KLEIN–NISHINA EFFECTS ON THE HIGH-ENERGY AFTERGLOW EMISSION OF GAMMA-RAY BURSTS

XIANG-YU WANG1, HAO-NING HE1, ZHUO LI2,3, XUE-FENG WU4,5, AND ZI-GAO DAI1

1 Department of Astronomy, Nanjing University, Nanjing 210093, China
2 Department of Astronomy, Peking University, Beijing 100871, China
3 Kavli Institute for Astronomy and Astrophysics, Peking University, Beijing 100871, China
4 Department of Astronomy and Astrophysics, Pennsylvania State University, 525 Davey Lab, University Park, PA 16802, USA
5 Purple Mountain Observatory, Chinese Academy of Sciences, Nanjing 210008, China

Received 2009 November 23; accepted 2010 February 11; published 2010 March 11

ABSTRACT

Extended high-energy ($\gtrsim$100 MeV) gamma-ray emission that lasts much longer than the prompt sub-MeV emission has been detected from quite a few gamma-ray bursts (GRBs) by Fermi–Large Area Telescope (LAT) recently. A plausible scenario is that this emission is the afterglow synchrotron emission produced by electrons accelerated in the forward shocks. In this scenario, the electrons that produce synchrotron high-energy emission also undergo inverse Compton (IC) loss and the IC scattering with the synchrotron photons should be in the Klein–Nishina (KN) regime. Here we study effects of the KN scattering on the high-energy synchrotron afterglow emission. We find that at early times the KN suppression effect on those electrons that produce the high-energy emission is usually strong and therefore their IC loss is small with a Compton parameter $Y \lesssim 1$ for a wide range of parameter space. This leads to a relatively bright synchrotron afterglow at high energies that can be detected by Fermi–LAT. As the KN suppression effect weakens with time, the IC loss increases and could dominate over the synchrotron loss in some parameter spaces. This will lead to a faster temporal decay of the high-energy synchrotron emission than is predicted by the standard synchrotron model, which may explain the observed rapid decay of the early high-energy gamma-ray emission in GRB090510 and GRB090902B.

Key words: gamma-ray burst: general – radiation mechanisms: nonthermal

Online-only material: color figures

1. INTRODUCTION

With the launch of the Fermi satellite, a new feature of gamma-ray bursts (GRBs) at high energies has been established, i.e., GRBs show extended high-energy ($\gtrsim$100 MeV) emission which lasts much longer than the prompt phase. This extended emission has been seen in both long and short GRBs and the flux usually decays monotonically with time after the initial peak. In some cases (e.g., GRB090510 and GRB090902B), the temporal decay is a simple power-law decay with a slope ranging from $-1.3$ to $-1.5$ (Abdo et al. 2009a, 2009b; Ghirlanda et al. 2010; De Pasquale et al. 2010). One of the proposed models for such emission is the hadronic cascade emission model, in which the high-energy photons produced by the accelerated ultra-high energy protons cannot escape the soft photon field and a cascade is induced (Abdo et al. 2009b). This model has been applied to the extended emission of GRB941017 (Dermer & Atoyan 2004), but whether it can explain the simple power-law decay of the Fermi–Large Area Telescope (LAT) bursts is unknown. The long-lived behavior and not very rapid decay of the high-energy emission from GRB090510 and GRB090902B cannot be easily explained by the reverse shock emission model either (Wang et al. 2001a, 2001b). On the other hand, the simple power-law decay with a modestly large slope is reminiscent of the afterglow emission. The self-inverse-Compton (IC) emission of the afterglow has long been thought to produce a high-energy component (e.g., Zhang & Mészáros 2001; Sari & Esin 2001; Fan et al. 2008; Gou & Mészáros 2007), but the light curve is expected to rise initially and start to decay minutes to hours after the burst. Kumar & Barniol Duran (2009a) proposed that the extended high-energy emission from GRB080916C is due to afterglow synchrotron emission. This mechanism has also been proposed to explain the extended high-energy emission from GRB090916C (Gao et al. 2009; Ghirlanda et al. 2010; Ghisellini et al. 2009; De Pasquale et al. 2010) and GRB090902B (Kumar & Barniol Duran 2009b).

In the latter synchrotron afterglow scenario, the high-energy extended emission is produced by the electrons in the forward shock via synchrotron emission. The shock-accelerated electrons are usually assumed to have a power-law form in energy distribution, i.e., $dN_e/d\gamma_e \propto \gamma_e^{-\alpha}$, where $\gamma_e$ is the Lorentz factor of electrons. These electrons also suffer IC loss by scattering synchrotron photons. Due to the fact that the scattering between large $\gamma_e$ electrons and the synchrotron photons could enter the Klein–Nishina (KN) scattering regime, higher energy electrons may suffer smaller IC loss and as a consequence, their synchrotron emission is stronger. Since the Lorentz factor of the electrons producing the high-energy afterglow emission are usually large, the KN scattering effect must be taken into account when one calculates the synchrotron high-energy afterglow emission. In this paper, we study the effect of the KN scattering on the high-energy afterglow emission of GRBs. Recently, Wang et al. (2009) studied the KN effect on the prompt emission spectrum of GRBs and Nakar et al. (2009) studied the KN effect on the optically thin synchrotron and synchrotron self-Compton spectrum in general. In this paper, we focus on the early high-energy afterglow emission and confront the theoretical results with the high-energy afterglow observations by Fermi–LAT.

The paper is organized as follows. First, we study how the KN scattering affects the electron distribution in the forward

---

6 The decay of the high-energy emission in GRB080825C is steeper than $r^{-1.7}$ (Abdo et al. 2009c), which is consistent with the synchrotron self-IC scattering emission from the reverse shock or cross IC scatterings between the reverse shock and forward shock (Wang et al. 2001a, 2001b, 2005; Granot & Guetta 2003; Pe’er & Waxman 2004).
the synchrotron cooling is considered ($\gamma_{e,\text{syn}}$ is the cooling Lorentz factor of electrons when only the synchrotron cooling is considered; see, e.g., Sari et al. 1998). Therefore, the number density of electrons of $\gamma_*$ is a factor of $1 + Y(\gamma_*)$ lower than that in the case where only the synchrotron cooling is considered. As a result, the synchrotron luminosity produced by $\gamma_*$ electrons is correspondingly reduced by the same factor. In the fast-cooling case, when the IC scatterings of electrons of $\gamma_* \lesssim \gamma_m$ with synchrotron photons are in the Thomson scattering regime (the case of Section 3.2.2), one can also obtain $N(\gamma_*) = N_{\text{syn}}(\gamma_*)/[1 + Y(\gamma_*)]$, so the synchrotron luminosity is also reduced by a factor of $1 + Y(\gamma_*)$.

3. KN EFFECT ON THE COMPTON PARAMETERS

Now we derive $Y(\gamma_*)$. As the electron distribution is different in the fast and slow cooling cases, we divide the following analysis into these two different cases.

3.1. The Slow-cooling Case

Whether the afterglow emission belongs to the slow-cooling or fast-cooling case depends on shock microphysics parameters (i.e., the magnetic field equipartition factor $\epsilon_B$ and electron energy equipartition factor $\epsilon_e$) and other parameters such as the burst energy $E$ and the circumburst density $n$. Among these parameters, the magnetic field equipartition factor $\epsilon_B$ is the most poorly known. The circumburst density $n$ depends on the burst environment and may range from $10^{-5}$ cm$^{-3}$ to 10 cm$^{-3}$ (e.g., Kumar & Panaitescu 2001).

The condition for slow cooling is

$$n^{−1/2} < 400[1 + Y(\gamma_*)]^{-1} \epsilon_f^{-1} \epsilon_e^{-1} E_{54}^{−1/2} (1 + z)^{−1/2},$$

where $\epsilon_f \equiv (6p − 2)/(p − 1)$, $p$ is the power-law index of the electron energy distribution ($p = 2.2$ has been used in the following calculations), $E_{54}$ is the energy in units of $10^{54}$ ergs (hereafter we use the cgs units and denotation $Q_0 \equiv Q/10^{54}$), and $z$ is the redshift. The cooling Lorentz factor and the minimum Lorentz factor of electrons in forward shocks are given by

$$\gamma_c = 10^7 [1 + Y(\gamma_*)]^{-1} E_{54}^{−3/8} n_{−1}^{−5/8} \epsilon_e^{−1/8} \epsilon_B^{−5/8} (1 + z)^{−1/8}$$

and

$$\gamma_m = 2.5 \times 10^{4} \epsilon_f \epsilon_e E_{54}^{1/8} n_{−1}^{−3/8} \epsilon_B^{−1/8} (1 + z)^{3/8},$$

respectively, where $Y(\gamma_*)$ is the Compton parameter of the electrons of energy $\gamma_c$. The cooling frequency and minimum frequency of electrons corresponding to $\gamma_c$ and $\gamma_m$ are, respectively,

$$\nu_c = 8 \times 10^{12} [1 + Y(\gamma_*)]^{-2} \epsilon_f^{−3/2} E_{54}^{−3/2} n_{−1}^{−1/2} (1 + z)^{−1/2} \text{ Hz}$$

and

$$\nu_m = 5 \times 10^{17} \epsilon_f^{−3/2} E_{54}^{1/2} n_{−1}^{−3/2} (1 + z)^{1/2} \text{ Hz},$$

In the slow-cooling case, the synchrotron luminosity is dominated by $\gamma_*$ electrons and the ratio of the SSC luminosity to synchrotron luminosity is approximately given by $Y(\gamma_*)$. Depending on the location of $\nu_{\text{KN}}(\gamma_*)$, $U_{\text{pol}}[\nu < \nu_{\text{KN}}(\gamma_*)]$ is proportional to $\nu_{\text{KN}}(\gamma_*)^{3−p/2}$ or $\nu_{\text{KN}}(\gamma_*)^{p/3}$. So the value of
$Y(\gamma_c)$ can be obtained from

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \left(\frac{\gamma_c}{\gamma_m}\right)^{2-p},$$

and

$$\frac{v_{KN}(\gamma_c)}{v_c} = 7.5 \times 10^{-8} \left(1 + Y(\gamma_c)\right)^3 E_{54} n_{-1}^{-5/3} (1 + z)^{-1/3}$$

Since we are interested in the high-energy afterglow emission, we also need to know the ratio of $v_{KN}(\gamma_c)$ to $v_m$, which is

$$\frac{v_{KN}(\gamma_c)}{v_m} = \left(\frac{v_c}{v_m}\right)^{1/2} \frac{v_{KN}(\gamma_c)}{v_m}$$

$$= 2.8 \times 10^{-7} f_p^{-2} e^{-2} e^{-1} \epsilon_{B,-5}^{1/2} n_{-1}^{5/6} (1 + z)^{-1/2}.$$ (14)

According to the relations among $v_{KN}(\gamma_c)$, $v_m$ and $v_c$, we divide the discussion into three cases.

### 3.1.1. Case I: $v_{KN}(\gamma_c) < v_m < v_c$

This case typically happens at early times for the reference parameter values we used. Equation (11) can be simplified as

$$Y(\gamma_c)[1 + Y(\gamma_c)] = 0.09 \left(1 + Y(\gamma_c)\right)^{7/3}$$

while for $Y(\gamma_c) \gg 1$, the value of $Y(\gamma_c)$ can be obtained only numerically.

One can also obtain the Compton parameter for those electrons that produce high-energy afterglow emission with frequency $\nu_c$. Since usually $h \nu_c = 100$ MeV > $h \nu_c$, we only discuss the case of $v_{KN}(\gamma_c) < v_{KN}(\gamma_c)$ below. For $v_{KN}(\gamma_c) < v_{KN}(\gamma_c) < v_m$, we have

$$Y(\gamma_c) = \frac{v_{KN}(\gamma_c)}{v_c} \left(\frac{\gamma_c}{\gamma_m}\right)^{4/3} \left(\frac{\gamma_m}{\gamma_c}\right)^{-4/3} = Y(\gamma_c) \left(\frac{\gamma_c}{\gamma_m}\right)^{2/3}$$

$$= 0.3 f_p^{-5/3} e^{-2} e_{-1}^{1/2} e_{1/3}^{1/3} n_{-1}^{1/6} (1 + z)^{1/2}.$$ (17)

### 3.1.2. Case II: $v_m < v_{KN}(\gamma_c) < v_c$

As $v_{KN}(\gamma_c)/v_m$ increases with time, it is likely that $v_{KN}(\gamma_c) > v_m$ at later times. In this case, we have

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \frac{\gamma_c}{\gamma_m} \left(\frac{\gamma_m}{\gamma_c}\right)^{2-p} = 1.2 \left[1 + Y(\gamma_c)\right]^{3-p/2} f_p^{-2} e_{-1} e_{1/3} B_{-5}^{-6} E_{54} n_{-1}^{-1/2} (1 + z)^{1/2}. \quad (18)$$

### 3.1.3. Case III: $v_{KN}(\gamma_c) > v_c$

In this case, the KN effect on $\gamma_c$ electrons is not important and the Compton parameter is given by

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \left(\frac{\gamma_c}{\gamma_m}\right)^{2-p} = 2 f_p^{-5/3} e^{-2} e_{-1} e_{1/3} B_{-5}^{-6} E_{54} n_{-1}^{1/2} (1 + z)^{1/2}.$$ (19)

### 3.1.2. Case II: $v_m < v_{KN}(\gamma_c) < v_c$

As $v_{KN}(\gamma_c)/v_m$ increases with time, it is likely that $v_{KN}(\gamma_c) > v_m$ at later times. In this case, we have

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \frac{\gamma_c}{\gamma_m} \left(\frac{\gamma_m}{\gamma_c}\right)^{2-p} = 1.2 \left[1 + Y(\gamma_c)\right]^{3-p/2} f_p^{-2} e_{-1} e_{1/3} B_{-5}^{-6} E_{54} n_{-1}^{-1/2} (1 + z)^{1/2}. \quad (18)$$

### 3.1.3. Case III: $v_{KN}(\gamma_c) > v_c$

In this case, the KN effect on $\gamma_c$ electrons is not important and the Compton parameter is given by

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \frac{\gamma_c}{\gamma_m} \left(\frac{\gamma_m}{\gamma_c}\right)^{2-p} = 2 f_p^{-5/3} e^{-2} e_{-1} e_{1/3} B_{-5}^{-6} E_{54} n_{-1}^{1/2} (1 + z)^{1/2}.$$ (19)

### 3.1.3. Case III: $v_{KN}(\gamma_c) > v_c$

In this case, the KN effect on $\gamma_c$ electrons is not important and the Compton parameter is given by

$$Y(\gamma_c)[1 + Y(\gamma_c)] = \frac{\gamma_c}{\gamma_m} \left(\frac{\gamma_m}{\gamma_c}\right)^{2-p} = 2 f_p^{-5/3} e^{-2} e_{-1} e_{1/3} B_{-5}^{-6} E_{54} n_{-1}^{1/2} (1 + z)^{1/2}.$$ (19)
3. Case IIIc: \( v_m < v_e < v_{\text{KN}}(\gamma_e) < v_{\text{KN}}(\gamma). \) In this case,

\[
Y(\gamma) = Y(\gamma_e).
\]  

(24)

3.2. The Fast-cooling Case

The condition for the fast-cooling case is

\[
n_e^{-1/2} \epsilon_{B,-2} \gtrsim 0.4[1 + Y(\gamma_e)]^{-1} f_p^{-1} \epsilon_{e,-1} E_{54}^{-1/2} y_0^{-1/2} (1 + z)^{-1/2}.
\]  

(25)

Below we use a larger \( \epsilon_B \) as the reference value for the fast-cooling case. In this case, the synchrotron radiation is dominated by \( \gamma_m \) electrons and the critical frequency of interest is

\[
v_{\text{KN}}(\gamma_m) = 3.7 \times 10^{18} f_p^{-1} \epsilon_{e,-1}^{-1} \text{Hz}.
\]  

(26)

Similarly, one can find the ratios of \( v_{\text{KN}}(\gamma_m) \) to two characteristic frequencies,

\[
\frac{v_{\text{KN}}(\gamma_m)}{v_m} = 0.24 f_p^{-3} \epsilon_{e,-1} E_{54}^{-1/2} \epsilon_{B,-2}^{-1} y_0^{-1/2} (1 + z)^{-1/2},
\]  

(27)

\[
\frac{v_{\text{KN}}(\gamma_m)}{v_e} = 1.5 [1 + Y(\gamma_e)]^{1/2} f_p^{-1}
\times \epsilon_{e,-1}^{-1} E_{54}^{-1/2} \epsilon_{B,-2}^{-1/2} y_0^{-1/2} (1 + z)^{-1/2}.
\]  

(28)

Below we divide the discussion into two cases according to whether the KN effect of \( \gamma_m \) electrons is important or not, i.e., the cases of \( v_{\text{KN}}(\gamma_m) < v_m \) and \( v_{\text{KN}}(\gamma_m) > v_m \).

3.2.1. Case I: \( v_{\text{KN}}(\gamma_m) < v_m \)

Unlike the slow-cooling case, the electron distribution at low energies is affected by the KN effect in this case and therefore the corresponding synchrotron spectrum may be changed. Following Nakar et al. (2009), we define \( v_0 \) as the synchrotron frequency of electrons of \( \gamma_0 \) (i.e., \( v_0 = v_{\text{syn}}(\gamma_0) \)), where \( Y(\gamma_0) = 1 \). According to whether \( \gamma_m \) is greater or smaller than \( \gamma_0 \), there are two subcases, i.e., (1) \( \gamma_0 < \gamma_m \) and (2) \( \gamma_0 > \gamma_m \).

1. Case Ia: \( \gamma_0 < \gamma_m \). This case applies when \( \epsilon_B \) is large. Define \( \gamma_0 = \Gamma m_e c^2 / h \nu_0 \) and \( \gamma_m = \Gamma m_e c^2 / h \nu_m \). In the energy range \( \gamma_0 < \gamma < \gamma_0 \), the electron distribution is \( N(\gamma_e) \propto \gamma_e^{-1} \) and the synchrotron spectrum is \( \nu F(\nu) \propto \nu \) (Wang et al. 2009; Nakar et al. 2009). Hence we have \( Y(\gamma_0) = Y(\gamma_0)(\gamma_0/\gamma_0)^{-1} \). Since \( \nu F(\nu) = \nu F(\nu_0)(\nu_0/\nu_m)^{1/2} \) and \( Y(\gamma_m) = \nu_c/\epsilon_B \) in this case, \( Y(\gamma_0) = \nu_c/\epsilon_B \). Then we obtain \( \gamma_0 = \gamma_0(\nu_0/\nu_m)^{1/2} \).

\[
\frac{v_0}{v_m} = \left( \frac{\gamma_0}{\gamma_m} \right)^{2} = 0.06 f_p^{-3} \epsilon_{e,-1} E_{54}^{-1/2} \epsilon_{B,-2}^{-3/2} y_0^{-3/2} (1 + z)^{-1/2},
\]  

(29)

one can further obtain the corresponding synchrotron frequency of \( \gamma_0 \) electrons,

\[
v_0 = 3.6 \times 10^{18} f_p^{-1} \epsilon_{B,-1} \text{Hz}.
\]  

(30)

Similarly, one can obtain the synchrotron frequency of \( \gamma_0 \) electrons,

\[
\hat{v}_0 = \left( \frac{\gamma_0}{\gamma_m} \right)^{2} v_m = \left[ \frac{v_{\text{KN}}(\gamma_m)}{v_0} \right]^{2} v_m
\times 5 \times 10^{19} f_p^{-3/2} \epsilon_{B,-1} E_{54}^{-1/2} y_0^{-3/2} (1 + z)^{1/2} \text{Hz}.
\]  

(31)

If \( \max(v_e, \hat{v}_0) \lesssim v_{\text{KN}}(\gamma_m) \lesssim v_0 \lesssim v_m \), we have

\[
Y(\gamma_0)[1 + Y(\gamma_m)] = \frac{\epsilon_e}{\epsilon_B} \left( \frac{v_{\text{KN}}(\gamma_m)}{v_0} \right) \left( \frac{v_0}{v_m} \right)^{1/2}.
\]  

(32)

As \( Y(\gamma_m) < 1 \) when \( \gamma_0 < \gamma_m \), we obtain

\[
Y(\gamma_0) \lesssim \frac{\epsilon_e}{\epsilon_B} \left( \frac{v_{\text{KN}}(\gamma_m)}{v_0} \right) \left( \frac{v_0}{v_m} \right)^{1/2}
= 0.2 f_p^{3/2} \epsilon_{e,-1} E_{54}^{-1/4} \epsilon_{B,-1}^{-3/4} y_0^{-3/4} (1 + z)^{-1/4}.
\]  

(33)

If \( \max(v_e, \hat{v}_0) \lesssim v_{\text{KN}}(\gamma_m) \lesssim v_0 \), it is easy to obtain

\[
Y(\gamma_m) = Y(\gamma_m) \frac{v_{\text{KN}}(\gamma_m)}{v_{\text{KN}}(\gamma_m)} = Y(\gamma_m) \left( \frac{v_0}{v_m} \right)^{-1/2} = 0.01 f_p^{-1/2} \epsilon_{B,-1} \text{Hz}.
\]  

(34)

In other cases, the derivation of \( Y(\gamma_m) \) is complicated. However, we note that as long as \( \nu_e > v_0 = 3.6 \times 10^{18} f_p^{-1} \epsilon_{B,-1} \text{Hz} \),

\[
Y(\gamma_m) < 1.
\]  

(35)

For \( h \nu_e = 100 \text{ MeV} \), \( \nu_e > v_0 \) is satisfied given that the fast-cooling condition is satisfied. Therefore we conclude that \( Y(\gamma_m) < 1 \) in case Ia.

2. Case Ib: \( \gamma_0 > \gamma_m \). This case applies when \( \epsilon_B \) is smaller. In this case, \( v_{\text{KN}}(\gamma_m) \lesssim v_m \lesssim v_0 \), so

\[
Y(\gamma_0)[1 + Y(\gamma_m)] = \frac{\epsilon_e}{\epsilon_B} \frac{v_{\text{KN}}(\gamma_m)}{v_m}.
\]  

(36)

From \( Y(\gamma_0) = Y(\gamma_m) \left( \frac{v_0}{v_m} \right)^{-1} = 1 \), we obtain \( \gamma_0 = \frac{\nu_0}{\epsilon_B \nu_m} \gamma_m \) and

\[
v_0 = 3.6 \times 10^{19} f_p^{-1} \epsilon_{B,-2} \text{Hz}.
\]  

(37)

As \( Y(\gamma_m) > 1 \) in this case, we get

\[
Y(\gamma_0) = \left[ \frac{\epsilon_e}{\epsilon_B} \frac{v_{\text{KN}}(\gamma_m)}{v_m} \right]^{1/2}
= 1.2 f_p^{3/2} \epsilon_{e,-1} E_{54}^{-1/4} \epsilon_{B,-1}^{-3/4} y_0^{-3/4} (1 + z)^{-1/4}.
\]  

(38)

Similarly, if \( \max(v_e, \hat{v}_0) \lesssim v_{\text{KN}}(\gamma_m) \lesssim v_0 \),

\[
Y(\gamma_m) = Y(\gamma_m) \frac{v_{\text{KN}}(\gamma_m)}{v_{\text{KN}}(\gamma_m)} = 0.03 f_p^{-1/2} \epsilon_{B,-2} \text{Hz}.
\]  

(39)

In the other case, we also have

\[
Y(\gamma_m) < 1
\]  

(40)

as long as \( \nu_e > v_0 \) is satisfied.
3.2.2. Case II: $v_{\text{KN}}(\gamma_m) > v_m$

At later times, when $t > 3 f_p^2 \epsilon_e^2 - 1 E_5^{1/3} \epsilon_B^{1/3} E_2^{-1/3} (1 + z)^{1/3}$ s, $v_{\text{KN}}(\gamma_m) > v_m$. In this case, the KN effect of $v_m$ electrons is unimportant and

$$Y(\gamma_m) = \left( \frac{e_e}{e_B} \right)^{1/2} = 3 \epsilon_{e,-1} \epsilon_{B,-2},$$

as the synchrotron emission typically peaks at $v_m$ in this case (Nakar et al. 2009). In order to calculate $Y(\gamma_m)$, let us first derive the ratios of $v_{\text{KN}}(\gamma_e)$ to two critical frequencies, which are respectively

$$\frac{v_{\text{KN}}(\gamma_e)}{v_m} = \left( \frac{v_m}{v_e} \right)^{1/2} \frac{v_{\text{KN}}(\gamma_m)}{v_m} = 0.04 f_p^{-2} \epsilon_{e,-1} \epsilon_{B,-2}^{1/4} E_5^{-1/4} t_1^{3/4} (1 + z)^{-1/4},$$

and

$$\frac{v_{\text{KN}}(\gamma_e)}{v_e} = \left( \frac{v_m}{v_e} \right)^{1/2} \frac{v_{\text{KN}}(\gamma_m)}{v_e} = 0.25 \epsilon_{e,-1} \epsilon_{B,-2}^{3/4} E_5^{1/4} t_1^{-1/4} (1 + z)^{3/4},$$

where $Y(\gamma_e) = Y(\gamma_m) = \sqrt{\epsilon_e/\epsilon_B}$ has been used. According to the relations among $v_{\text{KN}}(\gamma_e)$, $v_m$ and $v_e$, there are three subcases:

1. Case Ia: $v_c < v_{\text{KN}}(\gamma_e) < v_m$. In this case,

$$Y(\gamma_e) = Y(\gamma_m) \left( \frac{v_{\text{KN}}(\gamma_e)}{v_m} \right)^{1/2} = 0.7 f_p^{-1} \epsilon_{e,-1} \epsilon_{B,-2}^{1/3} E_5^{1/3} t_1^{-1/3} (1 + z)^{-1/3},$$

2. Case Ib: $v_{\text{KN}}(\gamma_e) < v_c$. In this case,

$$Y(\gamma_e) = Y(\gamma_m) \left( \frac{v_{\text{KN}}(\gamma_e)}{v_m} \right)^{1/2} \left( \frac{v_{\text{KN}}(\gamma_m)}{v_e} \right)^{4/3} = 2.2 \left[ 1 + Y(\gamma_e) \right]^{-1} f_p^{-1} \epsilon_{e,-1}^{-5/6} \epsilon_{B,-2}^{1/2} E_5^{1/3} n_{-1} t_1^{1/3} (1 + z)^{1/3}$$

3. Case Ic: $v_{\text{KN}}(\gamma_e) > v_m$. In this case,

$$Y(\gamma_e) = Y(\gamma_m) = 3 \epsilon_{e,-1} \epsilon_{B,-2}^{1/3}.$$
4. KN EFFECT ON THE HIGH-ENERGY SYNCHROTRON
AFTERGLOW LUMINOUSITY

The above analyses give the dependence of the Compton parameters \( Y_{\gamma_e} \), \( Y_{\gamma_c} \) (in the slow-cooling case) and \( Y_{\gamma_m} \) (in the fast-cooling case) on parameters such as \( \epsilon_c \), \( \epsilon_B \), \( E \), and \( n \). Since \( \epsilon_B \) and \( n \) are the least known among these parameters for GRB afterglows, we explore the value of \( Y_{\gamma_e} \) and \( Y_{\gamma_c} \) (or \( Y_{\gamma_m} \)) as a function of these two parameters. In Figures 1 and 2, we show the result for two different times, i.e., at \( t = 1 \) s and \( t = 10 \) s, respectively. We find \( Y_{\gamma_e} \) is smaller than a few at \( t = 1 \) s for the parameters \( \epsilon_B \) in the range from \( 10^{-6} \) to \( 10^{-4} \) and \( n \) in the range from \( 10^{-3} \) cm\(^{-3} \) to 10 cm\(^{-3} \). At \( t = 10 \) s, \( Y_{\gamma_c} \) is also smaller than a few in a wide range of parameter space (it is larger than a few only when \( n \) is as high as 10 cm\(^{-3} \) and \( \epsilon_B \) close to \( 10^{-6} \)). On the other hand, \( Y_{\gamma_c} \) or \( Y_{\gamma_m} \) can be more than 1 order of magnitude higher in the same parameter space. This implies that SSC loss of high-energy electrons that produce high-energy (\( \gtrsim 100 \) MeV) afterglow photons is typically small. As a result, the synchrotron luminosity at high energies is correspondingly high, which enables the detection of early high-energy afterglow emission by Fermi–LAT.

In order to see whether the SSC emission contributes to the high-energy afterglow emission at Fermi–LAT energy band, we calculate the spectral energy distribution of the afterglow emission numerically at early times. Assuming adiabatic evolution of the blast wave and using the electron distribution given in

Section 2, we calculate the synchrotron radiation spectrum as well as the SSC spectrum with a full KN cross section taken into account (see Equations (2) and (11) of He et al. (2009) for the description of the dynamic and the full KN cross section). Figure 3 shows the \( \nu F_\nu \) spectra of the afterglow synchrotron emission and the SSC emission for the slow-cooling case at times \( t = 1 \) s and \( t = 10 \) s when the KN effect is taken into account. In the Fermi–LAT energy band, the synchrotron component is dominated at both times. The SSC component becomes dominated only at energies above the maximum synchrotron photon energy of shock-accelerated electrons, at which the flux usually becomes, however, too low to be detectable by Fermi–LAT. The spectral energy distribution of the afterglow emission for the fast-cooling case at times \( t = 10 \) s and \( t = 100 \) s is shown in Figure 4. Similarly, SSC contribution to the high-energy emission at energies below 100 GeV is negligible at these times. Figure 4 (see the bottom panel) also shows that the spectrum becomes harder at energies above \( 10^7 \) eV. This is caused by the decreased IC loss suppression on the synchrotron flux at high energies due to the KN effect.

The Compton parameters also vary with time. In the above analytic calculation in Section 3, we have shown that \( Y_{\gamma_e} \) increases with time as \( t^{1/2} \) in the slow-cooling case as long as \( \nu_{\text{KN}}(\gamma_e) < \nu_m \). When \( \nu_{\text{KN}}(\gamma_e) > \nu_m \), \( Y_{\gamma_e} \) starts to decrease with time. We calculate \( Y_{\gamma_c} \) numerically using Equation (11) and show the evolution of \( Y_{\gamma_c} \) and \( Y_{\gamma_m} \) with time in the top panel of Figure 5. If \( Y_{\gamma_e} > 1 \) as well, as in the case of some parameter spaces shown in Figures 1 and 2, the
decay of the synchrotron afterglow emission will be faster than what is predicted by the standard synchrotron afterglow theory (i.e., steeper than $t^{-3p/4}$ for $v_\gamma > v_\ell$), since the synchrotron luminosity at frequency $v_\gamma$ scales as $1/(1 + Y(\gamma_\ell))$. The decay could be steeper by a factor of $\Delta \alpha = 1/2$ at most. We numerically calculate the light curves of the afterglow emission assuming adiabatic evolution of the blast wave and using the electron distribution given in Section 2. The light curves of the synchrotron emission, the SSC emission, and the sum of them at $h\nu_\gamma = 100$ MeV are shown in the bottom panel of Figure 5 (the black lines). As a comparison, we also show the light curves for another set of parameters in the fast-cooling case when the KN effect is taken into account. In Figure 6, we also show the light curves of the synchrotron emission, the SSC emission, and the sum of them at $h\nu_\gamma = 100$ MeV for two sets of parameters. They clearly indicate that the temporal decay of the high-energy afterglow emission becomes steeper than $t^{-3p/4}$ when $Y(\gamma_\ell) \gtrsim 1$. So we conclude that the light curve of the high-energy afterglow emission could be also steeper in the fast-cooling case when the KN effect is taken into account.

5. IMPLICATIONS FOR FERMI–LAT OBSERVATION OF AFTERGLOW EMISSION

As has been shown in Figures 1 and 2, the Compton parameters for the electrons that produce high-energy gamma-ray afterglow emission are typically small at early time $t \lesssim 10^3$ s, i.e., $Y(\gamma_\ell) \lesssim 1$ for a wide range of parameter space. This has an important implication for the detectability of high-energy afterglow emission by Fermi–LAT, since a low $Y(\gamma_\ell)$ leads to a high synchrotron luminosity at high-energies. The low-energy electrons that produce early X-ray and optical afterglow emission, however, still suffer from the strong IC loss (i.e., $Y(\gamma_e)$ or $Y(\gamma_m)$ are typically high) and therefore the observed
The X-ray emission before respectively, the light curves with and without KN effect taken into account. and the sum of them at panel shows the light curves of the synchrotron emission, the SSC emission, gamma-ray emission seen in GRB090510 and GRB090902B. faster decay of the high-energy synchrotron afterglow emission, late times for some range of parameter space. This will lead to observations of the short burst GRB090510 are reported in Abdo (Ghirlanda et al. 2010). The high-energy emission above 100 MeV shows a simple power-law decay after the peak, with a decay slope of $\alpha = -0.99$ at high-energy frequency with $v_\gamma > v_c$ in the standard synchrotron scenario, since the different cooling behavior of electrons cause a difference of $\Delta \alpha = 0.25$ in the decay slope. This slope is much shallower than the observed slope. The high-energy gamma-ray emission observed from the long burst GRB090902B by Fermi/LAT is reported in Abdo et al. (2009b). LAT detected high-energy gamma-ray emission above 100 MeV on timescales much longer than the prompt phase. The time-integrated spectrum of the LAT detected emission after the prompt phase is consistent with $\beta_{LAT} = -1.1 \pm 0.1$. Its flux declines as a simple power law with a decay slope of $\alpha_{LAT} = -1.5 \pm 0.1$ from $t = 25$ s to 1 ks. Taking $p = -2\beta_{LAT} = 2.2 \pm 0.2$, the standard synchrotron emission predicts a decay slope of $\alpha = -(3p - 2)/4 = -1.15 \pm 0.1$, which is shallower than the observed decay slope, similar to the case in GRB090510. We suggest that one possible origin for the discrepancy in the theoretical and observed decay slopes seen in GRB090510 and GRB090902B is due to the KN effect on high-energy electrons, as discussed in Section 4.

The spectrum of the afterglow emission at high-energies can be changed as well when $Y(y_e)$ increases to be larger than 1 at later times, since the electron distribution is changed in this case. Therefore, the time-resolved spectrum at high energies at

\begin{equation}
\alpha_{x,2} = 2.18 \pm 0.10 \text{ after the break is consistent with a jet break model (Kumar & Barniol Duran 2009). Such an interpretation predicts a decay slope of } \alpha = \alpha_{x,1} - \Delta \alpha = -0.99 \text{ at high-energy frequency with } v_\gamma > v_c \text{ in the standard synchrotron scenario, since the different cooling behavior of electrons cause a difference of } \Delta \alpha = 0.25 \text{ in the decay slope. This slope is much shallower than the observed slope. The high-energy gamma-ray emission observed from the long burst GRB090902B by Fermi/LAT is reported in Abdo et al. (2009b). LAT detected high-energy gamma-ray emission above 100 MeV on timescales much longer than the prompt phase. The time-integrated spectrum of the LAT detected emission after the prompt phase is consistent with } \beta_{LAT} = -1.1 \pm 0.1. \text{ Its flux declines as a simple power law with a decay slope of } \alpha_{LAT} = -1.5 \pm 0.1 \text{ from } t = 25 \text{ s to } 1 \text{ ks. Taking } p = -2\beta_{LAT} = 2.2 \pm 0.2, \text{ the standard synchrotron emission predicts a decay slope of } \alpha = -(3p - 2)/4 = -1.15 \pm 0.1, \text{ which is shallower than the observed decay slope, similar to the case in GRB090510. We suggest that one possible origin for the discrepancy in the theoretical and observed decay slopes seen in GRB090510 and GRB090902B is due to the KN effect on high-energy electrons, as discussed in Section 4.}

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure7.png}
\caption{Figure 7. Top panel shows the evolution of Compton parameters $Y(y_e)$ and $Y(y_{\gamma})$ with time in the fast-cooling case for the parameters $\epsilon_\gamma = 0.1$, $\epsilon_\gamma = 0.01$, $n = 0.1 \, \text{cm}^{-3}$, $E = 10^{54} \, \text{erg}$, $p = 2.2$, $\Gamma_0 = 3000$, and $z = 1$. The bottom panel shows the light curves of the synchrotron emission, the SSC emission, and the sum of them at $\nu_\gamma > \nu_c$ and $Y > 1$. Black lines and red lines denote, respectively, the light curves with and without KN effect taken into account. (A color version of this figure is available in the online journal.)}
\end{figure}

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure8.png}
\caption{Figure 8. Same as Figure 7, but for $\epsilon_\gamma = 0.005$ and $n = 1 \, \text{cm}^{-3}$. (A color version of this figure is available in the online journal.)}
\end{figure}
later times can be harder than $F_\nu \sim \nu^{-p/2}$. However, the time-integrated spectrum of the high-energy emission will be still $F_\nu \sim \nu^{-p/2}$ since the time-integrated fluence is dominated by the contribution at early times (note that the flux decays usually faster than $t^{-1}$), at which times $Y(\gamma_*) < 1$ usually. The low-significance data of the time-resolved spectra of GRB090510 and GRB090902B prevent us from identifying the KN effect in the spectrum (De Pasquale et al. 2010; Abdo et al. 2009b). Obtaining high-significance time-resolved spectra at high-energies from bright GRBs in the future would be very useful to do this.

6. SUMMARY

We have studied the KN effect on early high-energy afterglow emission of GRBs. Our findings are summarized as follows.

1. The IC scatterings between high-energy electrons that produce early-time high-energy ($\gtrsim 100$ MeV) afterglow emission and synchrotron peak photons of the afterglow are generally in the deep KN scattering regime. As a result, the IC loss of these electrons is small and the synchrotron luminosity at $h\nu \gtrsim 100$ MeV is high, which is favorable for the detection of high-energy gamma-ray emission from the early afterglow by Fermi–LAT.

2. The high-energy gamma-ray emission at the early afterglow phase is dominated by the synchrotron emission. The SSC emission becomes dominated only at energies above the maximum synchrotron photon energy of shock-accelerated electrons, but at such energies the SSC flux is usually too weak to be detectable by Fermi–LAT.

3. The KN suppression effect of high-energy electrons weakens with time, so that the IC loss increases with time. In the parameter space where the Compton parameter is $Y(\gamma_*) > 1$, the increasing IC loss leads to a faster temporal decay of the synchrotron afterglow emission at high frequency. The decay slope could be steeper by a factor of $\Delta \alpha = 0.5$ at most under favorable conditions. This may explain the somewhat faster than expected decay of the early-time high-energy emission observed in GRB090510 and GRB090902B.

This work is supported by the NSFC under grants 10973008, 10873009, 10843007, 10503012, 10621303, and 10633040, the 973 program under grants 2009CB824800 and 2007CB815404, the Foundation for the Authors of National Excellent Doctoral Dissertations of China, the Program for New Century Excellent Talents in University, the Qing Lan Project, and the NASA grants (NNX09AT72G and NNX08AL40G).

REFERENCES

Abdo, A., et al. 2009a, arXiv:0908.1832
Abdo, A., et al. 2009b, ApJ, 706, L138
Abdo, A., et al. 2009c, ApJ, 707, 380
De Pasquale, M., Schady, P., & Kuin, N. P. M. 2010, ApJ, 709, L146
Dermer, C. D., & Atoyan, A. 2004, A&A, 418, L5
Fan, Y. Z., et al. 2008, MNRAS, 384, 1483
Gao, W. H., et al. 2009, ApJ, 706, L33
Ghirlanda, G., Ghisellini, G., & Nava, L. 2010, A&A, 510, L7
Ghisellini, G., Ghirlanda, G., & Nava, L. 2009, arXiv:0910.2459
Gou, L. J., & Mészáros, P. 2007, ApJ, 668, 392
Granot, J., & Guetta, D. 2003, ApJ, 598, L1
He, H. N., Wang, X. Y., Yu, Y. W., & Mészáros, P. 2009, ApJ, 706, 1152
Kumar, P., & Barniol Duran, R. 2009a, MNRAS, 400, L75
Kumar, P., & Barniol Duran, R. 2009b, arXiv:0910.5726
Kumar, P., & Panaitescu, A. 2001, ApJ, 554, 667
Li, Z., & Waxman, E. 2006, ApJ, 651, 328
Nakar, E., Ando, S., & Sari, R. 2009, ApJ, 703, 675
Pe'er, A., & Waxman, E. 2004, ApJ, 603, L1
Sari, R., & Esin, A. A. 2001, ApJ, 548, 787
Sari, R., Piran, T., & Narayan, R. 1998, ApJ, 497, L17
Wang, X. Y., Cheng, K. S., Dai, Z. G., & Lu, T. 2005, A&A, 439, 957
Wang, X. Y., Dai, Z. G., & Lu, T. 2001a, ApJ, 546, L33
Wang, X. Y., Dai, Z. G., & Lu, T. 2001b, ApJ, 556, 1010
Wang, X. Y., Li, Z., Dai, Z. G., & Mészáros, P. 2009, ApJ, 698, L98
Zhang, B., & Mészáros, P. 2001, ApJ, 559, 110