different (i.e., the forward shock is Newtonian and the reverse shock is relativistic) and the resulting two-component synchrotron emission is expected to provide a new scenario for some seldom GRBs that have two spectral peaks in the prompt emission.

As an example, we compare our scenario with the recently observed naked-eye GRB 080319B and constrain the model parameters by fitting the observations. We find that, on one hand, the optical and MeV $\gamma$-ray fluxes of this unique GRB could be explained in our two-component synchrotron emission scenario, if some unknown physical processes of the central engine (e.g., the picture described in Section 2.1) can give rise to a fireball where the Lorentz factors of $\sim 400$ and $\sim 10^5$ appear alternately in a variability timescale of a few seconds. On the other hand, although the internal shock-produced emission can roughly account for the fluctuating structure of the light curves and the mild temporal correlation between the optical and $\gamma$-ray emissions, it is still difficult to explain clearly why the observed optical emission varies relatively slower than the $\gamma$-ray emission. We speculate that this difference between the light curve variabilities may be due to a complicated distribution of Lorentz factors in the fireball, an inhomogeneous structure of each fireball shell, different shock-crossing times, and other more realistic properties of the system. So, a more detailed simulation is required to improve our present model.

Finally, for high energy emission, the synchrotron plus SSC scenario suggested by Kumar & Pannaitescu (2009) and Racusin et al. (2008b) predicts significant GeV $\gamma$-ray emission by considering the second-order IC-scattering, the flux of which is about 10 times higher than the observed MeV one. In contrast, our model predicts a relatively weaker GeV component, whose
flux is lower than or at most comparable to that of the synchrotron MeV emission. Therefore, future observations for high energy counterparts of GRBs by the *Fermi Space Telescope* are expected to be able to discriminate these two models.

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