GRB 080916C: ON THE RADIATION ORIGIN OF THE PROMPT EMISSION FROM keV/MeV TO GeV

XIANG-YU WANG1, ZHUO LI2,3, ZI-GAO DAI1, AND PETER MÉSZÁROS4,5
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, University Park, PA 16802, USA
5 Department of Physics, Pennsylvania State University, University Park, PA 16802, USA

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ABSTRACT

Fermi observations of high-energy gamma-ray emission from GRB 080916C shows that its spectrum is consistent with the band function from MeV to tens of GeV. Assuming one single emission mechanism dominates in the whole energy range, we show that this spectrum is consistent with synchrotron origin by shock-accelerated electrons. The simple electron inverse-Compton model and the hadronic model are found to be less viable. In the synchrotron scenario, the synchrotron self-Compton scattering is likely to be in the Klein–Nishina (KN) regime and therefore the resulting high-energy emission is subdominant, even though the magnetic field energy density is lower than that in relativistic electrons. The KN inverse-Compton cooling may also affect the low-energy electron number distribution and hence results in a low-energy synchrotron photon spectrum below the peak energy. Under the framework of the electron synchrotron interpretation, we constrain the shock microphysical parameters and derive a lower limit of the upstream magnetic fields. The detection of synchrotron emission extending to about 70 GeV in the source frame in GRB 080916C favors the Bohm diffusive shock acceleration if the bulk Lorentz factor of the relativistic outflow is not significantly greater than thousands.

Key words: gamma rays: bursts – shock waves

1. INTRODUCTION

It was recently reported that the Fermi satellite has detected strong >100 MeV emission from a very energetic long-duration burst GRB 080916C (Abdo et al. 2009). At a redshift of \(z = 4.35 \pm 0.15\) (Greiner et al. 2009), the burst is the most energetic one ever, with an isotropic gamma-ray energy \(E_{\gamma} \approx 8 \times 10^{54}\) erg, which is released over a duration of \(T_{90} \approx 60\) s. Equally remarkably, more than 10 photons with energy above GeV are detected, with the highest energy one at 13 GeV (in the observer frame). The spectra of all five time intervals, designated as times \(a–e\) in the light curves of GRB 080916C (Abdo et al. 2009), are well fit by the empirical Band function (Band et al. 1993), which smoothly joins low- and high-energy power laws. The high-energy power law extends to GeV (Band et al. 1993), which smoothly joins low- and high-energy power laws. The high-energy power law extends to GeV energies, without any additional spectral component visible. The peak energy of the spectra during these intervals is around \(E_p \approx 400\) keV–1 MeV. Except during the first time interval, the low-energy and high-energy photon spectral indices of the prompt emission are constant and consistent with \(\alpha \approx -1.0\) and \(\beta \approx -2.2\), respectively. With its high temporal resolution, INTEGRAL detected the temporal variability of the keV/MeV emission on timescales as short as 100 ms with high statistical significance (Greiner et al. 2009). So the variability timescale in the local source frame is \(t_v \approx 100\) ms/(1 + \(z\)) = 20 ms.

The nonthermal synchrotron radiation by electrons has been suggested to be a possible mechanism for the 10 keV–MeV emission (see Mészáros 2006 and Zhang 2007 for recent reviews), but one famous problem remains so far, i.e., the low-energy photon spectral index \(\alpha\) is incompatible with the index \(-3/2\) that is expected from fast-cooling electrons (e.g., Preece et al. 1996; Ghisellini et al. 2000). Electron inverse Compton emission has been a competitive mechanism (e.g., Panaitescu & Mészáros 2000). The fact that one single spectral component fits the spectrum of the prompt emission from

10 keV to GeV in GRB 080916C suggests that one emission mechanism dominates in this whole energy range.6 In this Letter, we study the constraint that this puts on the emission mechanism. Abdo et al. (2009) mentioned as one of the possibilities that the delay of high-energy gamma-ray emission relative to low-energy emission in GRB 080916C could be a result of longer acceleration time needed for higher energy protons or nuclei in hadronic emission models. In accordance with this, we also study whether such hadronic models could be a possible mechanism that produces the keV/MeV to GeV emission in GRB 080916C.

2. THE SYNCHROTRON MODEL AND PARAMETER CONSTRAINTS

Assuming that in GRB shocks fractions of \(\epsilon_B\) and \(\epsilon_e\) of the shock internal energy are converted into the energy in the magnetic field and electrons, respectively. To ensure a high radiation efficiency for the prompt emission, it is usually assumed that the electrons are rapidly cooling, so the energy density in gamma-ray emission \(U_{\gamma}\) is equal to the electron energy density \(U_e\). The magnetic field is given by

\[
\frac{B^2}{8\pi} = \left(\frac{\epsilon_B}{\epsilon_e}\right) U_{\gamma} = \left(\frac{\epsilon_B}{\epsilon_e}\right) \frac{L_{\gamma}}{4\pi R^2 c \Gamma^2}
\]

where \(R\) is the radius of the shock, \(L_{\gamma}\) is the luminosity in gamma-ray emission, and \(\Gamma\) is the bulk Lorentz factor. The detection of GeV photons suggests that the emission region has a bulk Lorentz factor \(\Gamma \gtrsim 10^3\) (Greiner et al. 2009; Abdo

6 Dropping the assumption of the same origin for both MeV and high-energy emission, Li (2008) explained the high-energy emission in GRB 080916C as delayed, “residual” emission from subsequent collisions at larger and larger radii in the baryonic outflow.
et al. 2009. From the casualty constraint, the emission radius is \( R = 15^2 \epsilon c t_v \). In the synchrotron model for the 10 keV–GeV emission, by the use of \( \epsilon_p = h \nu_p = \phi_v \frac{3 \pi B}{3 \pi m_e} \epsilon^2 \Gamma \), one can derive the Lorentz factor of electrons that radiate at the GRB peak energy \( \epsilon_p \),

\[
y_m = \left( \frac{4 \pi m_e c^2 \epsilon_p}{3 \phi_v h q} \right)^{1/2} \left( \frac{\epsilon}{\epsilon_B} \right)^{1/4} \left( \frac{2 L_\Gamma}{R^2 c} \right)^{-1/4} \epsilon^{1/2} = 2.5 \times 10^7 \left( \frac{\epsilon}{\epsilon_B} \right)^{1/4} L_{\gamma,54}^{-1/4} \Gamma_{\gamma,1/2}^{1/2} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right)^{1/2}. \tag{2}
\]

where \( \phi_v \simeq 0.5 \) is the coefficient defined in Wijers & Galama (1999) and \( q \) is the electron charge. Define \( \gamma_T \) as the Lorentz factor of electrons below which the scattering with peak energy photons is in the Thomson scattering regime, i.e.,

\[
\gamma_T = \Gamma m_e c^2 / \epsilon_p = 250 \Gamma_5 (\epsilon_p / 2 \text{ MeV})^{-1}. \tag{3}
\]

Unless \( \epsilon_c \lesssim 10^{-4} \epsilon_B \), which is unreasonable in terms of the burst energetics, the IC scattering between \( \gamma_m \) electrons and the bulk of the gamma-ray emission should be in the Klein–Nishina (KN) regime.

The KN Compton cooling of electrons may affect the low-energy electron distribution at \( \gamma < \gamma_m \) and hence affect the low-energy spectral slope below \( \epsilon_p \) (e.g., Rees 1967; Derishev et al. 2003), as we show below. Consider a population of electrons injected into a box with magnetic field \( B \) in a power-law form \( N(\gamma) \propto \gamma^{-p} \) for \( \gamma \geq \gamma_m \). These electrons will cool down rapidly through synchrotron and/or IC radiation. As the electron cools its energy changes as \( \gamma \), so we have

\[
v F_v [1 + k(\gamma) U_{\gamma} / U_B] \propto \gamma \tag{4}
\]

for \( \gamma < \gamma_m \), where \( v \) is the synchrotron frequency of \( \gamma \)-electrons and \( k(\gamma) \) accounts for the reduction of the effective photon energy density for IC scattering of the \( \gamma \)-electrons due to the KN effect. Define \( h_{\text{KN}} = \Gamma m_e c^2 / \gamma \) as the critical energy of the photons with which the IC scattering of \( \gamma \)-electrons is just in the KN regime. For a low-energy photon spectrum of the form \( \nu F_\nu \propto \nu^\delta \nu (\nu < \nu_p) \), we have

\[
k(\gamma) \simeq \frac{U_{\nu}(\nu < v_{\text{KN}})}{U_{\gamma}} = \left( \frac{v_{\text{KN}}}{\nu_p} \right)^\delta = \left( \frac{\nu}{\nu_T} \right)^\delta \tag{5}
\]

for \( \gamma_T < \gamma \leq \gamma_m \) and \( k \simeq 1 \) for \( \gamma < \gamma_T \). In the case of \( \gamma>T < \gamma \leq \gamma_m \), if \( kU_{\gamma} / U_B \gg 1 \), i.e., the IC cooling is still dominant even though the scatterings are in the KN regime, one can obtain \( v F_v \propto \nu^{3+1/2} \propto \nu^{(\delta+1)/2} \), where we have used \( \nu \propto \gamma^2 \) in the last step. Equating this synchrotron spectral index with the initially assumed spectral index \( F_v \propto \nu^{\delta-1} \), one can derive

\[
\delta = 1, \ F_v \propto \nu^0. \tag{6}
\]

This can explain the observed low-energy photon index of \( \alpha = -1.0 \) in GRB 080916C if the condition, \( U_B \lesssim k(\gamma) U_{\gamma} \) for \( \gamma_T < \gamma < \gamma_m \), is satisfied.

On the other hand, the high-energy spectrum of GRB 080916C above \( \epsilon_p \) is consistent with the synchrotron spectrum produced by fast-cooling electrons above \( \gamma_m \), i.e., \( F_v \propto \nu^{-p/2} \) with \( p = -2(1 + \beta_2) = 2.4 \). The dominance of synchrotron cooling above \( \epsilon_p \) implies that \( U_B \gtrsim k(\gamma) U_{\gamma} \) for electrons with Lorentz factor \( \gamma \gtrsim \gamma_m \). Therefore, we find that, at \( \gamma = \gamma_m \), \( U_B \simeq kU_{\gamma} \). Since \( U_{\gamma} \simeq U_e \) for fast-cooling electrons, the requirement \( U_B \simeq k(\gamma_m) U_{\gamma} \) translates into \( \epsilon_c / \epsilon_B \approx \gamma_m \epsilon_p / \Gamma m_e c^2 \), which gives

\[
\frac{\epsilon_c}{\epsilon_B} \approx 20 L_{\gamma,54}^{-1/3} \Gamma_{\gamma,1/2}^{-2/3} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right)^2. \tag{7}
\]

Note that the transition region between the two asymptotic power laws at low- and high-energy ends in the Band function is rather wide, so the above requirement, \( U_B \simeq k(\gamma_m) U_{\gamma} \), should be regarded as an order of magnitude of estimate. In addition, this requirement applies only to large \( \epsilon_p \) bursts, because for low \( \epsilon_p \) bursts, the IC scattering may be no longer in the KN regime. A signature that high \( \epsilon_p \) bursts have \( \alpha \) preferentially close to \(-1\) can be seen in the analysis of Preece et al. (1996). For some low \( \epsilon_p \) bursts that have \( \alpha \lesssim -1 \), some other mechanisms may be at work.

For electrons with \( \gamma \lesssim \gamma_T \), the IC scatterings with peak energy photons are in the Thomson scattering regime, so \( k(\gamma) = 1 \) and \( N(\gamma) \propto \gamma^{-p} \), leading to a conventional fast-cooling photon spectrum of \( F_\nu \propto \nu^{-p/2} \). Observations show a single power-law spectrum \( F_\nu \propto \nu^\beta \) from 10 keV to \( \sim \text{MeV} \) in GRB 080916C, implying that \( \gamma_m / \gamma_T \gtrsim 10 \), and one can therefore obtain a constraint

\[
\frac{\epsilon_c}{\epsilon_B} \gtrsim 1 L_{\nu,54}^{-1/3} \Gamma_{\nu,1/2}^{-2} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right)^{-6}. \tag{8}
\]

Due to the fact that the IC scatterings between \( \gamma_m \) electrons and the peak energy photons with energy \( \epsilon_p \) are in the KN regime, the IC emission peaks at

\[
h_{\nu,\text{IC}} = \Gamma m_e c^2
\]

\[
\simeq 1 \left( \frac{\epsilon}{\epsilon_B} \right)^{1/4} L_{\gamma,54}^{-1/4} \Gamma_{\gamma,1/2}^{1/2} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right)^{1/2} \text{ TeV}, \tag{9}
\]

with a flux \( v F_v \nu (\epsilon_p = h_{\nu,\text{IC}}) = Y (\gamma_m) \nu_p F_{v_p} \simeq v_p F_{v_p} \), where \( Y \) is the Compton parameter. For a flat synchrotron spectrum with \( \beta \simeq -2 \) above \( \nu_p \), it is natural that the IC component is not seen at high energies since \( v F_v \nu (\epsilon_p = 70 \text{ GeV}) \gtrsim v F_{v,\text{IC}} (\epsilon_p = 70 \text{ GeV}) \) for GRB 080916C.

In the above, we have not assumed any model for the dissipation mechanism of the shocks. In the popular internal shock scenario, the typical Lorentz factors of the shocked electrons is \( \gamma_m = \gamma_m (m_p/m_e) \Gamma_{sh} \), where \( \Gamma_{sh} \) is the shock Lorentz factor, which is equal to the relative Lorentz factor of the two colliding shells. For GRB 080916C, we have obtained a constraint \( \gamma_m \simeq 5 \times 10^4 L_{\gamma,54}^{-1/3} \Gamma_{\gamma,1/2}^{-2/3} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right)^{-1} \). So if internal shock applies to GRB 080916C, we would need a large relative Lorentz factor, \( \Gamma_{sh} \simeq 8 (\epsilon_p)^{-1} L_{\gamma,54}^{-1/3} \Gamma_{\gamma,1/2}^{-2/3} \left( \frac{\epsilon_p}{2 \text{ MeV}} \right) \). This could be caused by the interaction among the shells whose Lorentz factors have a large contrast (Yu et al. 2009). Of course, the shock could also arise from the magnetic reconnection or turbulence (e.g., Thompson 1994; Mészáros & Rees 1997; Lyutikov & Blandford 2003; Narayan & Kumar 2008; Zhang & Pe’er 2009) and we do not have the estimate of the shock Lorentz factor from the first principle.
3. ALTERNATIVE MODELS FOR MEV–10 GEV EMISSION?

3.1. The One-Zone SSC Scenario

Let us explore whether the simple IC scenario (i.e., one-zone synchrotron self-Compton (SSC) scenario) can explain the single power-law spectrum from MeV to 10 GeV in GRB 080916C. Suppose that the first-order SSC of electrons with energy \( \gamma_m \) produce the peak emission \( \epsilon_p = 2 \text{ MeV} \). Since this IC emission is not hidden by the synchrotron emission, one will expect that the second-order IC emission appears at high energy if the second-order IC peak is located within the observation energy window and that the second-order IC scattering is still in the Thomson scattering regime.

The fact that we did not see the second-order IC component implies that \( \gamma_m^2 \epsilon_p > 70 \text{ GeV} \) or \( \gamma_m \epsilon_p \gg m_e c^2 \), so we have \( \gamma_m > 190 (\epsilon_p/2 \text{ MeV})^{-1/2} \) or \( \gamma_m > 250 \Gamma_s (\epsilon_p/2 \text{ MeV})^{-1} \). Since \( \epsilon_p = h \nu_{\text{syn},p} \gamma_m^2 \), one obtains the synchrotron peak frequency at \( h \nu_{\text{syn},p} = 55(\gamma_m/190)^{-2}(\epsilon_p/2 \text{ MeV}) \text{ eV} \). Then one can obtain an upper limit of the magnetic field \( B = h \nu_{\text{syn},p} / (\phi_{\text{4MeV}} \gamma_m \Gamma_s) = 140(\gamma_m/190)^{-4} \Gamma_s^{-1}(\epsilon_p/2 \text{ MeV}) \text{ G} \).

With this magnetic field, we can derive an upper limit for \( \epsilon_B/\epsilon_e \).

In the case that second-order IC field is in the Thomson regime, \( U_e = (Y + 1) U_y \), so from \( \epsilon_B/\epsilon_e = U_B/U_e = U_B/(U_y + 1) U_y, \) we obtain \( \gamma_m \epsilon_e/\epsilon_p = 3 \times 10^{-3} \gamma_m/190 \) for a spectrum \( \nu F_\nu \propto \nu^{\delta_B} \) below \( \epsilon_p = (\epsilon_e/\epsilon_p)^{1/3} \) (Kobayashi et al. 2007), one gets \( \gamma_m = 170(\gamma_m/190)^{4}(\epsilon_p/2 \text{ MeV})^{-1} \nu_\gamma^{-1} \). As \( \gamma_m > 190 \), so the radio emission in the second-order IC will be \( E_{\text{2nd,IC}} = Y E_y \gamma \geq 1.3 \times 10^{57} \text{ erg} \), which is too large to be realistic. Such an energy crisis problem has also been found in the case of GRB 080319B for the IC scenario of the prompt MeV emission (Piran et al. 2009).

On the other hand, if the second-order IC is already in the deep KN regime (for \( \gamma_m \gg \gamma_T = 250 \Gamma_s \)), \( \epsilon_B/\epsilon_e = U_B/U_\gamma = U_B/(\Gamma_s + 1) U_y \), the second-order Compton \( Y \) parameter is \( Y_{\text{2nd}} = (\epsilon_e/\epsilon_B)^{1/3} (\gamma_e/\gamma_T)^{-1} \) for a spectrum \( \nu F_\nu \propto \nu^{\delta_B} \) below \( \epsilon_p \). (see Equation (5)). To get \( Y_{\text{2nd}} \leq 1 \), one needs \( \gamma_m \geq (\epsilon_e/\epsilon_B)^{1/2} (\gamma_e/\gamma_T)^{-1} = 250(\epsilon_p/2 \text{ MeV})^{-1} \Gamma_s^{-1} \). For a spectrum \( \nu F_\nu \propto \nu^{\delta_B} \) below \( \epsilon_p \), which is in conflict with the precondition \( \epsilon_e/\epsilon_B = Y_2 \leq 1 \). This means that significant suppression of the second-order IC emission by KN scattering cannot be fulfilled. So we conclude that the simple one-zone SSC model does not work for the MeV to 10 GeV emission in GRB 080916C.

3.2. The Hadronic Scenario

We first study whether the proton synchrotron emission can produce the MeV–10 GeV emission of GRB 080916C. The photon spectrum index above \( \epsilon_p, \beta = -2.2 \), we imply that the proton distribution index is \( p \simeq 2.4 \) for fast-cooling protons or a very steep index \( p \simeq 3.4 \) for slow-cooling ones. In the proton synchrotron scenario, the observed peak emission at \( \epsilon_p \) is produced by protons with a Lorentz factor of \( \gamma_p = \left( \frac{4 \pi e^2 m_p c^3}{3 m_p q^2 B^2 H} \right)^{1/2} \). The synchrotron cooling time of these protons in the comoving frame is \( t_{\text{syn}} = 6 \pi m_p c/\varsigma(\sigma m_T^2 \gamma_p B^2) \).

Define that the magnetic field energy density is a fraction of \( \xi_B \) of the comoving frame photon energy density, i.e., \( U_B = \xi_B U_y \). So the synchrotron radiation efficiency of the proton protons is \( \eta = \min[1, t_\gamma/t_{\text{syn}}] \), where \( t_\gamma = R/\Gamma c \) is the dynamic time in the comoving frame, which is equal to the comoving frame variability time, \( t_d = t_\gamma = \Gamma \tau_\gamma \).

As long as \( t_\gamma/t_{\text{syn}} \leq 1 \), we have a radiation efficiency for \( \gamma_p \) protons

\[
\eta(\gamma_p) = \frac{t_\gamma}{t_{\text{syn}}} = 3 \times 10^{-4} \frac{\gamma_p^{3/4} \Gamma_m^{3/4} \epsilon_p^{1/2}}{2 \text{ MeV}} \frac{1}{\Gamma_3}.
\]

Such a low radiation efficiency implies an unrealistically large energy in protons, a factor of \( 1/\eta \simeq 3 \times 10^5 \) higher than the energy in gamma-rays. The radiation efficiency is quite low \( (\eta \sim 2 \times 10^{-5}) \) even for protons that produce the high-energy gamma-rays of energy \( \gtrsim 100 \text{ MeV} \). So there is no room for the proton synchrotron model even in the assumption that the high-energy gamma-ray emission belongs to a different component than the MeV component. If the spectrum above \( \epsilon_p \) is interpreted as arising from fast-cooling protons, as required in the case of \( 2 \lesssim p \lesssim 3 \), one would need \( \xi_B = U_B/U_y \gtrsim 5 \times 10^4 \left( \gamma_p/10^2 \right)^{-3/2} \), which is also unreasonable, as the total energy in the magnetic field is too large to be realistic for a GRB.

Let us also explore the scenario of the secondary emission from the hadronic photoprocess. Detection of high-energy gamma-rays of energy greater than 10 GeV puts a constraint on the opacity of \( \gamma\gamma \) absorption. As the hadronic \( p\gamma \) opacity is related with the \( \gamma\gamma \) opacity, a higher maximum photon energy, hence a lower \( \gamma\gamma \) absorption opacity, would imply a lower hadronic radiation efficiency (e.g., Dermer et al. 2008). It is useful to express the hadronic \( p\gamma \) efficiency as a function of the pair production optical depth \( \tau_{\gamma\gamma} \). Following Waxman & Bahcall (1997), the optical depth for pair production of a photon of energy \( \epsilon_h \) is \( \tau_{\gamma\gamma}(\epsilon_h) = \frac{R}{\Gamma_{\gamma\gamma}} = \frac{R \sigma_T U_e \epsilon_h}{16 \Gamma_s m_e c^2} \), where \( \Gamma_{\gamma\gamma} \) is the mean free path. For the simplicity of calculation, we have assumed a photon spectrum \( \beta_\gamma = -2 \) above \( \epsilon_p \), which is a good approximation for GRB 080916C. The fraction of energy lost by protons to pions is \( f_\pi = \frac{R \sigma_T U_e \epsilon_h}{16 \Gamma_s m_e c^2} \), where \( \sigma_T \) is the cross section of the \( \gamma\gamma \) reaction at the resonance and \( \epsilon_h \) is the mean energy loss in one interaction. So the maximum photon efficiency is

\[
f_\pi = 2 \times 10^{-3} \Gamma_s^{-1} \left( \frac{\epsilon_h}{70 \text{ GeV}} \right)^{-1} \Gamma_s^{-1} \frac{1}{2 \text{ MeV}} \frac{1}{\Gamma_{\gamma\gamma}(\epsilon_h)}
\]

\[
= 2 \times 10^{-3} \Gamma_s^{-1} \Gamma_{\gamma\gamma}(\epsilon_h)^{-6} \frac{1}{\Gamma_{\gamma\gamma}(\epsilon_h)} \left( \frac{\epsilon_h}{70 \text{ GeV}} \right)^{-1} \frac{1}{2 \text{ MeV}} \frac{1}{\Gamma_{\gamma\gamma}(\epsilon_h)}
\]
motion between the upstream and downstream plasma is only mildly relativistic (such as in internal shocks). For an electron being shock-accelerated, the residence time in downstream and upstream regions are, respectively, \( t'_f = \kappa_d e \varepsilon_f / q B_d c \) and \( t'_u = \kappa_u e \varepsilon_u / q B_u c \), where \( \varepsilon_f \) is the energy of accelerated electrons, \( B_d \) and \( B_u \) are respectively the magnetic fields in the shock downstream and upstream, and \( \kappa_d, \kappa_u \gtrsim 1 \) parameterizes the efficiency of shock acceleration, with \( \kappa_d, \kappa_u \approx 1 \) corresponding to the fastest shock acceleration—the Bohm diffusive shock acceleration with the scattering mean free path equal to the particle gyroradius. It is generally assumed that the downstream magnetic field is close to the equipartition with the shock internal energy, while the value of upstream magnetic field is less clear. As \( B_u \lesssim B_d \), the total acceleration time is dominated by upstream residence time, so \( t'_{\text{acc}} \simeq \kappa \varepsilon_f / q B_u c \). The maximum energy of accelerated electrons in each region is determined by equating the residence time with the shorter one of the cooling time and the available dynamic time, i.e., \( t'_{\text{acc}} = \min(t'_{\text{cool}}, t'_{\text{dyn}}) \).

The cooling time in the downstream and upstream are, respectively, \( t'_{\text{cool}, d} = 3 \varepsilon_d / (4 \pi \gamma T U_{\text{acc}} + k (\gamma U_f U_d)) \), where \( U_{\text{acc}} \) and \( U_d \) represent the magnetic field energy density in downstream and upstream, respectively, and \( \gamma U_f \) is the photon energy density. In the downstream region, \( U_{\text{acc}} \gg k (\gamma U_f U_d) \), so the maximum electron energy is \( \gamma_{\text{M,d}} = \left( \frac{6 \varepsilon_p B_0}{q B_d} \right)^{1/2} \). In the upstream region, the magnetic field energy density could be lower than \( k (\gamma U_f U_d) \), and in this case, the maximum electron energy is \( \gamma_{\text{M,u}} = \left( \frac{3 \varepsilon_p B_0}{q \kappa_u e \sigma_T U_d} \right)^{1/2} \), where \( k (\gamma U_f U_d) = (\gamma T / \gamma U_f U_d)^{1/2} \).

So \( \gamma_{\text{M,u}} = \left( \frac{3 \varepsilon_p B_0}{4 \kappa_u e \sigma_T U_d} \right)^{2/3} \). As \( B_d \gtrsim B_u \), the electrons radiate more efficiently in the downstream and therefore the relevant maximum Lorentz factor with the observed radiation is \( \gamma = \min(\gamma_{\text{M,u}}, \gamma_{\text{M,d}}) \). Depending on which of \( \gamma_{\text{M,u}} \) and \( \gamma_{\text{M,d}} \) is larger, we divide the discussion into two cases.

1. The \( \gamma_{\text{M,d}} \lesssim \gamma_{\text{M,u}} \) case. The maximum synchrotron photon energy is

\[
h_{\gamma_{\text{syn}, M}} = \frac{0.2294}{4 \pi m_e c} \frac{3 q B_d}{4 \pi m_e c} \frac{\gamma_{\text{M,d}}^2}{\gamma_{\text{M,d}}^2} \Gamma_3 \geq 55 \left( \frac{1}{\kappa_d} \right) \Gamma_3 \text{ GeV},
\]

which is only dependent on the bulk Lorentz factor \( \Gamma \) of the relativistic outflow (0.2294 is the coefficient quoted from Wijers & Galama 1999). If the observed highest energy photon is produced by synchrotron radiation, from \( h_{\gamma_{\text{syn}, M}} \gtrsim \varepsilon_h \), we obtain

\[
\kappa_d \lesssim 0.8 \Gamma_3 \left( \frac{\varepsilon_h}{70 \text{ GeV}} \right)^{-1},
\]

which favors the Bohm diffusive acceleration if the bulk Lorentz factor \( \Gamma \lesssim 1 \) is fawful.

From the precondition, \( \gamma_{\text{M,d}} \lesssim \gamma_{\text{M,u}} \), we obtain a lower limit of the upstream magnetic field in this case,

\[
B_u \gtrsim \left( \frac{4 \kappa_u e \sigma_T U_d \gamma_{\text{M,d}}^{1/2}}{3 q} \frac{6 \varepsilon_f q}{\kappa_u e \sigma_T B_d} \right)^{3/4}
\]

\[
= 500 \left( \frac{\varepsilon_f}{\varepsilon_h} \right)^{3/8} \kappa_u^{3/4} q B_d / \gamma_{\text{M,d}}^{13/4} \frac{1}{\gamma_{\text{M,d}}^{1/2}} \Gamma_3 \frac{5/4}{\varepsilon_h / 2 \text{ MeV}}
\]

2. For the case \( \gamma_{\text{M,u}} \lesssim \gamma_{\text{M,d}} \), the maximum synchrotron photon energy is \( h_{\gamma_{\text{syn}, M}}(\gamma_{\text{M,u}}) = 0.2294 \frac{3 q B_d}{4 \pi m_e c} \gamma_{\text{M,u}} \Gamma_3 \). So from \( h_{\gamma_{\text{syn}, M}}(\gamma_{\text{M,u}}) \gtrsim \varepsilon_h \), we obtain a lower limit of the upstream magnetic field,

\[
B_u \gtrsim \left( \frac{4 \kappa_u e \sigma_T U_d \gamma_{\text{M,d}}^{1/2}}{3 q} \frac{6 \varepsilon_f q}{\kappa_u e \sigma_T B_d} \right)^{3/4}
\]

\[
\approx 600 \left( \frac{\varepsilon_f}{\varepsilon_h} \right)^{3/8} \frac{1}{\gamma_{\text{M,d}}^{1/2}} \frac{1}{(70 \text{ GeV})^{3/4}} \frac{L_{5/4}}{\gamma_{\text{M,d}}^{1/4} \Gamma_3} \left( \frac{\varepsilon_p}{2 \text{ MeV}} \right)^{-1/2} \Gamma_3^{-5/4} \frac{1}{\gamma_{\text{M,d}}^{1/2}} \frac{1}{\varepsilon_h / 2 \text{ MeV}}
\]

Combining this lower limit with the precondition \( \gamma_{\text{M,u}} \lesssim \gamma_{\text{M,d}} \), we find that Equation (15) is applicable only when \( \varepsilon_h \lesssim 55 \kappa_d^{1/8} \Gamma_3 \text{ GeV} \). In both cases, the GRB shells that produce the prompt emission must have a preshock magnetic field greater than \( \sim 500 \text{ G} \) at a radius of \( R \sim 3 \times 10^{14} \text{ cm} \). If the field lines in the expanding shell are frozen and the width of the shell is constant, the components then vary with distance as \( B \propto r^{-1} \) and \( B \sim B_0 \sim r^{-2} \). For an initial magnetic field of \( B_0 \sim 10^{15} \text{ G} \) within a volume of radius of \( 10^6-10^7 \text{ cm} \), the above limit is larger than the \( B_0 \) component, but still within the \( B_0 \) or \( B_d \) component. Of course, the above limit is also consistent with the hypothesis that the upstream magnetic field is significantly amplified by the particle streaming instability (Bell 2004),\(^9\) Interestingly, the shock-compressed upstream magnetic field, \( B \sim 4 \Gamma_3 B_d \gtrsim 1500 \text{ G} \), is similar to the assumed equipartition magnetic field in the downstream (i.e., Equation (1)), which means that the field compression due to the shock is enough to explain the downstream magnetic field.

5. SUMMARY

The single-component spectrum of GRB 080916C from MeV to GeV puts useful constraints on the emission mechanism. We found that the synchrotron mechanism from relativistic electrons is consistent with the observed spectrum, while the simple one-zone electron IC and hadronic models are less viable. In the synchrotron interpretation, the SSC emission is found to be in the KN scattering regime and as a consequence, the IC component is not visible at high energies even though the magnetic field energy density is smaller than that in the relativist electrons, i.e., \( \varepsilon_B < \varepsilon_e \), as obtained in our case. We also suggest a scenario in which such a KN IC emission-dominated regime can explain the low-energy photon spectral index of GRB 080916C.

The delay of high-energy gamma-ray emission relative to the low-energy emission in GRB 080916C is still a mystery in the electron synchrotron scenario. It could be due to that the energy distribution slope \( p \) of electrons during the first time interval (time \( a \)) is rather steep so that the high-energy emission is suppressed or that the emission region has not become transparent for high-energy gamma-rays at early times.

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\(^9\) Recently, Li & Waxman (2006) constrained the preshock magnetic fields of GRB afterglow shocks by synchrotron X-ray afterglows, which also implies that the preshock magnetic fields may be amplified.
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