TeV BLAZAR GAMMA-RAY EMISSION PRODUCED BY A COOLING PILEUP
PARTICLE ENERGY DISTRIBUTION FUNCTION

Ludovic Sauge and Gilles Henri
Laboratoire d’Astrophysique de Grenoble, Université Joseph-Fourier, BP 53, F-38041 Grenoble, France;
ludovic.sauge@obs.ujf-grenoble.fr, gilles.henri@obs.ujf-grenoble.fr
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ABSTRACT

We propose a time-dependent one-zone model based on a quasi-Maxwellian “pileup” distribution in order to explain the time-averaged high-energy emission of TeV blazars. The instantaneous spectra are the result of the synchrotron and synchrotron self-Compton emission of ultrarelativistic leptons. The particle energy distribution function is computed in a self-consistent way, taking into account an injection term of fresh particles, a possible pair creation term, and the radiative cooling of the particles. The source term is not a usual power law but rather a pileup distribution, which can result from the combination of stochastic heating via second-order Fermi processes and radiative cooling. To validate this approach, we have performed time-averaged fits of the well-known TeV emitter Mrk 501 during the 1997 flaring activity period, taking into account the attenuation of the high-energy component by cosmic diffuse infrared background and intrinsic absorption via the pair creation process. The model can reproduce very satisfactorily the observed spectral energy distribution. A high Lorentz factor is required to avoid strong pair production; in the case of smaller Lorentz factor, an intense flare in the GeV range is predicted because of the sudden increase of soft photon density below the Klein-Nishina threshold. The possible relevance of such a scenario is discussed.

Subject headings: acceleration of particles — BL Lacertae objects: individual (Markarian 501) — galaxies: active — galaxies: jets — gamma rays: theory — radiation mechanisms: nonthermal

1. INTRODUCTION

It is now widely accepted that radio-loud active galactic nuclei (AGNs) harbor magnetized accretion-ejection structures involving a supermassive black hole as a central engine. The EGRET experiment aboard the Compton Gamma Ray Observatory discovered more than 80 gamma-ray–emitting AGNs, all of them belonging to the blazar class (nonthermal continuum spectrum, optical polarization, flat radio spectrum, and strong variability in all frequency bands). Some of these objects have been also firmly detected by atmospheric Cerenkov telescopes (ACTs) with an emission above 1 TeV. The two prototypes of TeV blazars are Mrk 421 (Punch et al. 1992) and Mrk 501 (Quinn et al. 1996), objects relatively close to us and at roughly the same distance, $z \sim 0.031$ and 0.034, respectively. Thanks to the development of ground-based gamma-ray astronomy, the sample of TeV emitters is increasing. During the last decade, several ACT teams have reported the detection or the confirmation of new sources: 1ES 1426+428 (Horan et al. 2002; Djannati-Ataï et al. 2002; Aharanion et al. 2002), 1ES 1959+650 (Nishiyama et al. 2000; Aharanion et al. 2003; Holder et al. 2003), 1ES 2344+514 (Catanese et al. 1998), and PKS 2155–304 (Chadwick et al. 1999). A characteristic feature of blazars is the strong nonthermal emission from the radio to the gamma-ray range, attributed to a relativistic jet supposed to be closely aligned with the observer’s light of sight. Their spectral energy distribution (SED) is quite typical and consists of two broad bumps. In the context of the synchrotron self-Compton (SSC) models, the low-energy component, peaking in the X-ray domain for TeV blazars, is commonly attributed to the synchrotron emission of ultrarelativistic particles plunged into a magnetic field. The second component is thought to be the result of the upscattering of the synchrotron photon field by the same population of ultrarelativistic particles via the inverse Compton (IC) mechanism (Jones et al. 1974; Konigl 1981; Ghisellini et al. 1985). This model provides a good framework to explain the correlated variability for the high- and the low-energy components. Even if the spectral properties of these objects seem to be understood, the different models do not discuss the origin and the physical mechanism of particle acceleration. To reproduce the curved shape of the synchrotron and IC spectra on a wide energy domain, several authors have chosen a particle energy distribution function (EDF) parameterized by a simple or a broken power law in a prescribed energy range $[\gamma_{\text{min}}, \gamma_{\text{max}}]$. This choice is purely phenomenological and has no theoretical justification, even if in some special cases of shock acceleration (first-order Fermi process) power-law EDFs are expected (Jones 1994). For example, to reproduce the X-ray synchrotron bump several authors use a simple power law, $n(\gamma) \propto \gamma^{-\alpha}$, $\gamma \in [\gamma_{\text{min}}, \gamma_{\text{max}}]$, but the dynamical range, i.e., the ratio of $\gamma_{\text{max}}/\gamma_{\text{min}}$, is less than 10 (Pian et al. 1998). In this case, it seems more appropriate to consider a quasi-monoenergetic distribution. In this work, we propose another primary type of EDF for emitting particles in order to reproduce the peculiar SED of TeV blazars. We assume that the acceleration mechanism combined with radiative losses or an escape process produces a quasi-Maxwellian or pileup distribution, which is injected into a spherical region in which it cools freely. The effect of cooling is to produce naturally a $E^{-2}$ power law in some limited range of energy. We also take into account the time dependence of the EDF to compare with the observations, considering that the observed spectra are always time-averaged spectra of intrinsically highly variable objects. In § 2, we present the kinetic scenario used to obtain the energy spectrum of...
the particles in a self-consistent way and we briefly describe the emission processes used to reproduce the blazar spectra. Finally, we illustrate our approach in § 3, giving some results of SED fitting before concluding.

2. THE MODEL

2.1. Stochastic Particles Acceleration

In the following, we consider only a homogeneous one-zone model in which all physical quantities are assumed to be averaged over the volume of the emission region. All spatial dependences are dropped from the equations. The particle distribution function $f(p; t)$ is assumed to be isotropic in some frame, called the blob frame, moving relativistically with a bulk Lorentz factor $\Gamma_B$. In this frame, the particle distribution function depends only on the modulus of the momentum $p = |p|$ and the time $t$. For relativistic particles, the energy is given by $E = \gamma m_c c^2 \sim pc$ and the differential number density of pairs $n_\pm(\gamma; t)$ of reduced energy $\gamma$ is related to the EDF $f(\gamma)$ by the usual relation (time is implicit), $dn_\pm = n_\pm(\gamma) d\gamma = 4\pi p^2 f(p) dp \sim 4\pi (mc^2)^2 \gamma^2 f(\gamma) d\gamma$.

We assume that the particles are accelerated stochastically by energy exchanges with resonant plasma waves in a weak turbulent medium. In our model, the acceleration zone must be localized: it could be the basis of a jet, localized reconnection sites, or the interface between a relativistic beam and a confining jet as proposed, for example, by Henri & Pelletier (1991) in the framework of the “two-flow model” (Pelletier 1985; Pelletier & Sol 1992). This insures that the particles will spend only a tiny fraction of time in the acceleration zone before being injected in a larger region, where they cool freely. According to quasi-linear theory, the acceleration process can be described by a diffusion equation in the momentum space, leading to a Fokker-Planck equation. This equation gives the time-dependent evolution of any initial particle density submitted to deterministic continuous energy changes or diffusive Markovian processes. We suppose that the characteristic acceleration timescale is short compared with the other timescales in the problem and we will focus our attention onto the stationary solution $f(\gamma)$ of the Fokker-Planck equation. The diffusion coefficient $D_{\gamma\gamma}(\gamma)$ in phase space can be chosen as a power law in terms of the Lorentz factor $\gamma$ (Lacombe 1977; Henri & Pelletier 1991; Dermer et al. 1996),

$$D_{\gamma\gamma}(\gamma) = D_0 \gamma^r,$$

where $r \in [1, 2]$ is the index of wave turbulent spectra, assumed to be itself a power law (e.g., $r = 5/3$ for a Kolmogorov turbulence, $r = 3/2$ for Kraishman one). The steady state differential energy spectrum resulting from a competitive balance between usual radiative cooling processes and stochastic acceleration is a relativistic Maxwellian function, also called pileup distribution (Schlickeiser 1985; Aharonian et al. 1986; Henri & Pelletier 1991),

$$n(\gamma) \propto \gamma^2 \exp \left[ - \left( \gamma / \gamma_{\max} \right)^{3-r} \right],$$

where $\gamma_{\max}$ is simply the value of the individual Lorentz factor of the particles for which the acceleration time is equal to the cooling time. It corresponds to an energy distribution function of particles that is homogeneous and isotropic in the momentum space with an exponential cutoff at $\gamma_{\max}$. Note that in the case of a power-law distribution function (with spectral index $s > 2$), the enthalpy of the plasma is dominated by the lower bound of the particle energy range $\gamma_{\min}$. For a pileup, particles are mostly concentrated near $\gamma_{\max}$ and the dynamics of the plasma is mainly controlled by the high-energy particles. The inclusion of an escape term will modify the above solution. The model presented here will break down if the escape time is much smaller or much larger than the characteristic acceleration time at the critical Lorentz factor $\gamma_{\max}$. In the first case, acceleration will be much slower than the escape and no relativistic pileup can be formed. In the second case, the relativistic particles will remain a long time before escaping (and cooling) and the emission of the acceleration zone will be important. In the following, we exclude these two cases and we assume that the escape time is comparable to other times at $\gamma_{\max}$, neglecting the emission of the acceleration zone. A proper inclusion of the escape term would modify the solution of the type given by equation (1), but the general shape would be the same; a low-energy part behaving like $\gamma^2$ when the acceleration/diffusion is very fast, followed by an energy cutoff. For the sake of simplicity, we thus use equation (1) and replace the $3 - r$ exponent by 1. The shape of the SED high-energy tail is only weakly dependent on this approximation, and is not strongly constrained by the observations.

2.2. The Cooling Zone

In order to obtain the energy spectrum of emitting particles, we assume that the particles are accelerated as previously described in some localized region and are injected during some time into a spherical zone in which they cool freely. In this zone, we consider the standard kinetic equation in the continuous loss approximation with no escape term. It gives the evolution of the differential energy density of the particles $n_{\pm}(\gamma; t)$ with a Lorentz factor between $\gamma$ and $\gamma + d\gamma$,

$$\frac{\partial}{\partial t} n_{\pm}(\gamma; t) + \frac{\partial}{\partial \gamma} \gamma (\gamma; t) n_{\pm}(\gamma; t) = Q(\gamma; t).$$

The particle source term $Q(\gamma; t)$ will in fact include both the fresh particle injection term $Q_{\text{inj}}(\gamma; t)$ and the production rate $Q_{\text{prod}}(\gamma; t)$ because of pair creation via photon-photon annihilation, which is developed in § 2.4. We take the following approximate form for the injection term:

$$Q_{\text{inj}}(\gamma; t) = \begin{cases} n_0 \gamma^2 \exp (-\gamma / \gamma_{\max}) & \text{if } 0 \leq t \leq t_{\text{inj}}, \\ 0 & \text{otherwise}. \end{cases}$$

The factor $\gamma (\gamma; t)$ in the energy advective part of equation (3) is the continuous particle cooling rate. As mentioned above, charged particles can cool both via the synchrotron process or via the IC scattering of the previous synchrotron radiation field. We can thus write

$$\gamma (\gamma; t) = \gamma_{\text{syn}}(\gamma; t) + \gamma_{\text{ic}}(\gamma; t).$$

2.3. The Radiative Processes

In the following, we detail the equations used to compute the radiative processes. A tilde accent denotes a parameter expressed in the observer frame; otherwise it is in the blob frame.

COOLING PILEUP EDF GAMMA-RAY EMISSION 137
2.3.1. The Synchrotron Emission

Assuming an isotropic particle distribution, the synchrotron cooling rate is given by the well-known formula
\[ \dot{\gamma}_{\text{syn}}(\gamma; t) = -k_{\text{syn}} \gamma^2, \quad k_{\text{syn}} = \frac{4}{3} \frac{\sigma_{\text{Th}}}{m_e c} U_B, \]
(5)
where \( U_B = B^2 / 8\pi \) is the magnetic energy density. The synchrotron emission coefficient \( j_s(\nu) \) is obtained by performing the integration over the whole differential particle density of the mean emission coefficient for a single lepton averaged over an isotropic distribution of pitch angles \( R_{\text{CS}}(\nu) \) (Crusius & Schlickeiser 1986; Ghisellini et al. 1988),
\[ j_s(\nu; t) = \frac{3eB}{4\pi m_e c^2} \int d\gamma n_{\pm}(\gamma; t) R_{\text{CS}}(\gamma), \]
(6)
\[ R_{\text{CS}}(\nu) = 2z \left\{ K_{3/2}(z)K_1(z) - \frac{z}{3} \left[ K_{5/2}(z) - K_1^2(z) \right] \right\}, \]
(7)
with \( z = \nu / \gamma^2 \nu_B \) and \( \nu_B = eB / 2\pi m_e c, K_n \) being the McDonald function of order \( n \). An accurate approximation of the function \( R_{\text{CS}}(\nu) \) is given in the Appendix.

2.3.2. The Inverse Compton Emission

In the same way, the IC scattering cooling rate reads
\[ \dot{\gamma}_{\text{IC}}(\gamma; t) = \int d\epsilon_1 \epsilon_1 \int d\epsilon K_{\text{Jones}}(\epsilon_1, \epsilon, \gamma) n_{\text{syn}}(\epsilon; t), \]
(8)
where we use the Compton kernel \( K_{\text{Jones}} \) computed by Jones (Jones 1968; Blumenthal & Gould 1970) for an isotropic source of soft photons, considering the full Klein-Nishina section in the head-on approximation. More precisely, we have
\[ K_{\text{Jones}}(\epsilon_1, \epsilon, \gamma) = \frac{3}{4} \frac{e \gamma}{\epsilon_1^2} \frac{1}{\epsilon_1} \Omega(\gamma - 1 - q), \]
\[ \Gamma_\epsilon = 4e \gamma, \quad q = \frac{\epsilon_1}{4e \gamma / (\gamma - 1)}, \]
\[ f(q, \Gamma_\epsilon) = \left[ 2 \ln q = (1 + 2q)(1 - q) + \frac{1}{2} \frac{(\Gamma_\epsilon q)^2}{1 + \Gamma_\epsilon q}(1 - q) \right], \]
(9)
(10)
(11)
where \( \Omega(t) \) is the usual Heaviside one-step function and \( \sigma_{\text{Th}} \) is the Thomson cross section. We derive the IC emission coefficient \( j_c(\nu) \) in a similar way. The Compton kernel (eq. [10]) is this time integrated over the synchrotron emission spectrum:
\[ j_c(\nu_1; t) = \frac{h}{4\pi} \epsilon_1 \int \int K_{\text{Jones}}(\epsilon_1, \epsilon, \gamma) n_{\text{syn}}(\epsilon; t) n_{\pm}(\gamma; t) d\gamma d\epsilon, \]
(12)
where \( n_{\text{syn}}(\epsilon) \) is the differential synchrotron photon density and \( \epsilon_1 = h \nu_1 / m_e c^2 \). These equations are integrated numerically following the time evolution of the particle spectrum to find the time-dependent emission spectrum of the plasma.

2.3.3. The Pair Creation Process

Gamma-ray photons produced in the blob can be absorbed by the photoannihilation/pair creation process \( \gamma + \gamma \rightarrow e^+ + e^- \)
(Gould & Schréder 1967a) for which the total cross section is
\[ \sigma(x) = \frac{3\sigma_{\text{Th}}}{16} (1 - x^2) \left( 3 - x^4 \right) \ln \left( \frac{1 + x}{1 - x} \right) - 2x (2 - x^2) \]
(13)
The \( \gamma - \gamma \) annihilation optical depth per unit length \( dl \) reads
\[ \frac{d}{dl} \tau_{\gamma\gamma}(\epsilon_1) = \int \sigma(\beta, \Omega) n_{\text{ph}}(\epsilon_2, \Omega)(1 - \mu) d\epsilon_2 d^2\Omega, \]
(14)
where \( \beta \equiv \beta(\epsilon_1, \epsilon_2, \mu) = \left[ 1 - 2/\epsilon_1 \epsilon_2 (1 - \mu) \right]^{1/2} \) is the velocity of the pairs in the center-of-mass system, \( \epsilon_1 (\epsilon_2) \) is the energy of the high (low) energy photon, and \( \mu \) is the cosine of the collision angle (Coppi & Blandford 1990). For gamma rays in the TeV range, the pair production cross section is maximized when the soft photon energy is in the infrared range,
\[ \lambda = \lambda_{\text{ce}} \sim 2.4 \frac{E_\gamma}{1 \text{ TeV}} \text{ cm}^{-2} \text{ s}^{-1} \text{ m}^{-2} \text{ TeV}^{-1}, \]
(15)
According to the previous relation, we can distinguish two different sources of soft photons able to absorb the high-energy tail of blazars.

First, through intrinsic attenuation gamma-ray photons can interact with photons of the synchrotron bump in the source. Assuming again the synchrotron photon field to be isotropic in the blob frame, the integration over solid angle in equation (14) can be analytically computed. More precisely, one gets
\[ \frac{d}{dl} \tau_{\gamma\gamma}(\epsilon_1) = \frac{1}{c} \int n_{\text{ph}}(\epsilon_2) R_{\text{pp}}(\epsilon_1, \epsilon_2) d\epsilon_2, \]
(16)
where \( R_{\text{pp}}(\epsilon) \) is the angle-averaged pair production rate (cm3 s−1),
\[ R_{\text{pp}}(\epsilon) = c \int \mu_{\text{kin}} d\mu \frac{1 - \mu}{2} \sigma(\beta), \quad \mu_{\text{kin}} = \max (-1, 1 - 2/x), \]
\[ = 3 \frac{\sigma_{\text{Th}}}{4 x^2} \psi \left( \frac{1 + \sqrt{1 - 1/x}}{1 - \sqrt{1 - 1/x}} \right) \Theta(x - 1), \]
(17)
and introducing the function
\[ \psi(u) = -\frac{1}{2} \ln^2 (u) + \frac{2 u (2 + u)}{(u + 1)^2} \left[ \frac{u}{4} - \frac{2 u}{u + 1} + 2 \ln (1 + u) - \frac{1}{2} + \frac{1}{4u} \right] \ln (u) \]
\[ + 2 \text{dilog} (u + 1) - \frac{u}{2} + \frac{1}{2u} - \frac{2}{u + 1} + 1 + \frac{\pi^2}{6}, \]
(18)
with the dilogarithm function as (Abramowitz & Stegun 1965; Gould & Schréder 1967a)
\[ \text{dilog}(x) = - \int_{1}^{x} \frac{\ln (u)}{u - 1} du \]
\[ = - \frac{\pi^2}{6} - \frac{1}{2} \ln^2 (x - 1) + \sum_{k=1}^{\infty} \frac{(-1)^{k+1}}{k^2} (x - 1)^k, \quad x > 1. \]
To compute the photon escape probability, $P_{\text{esc}}(\nu_1; t)$, of a photon of energy $\epsilon_1 = h\nu_1/m_\gamma c^2 > 1$ [or the spectrum attenuation coefficient $C_{\text{abs}}(\nu_1; t)$], we use the approximate expression (Marcowith et al. 1995)

$$P_{\text{esc}}(\nu_1; t) = C_{\text{abs}}(\nu_1; t) = \left(1 - e^{-\tau_{\gamma\gamma}}\right) e^{-\tau_{\gamma\gamma}},$$

(19)

with $\tau_{\gamma\gamma} \sim R d\tau_{\gamma\gamma}/dl$. The factor in parentheses is the usual solution of the transfer equation in the plane-parallel geometry approximation and can approximate the photon escape probability in a spherical source of size $R$. The extra exponential factor in equation (19) has been introduced by Marcowith et al. (1995) to account for the possibility that high-energy photons annihilate outside the source because the soft target photons are not confined in the source as the particles are; rather, their density decreases slowly on a typical length scale equal to the source radius. Second, the high-energy photons can also interact with the photon field of the diffuse infrared background (DIRB; also called CIB for cosmic infrared background) during their travel through the universe from the source to the observer (Gould & Schreder 1967b). Hereafter, we call this effect extrinsic absorption. DIRB is the extragalactic light from the server (Gould & Schreder 1967b). Hereafter, we call this effect

$$C_{\text{abs}}^\text{ext}(\nu, z) = \exp\left[-\tau_{\gamma\gamma}^{\text{ext}}(\epsilon_1, z_0)\right], \quad \epsilon_1 = h\nu/m_\gamma c^2.$$  

(20)

For close sources ($z_0 \ll 1$), the expression (14) gives

$$\tau_{\gamma\gamma}^{\text{ext}}(\epsilon_1, z_0) \sim \frac{cz_0}{H_0} \int_1^\infty d\mu \frac{1 - \mu}{2} \int_{\epsilon_0}^\infty d\epsilon_2 \sigma(\beta) n_{\text{DIRB}}(\epsilon_2),$$

(21)

where $\epsilon_0 = 2/(1 - \mu)\epsilon_1, H_0$ is the Hubble parameter (assumed to be equal to 65 km s$^{-1}$ Mpc$^{-1}$ throughout this paper) and $n_{\text{DIRB}}$ is the density of the DIRB photon field. We have estimated this density by performing a Chebyshev interpolation using the measurements data compiled by Hauser & Dwek (2001), excluding the two points from COBE DIRBE at 60 and 100 $\mu$m. However, note that with or without these points, the resulting absorption coefficients are quite similar above 9 TeV, as shown in Figure 1, and do not change the main results of this work.

2.4. Pair Production Rate and Pair Cascade

As mentioned above, the kinetic equation source term includes the contribution of the population of created particles in the pair production process as calculated above. Note that a hard photon with a reduced energy $\epsilon > 1$ interacts preferentially with a soft photon of energy $\sim 1/\epsilon < 1$ to form a pair $e^+/e^-$ close to the production threshold. Consequently, both particles have a similar energy $\gamma$ and we formally write the conservation of energy as $\epsilon + 1/\epsilon \approx \epsilon = 2\gamma$. Thus, the pair production rate reads

$$Q_{\text{prod}}(\gamma; t) = \frac{dN}{dt d\gamma d\nu} \approx 2\epsilon_{\text{abs}}(\epsilon = 2\gamma; t).$$

(22)

Assuming that the IC emission is isotropic in the plasma rest frame, the differential photon absorption rate density per energy and time unit is given by

$$\dot{n}_{\text{abs}}(\epsilon) = 4\pi m_\gamma c^2 \frac{f_\nu}{h^2 \nu} P_{\text{abs}}(\nu; t),$$

(23)

and we finally obtain

$$Q_{\text{prod}}(\gamma; t) = \frac{8\pi m_\gamma c^2}{h^2} \frac{f_\nu}{\nu} P_{\text{abs}}(\nu; t) \left[\frac{f_\nu}{\nu} P_{\text{abs}}(\nu; t)\right]_{\nu = 2\gamma m_\gamma c^2/h}.$$  

(24)

2.5. Time-averaged Spectra

At time $t$, the whole specific intensity in the plasma rest frame reads

$$I_\nu(\nu; t) \sim R \left[f_\beta(\nu; t) + f_\gamma(\nu; t)C_{\text{abs}}(\nu; t)\right]C_{\text{abs}}^{\text{ext}}(\nu, z_0),$$

(25)

where all parameters and emission coefficients are expressed above. All the physical quantities must be converted from the blob frame to the observer frame, taking into account the Doppler boosting effect and the cosmological corrections according to

$$\tilde{I}_\nu(\tilde{\nu}; t) = \frac{R^2}{d_l^2} \delta_\beta(1 + z_\nu)I_\nu(\nu; t), \quad \nu = \frac{1 + z_\nu}{\delta_\beta} \tilde{\nu},$$

(26)

where $d_l$ is the usual luminosity distance, $\delta_\beta = 1/\Gamma_\beta(1 - \beta_\gamma \cos \theta)$ is the Doppler beaming factor of the source, and $\theta$ is the viewing angle. The observed spectrum is finally obtained by assuming that an observation takes place in the interval $[\tilde{t}_{\text{obs}}, \tilde{t}_{\text{obs}} + \Delta \tilde{t}_{\text{obs}}]$ (time $t = 0$ is related to the beginning
of the injection of fresh particles). The time-averaged spectrum is then

\[
F_\gamma^\text{obs}(\nu) = \langle \tilde{F}_\nu(\nu) \rangle_t = \frac{1}{\Delta t_{\text{obs}}} \int_{t_{\text{obs}}}^{t_{\text{obs}} + \Delta t_{\text{obs}}} \tilde{F}_\nu(\nu, t) \, dt.
\]  

(27)

3. RESULTS AND DISCUSSION

3.1. General Behavior

Our model requires eight parameters. Three of these are related to the properties of the source, namely, the magnetic field strength \(B\), the radius \(R\) (or \(R_{15}\), when expressed in units of \(10^{15}\) cm), and the bulk Doppler factor \(\delta_B\). Three others characterize the injected plasma: these are the characteristic Lorentz factor of the pileup EDF \(\gamma_{\text{max}}\), the number of injected particles that can be characterized by the integrated Thomson optical depth \(\tau_{\text{Th}} = \sigma_{\text{Th}} \int dt \int d\gamma \, Q_{\text{inj}}(\gamma, t)\), and the injection time \(t_{\text{inj}}\). The remaining parameters are the observational ones, \(t_{\text{obs}}\) and \(\Delta t_{\text{obs}}\). In a steady state model, there would only be five parameters, since the last three would be irrelevant. More exactly, the integrated Thomson optical depth \(\tau_{\text{Th}}\) should be replaced by the constant particle optical depth \(\tau_{\text{Th}} = nR \sigma_{\text{Th}}\), where \(n\) is the particle density in the source. For a pileup distribution, and neglecting pair creation, the whole spectrum is entirely characterized by two peak energies and two corresponding fluxes, corresponding to the synchrotron and the IC bumps, respectively. Thus, there would be only one free parameter left, which can be taken, for instance, as the unknown Doppler factor \(\delta_B\). The same spectrum would be obtained by varying \(\delta_B\) and adjusting the other parameters accordingly. Since the TeV emission is dominated by the Klein-Nishina cutoff, the IC peak energy is simply proportional to \(\delta_B \gamma_{\text{max}}^2\), whereas the synchrotron peak energy is proportional to \(\delta_B \gamma_{\text{max}}^2 B\). Thus, the following scaling laws would apply:

\[
\gamma_{\text{max}} \propto \delta_B^{-1}, \quad B \propto \delta_B.
\]  

(28)

The integrated synchrotron luminosity scales as \(N \gamma_{\text{max}}^2 B^2 \delta_B^4\), where \(N = 4\pi R^3 n / 3 \cdot 4\pi \tau_{\text{Th}} R^2 / (3 \sigma_{\text{Th}})\) is the total number of particles in the source. So one gets another scaling law,

\[
R^2 \tau_{\text{Th}} \propto \delta_B^{-4}.
\]  

(29)

The final condition must be determined by the fact that the IC luminosity, \(L_{\text{IC}}\), is directly related to the synchrotron photon density in the source, which is itself directly related to the radius of the source (the synchrotron luminosity does not depend on this radius for a given number of particles). The magnetic energy density scales as \(B^2\), so as \(\delta_B^2\) (eq. [33]), whereas the synchrotron photon energy density scales as \(L_{\text{IC}} \delta_B^{-4} R^{-2}\). If one neglects the Klein-Nishina correction, the ratio of synchrotron to Compton luminosity is simply the ratio of magnetic to soft photon density, and is fixed by the observations. One would thus expect the scaling law

\[
R \propto \delta_B^{-3}.
\]  

(30)

In fact, the real condition is more involved because the Klein-Nishina cutoff strongly diminishes the number of photons effectively available for IC scattering in a way depending on the Doppler factor and the shape of the synchrotron spectrum. However, qualitatively one can always choose the radius of the source to adjust the IC luminosity to the observed value. There

is, however, a limitation to the Doppler factor due to the existence of pair creation process. For small values of \(\delta_B\), the \(\gamma-\gamma\) optical depth can increase so much that it becomes on the order of unity. Then, the IC luminosity stops increasing and instead starts to decrease with decreasing radius because of the \(\gamma-\gamma\) absorption. For a given Doppler factor, there is thus a maximum reachable IC luminosity. Conversely, for a given IC luminosity there is a minimum Doppler factor (which can be of course be 1 in some cases in which pair creation is never important). Note that the variability timescale is also a possible limitation, rapid variability requiring also high Lorentz factors. In principle, the time-dependent model is more constrained. It requires three more free parameters (the injection time and the two observation times), but the entire shape of the synchrotron spectrum is dependent on these times. One can see that by varying \(\delta_B\), the cooling time of particles with the typical energy \(\gamma_{\text{max}} B^4\) varies as \((\gamma_{\text{max}} B^4)^{-1} \propto \delta_B^{-1}\) in the blob frame, so as \(\delta_B^{-2}\) in the observer frame. The shape of the synchrotron spectrum will remain unchanged if one scales all times proportionally to \(\delta_B^{-1}\) in the blob frame, or \(\delta_B^{-2}\) in the observer frame. However, one of these times, namely, the observation time \(\Delta t_{\text{obs}}\), is not a free parameter. Thus, if the model fits the data perfectly well for all values of \(\delta_B\), only one of these values is compatible with the actual value of \(\Delta t_{\text{obs}}\). So, theoretically a unique set of parameters (if any) can fit the data. Of course, things are not so ideal: because of observation error bars, data will be fitted by a set of possible values with a satisfactory \(\chi^2\) test.

3.2. Approximated Analytical Solution of the Kinetic Equation

For illustrative purposes, we develop here the simplest case, in which one can neglect IC cooling in comparison with synchrotron cooling and pair production is unimportant. We can express analytically the general solution of equation (2), which satisfies the boundary condition \(n_{\pm}(\gamma; t = 0) = 0\) for any arbitrary injection term \(Q_{\text{inj}}(\gamma; t)\) by

\[
n_{\pm}(\gamma; t) = \frac{1}{[\gamma(\gamma)]} \int_\gamma^{\infty} d\gamma_0 Q_{\text{inj}}(\gamma_0; t) - \tau(\gamma_0, \gamma),
\]  

(31)

where \(\tau(\gamma_0, \gamma)\) is the energy drift time, i.e., the time needed for a particle of energy \(\gamma_0\) to cool down to energy \(\gamma\),

\[
\tau(\gamma_0, \gamma) = \int_{\gamma_0}^\gamma \frac{du}{\gamma_0(\gamma)}. \quad (32)
\]

For the injection term chosen as in equation (3), equation (31) can be analytically integrated and we finally obtain

\[
n_{\pm}(\gamma; t) = n_0 \frac{\gamma^3_{\text{max}}}{\kappa_{\text{syn}} \gamma_0^2} \left\{ \Gamma \left[ 3, \frac{a(\gamma; t)}{\gamma_{\text{max}}} \right] - \Gamma \left[ 3, \frac{b(\gamma; t)}{\gamma_{\text{max}}} \right] \right\},
\]  

(33)

where \(\Gamma(\cdot, x)\) denotes to the incomplete gamma function. Here the parameters \(a(\gamma; t)\) and \(b(\gamma; t)\) represent, respectively, the lower and upper bounds of the integral (31), where the integrand does not vanish. To evaluate them, we distinguish three time intervals for each value of \(\gamma\). Let us define the parameter \(t_{\text{cool}}(\gamma) = |d\gamma / d\nu|^{-1}/\gamma_{\text{syn}}\). Note that \(t_{\text{cool}}(\gamma) = \tau(\infty, \gamma)\), i.e., it also represents the time spent by an initial infinite energy particle to cool down to \(\gamma\).

1. **Initial stage**, where \(t < \min[t_{\text{cool}}(\gamma), t_{\text{inj}}]\),

\[
b(\gamma; t) = \frac{\gamma}{1 - \kappa_{\text{syn}} \gamma t}, \quad a(\gamma; t) = \gamma.
\]
Particles are still injected at energy $\gamma$ but high-energy particles have not yet cooled down to $\gamma$. Particles injected between $\gamma$ and some finite upper bound contribute to the integral.

2. Cooling stage, where $\max (t_{\text{inj}}, t_{\text{cool}}(\gamma)) < t < t_{\text{inj}} + t_{\text{cool}}(\gamma)$,

$$b(\gamma; t) \rightarrow \infty; \quad a(\gamma; t) = \frac{\gamma}{1 - k_{\text{syn}}\gamma(t - t_{\text{inj}})}.$$  

Particles are no longer injected but some high-energy particles are still cooling down to $\gamma$; particles injected above a finite energy ($>\gamma$) contribute to the integral.

3. Intermediate stage. For intermediate values of $t$, we must distinguish two specific energy ranges. We define a critical energy ($>\gamma_{\text{lim}}$) and the pair production term.

3a. Low energy range, where $\gamma < \gamma_{\text{lim}}$ (or $t_{\text{inj}} < t_{\text{cool}}(\gamma)$),

$$b(\gamma; t) = \frac{\gamma}{1 - k_{\text{syn}}\gamma t}, \quad a(\gamma; t) = \frac{\gamma}{1 - k_{\text{syn}}\gamma(t - t_{\text{inj}})}.$$  

Injection is finished, but very high energy (VHE) particles have not yet cooled down to $\gamma$. Particles injected in some interval above $\gamma$ contribute to the integral.

3b. High energy range, where $\gamma \geq \gamma_{\text{lim}}$ (or $t_{\text{inj}} \geq t_{\text{cool}}(\gamma)$),

$$b(\gamma; t) \rightarrow \infty, \quad a(\gamma; t) = \gamma.$$  

Conversely, VHE particles have time to cool down to $\gamma$ while injection of fresh particles still takes place. All particles injected above $\gamma$ contribute. Note that this is the only stage for which $n(\gamma; t)$ does not depend on time. A steady state is set during this stage (although not in the spectrum because the whole distribution is not steady).

We could also define an end stage, for which $t > t_{\text{cool}}(\gamma) + t_{\text{inj}}$ and where $n_{\infty}(\gamma) = 0$.

Introducing the reduced variables, $\varepsilon = \gamma/\gamma_{\text{max}}$ and $\tau = t/t_{\text{cool}}(\gamma_{\text{max}})$, the previous equations can be collected into the relation

$$n(\varepsilon; \tau) = \frac{n_{0}(\gamma_{\text{max}})}{k_{\text{syn}}\varepsilon^{2}} \varpi_{\text{inj}}(\varepsilon; \tau),$$

where

$$\varpi_{\text{inj}}(\varepsilon; \tau) = \varepsilon^{-2} \Theta \left[ 1 - \varepsilon \max (0, \tau - \tau_{\text{inj}}) \right] \times \left\{ \Gamma \left[ 3, \frac{\varepsilon}{1 - \varepsilon \max (0, \tau - \tau_{\text{inj}})} \right] - \Theta(1 - \varepsilon \tau) \varepsilon \right\}.$$  

An example of a resulting cooling pair EDF at different times is plotted in Figure 2. One clearly sees the initial stage where the EDF is built, the formation of a $\gamma^{-2}$ EDF because of cooling, and the subsequent cooling of the whole distribution after the injection has stopped. As we see in realistic simulations, the shape is, however, strongly modified when taking into account the IC cooling process (which is not simply dependent on the energy because of the Klein-Nishina cutoff) and the pair production term.

![Figure 2](image-url)

**Figure 2.** Example of resulting cooling pileup energy distribution function $\varpi_{\text{inj}}(\varepsilon; \tau)$ for $\tau_{\text{inj}} = 4$ at time $\tau = 1, 2, 3, 4, 5$. Also represented (dashed straight line) is a typical power law of index 2, which results from the radiative cooling of a monoenergetic particle energy distribution.

### 3.3. Application to Mrk 501 Data

#### 3.3.1. Observations

We have applied the model to fit the SED of Mrk 501 during the period of 1997 April when this source experienced an intense period of activity. From this period, we distinguish two different activity states, namely, the “high state” from April 16 and the “medium state” from April 7. Simultaneous data are taken from the French atmospheric Cerenkov telescope CAT for the spectrum in the TeV energy regime (Djannati-Ataï et al. 1999; Barrau et al. 1998). In a first step, we have corrected the high-energy spectra using the attenuation coefficient computed previously. Note that the last corrected data point of the high state may be not meaningful, leading to a concave corrected high-energy spectrum. The most simple and obvious explanation of this problem is an overestimation of the measured high-energy tail of the blazar spectrum or/and of the DIRB density itself. The value of $\Delta t_{\text{obs}}$ is constrained by the observing duration. However, there are some discontinuities in the observing time that make the global observation period length different from the real exposure time. We have assumed that on average our $\Delta t_{\text{obs}}$ corresponds to the global observing time of the BeppoSAX instruments LECS/MECS and PDS ($\sim 40,000$ s for April 16, 37,845 s for April 7).

#### 3.3.2. Solutions without Pair Creation

We can reproduce the data with parameters corresponding to almost no pair creation. The parameters used to fit the data are reported in Table 1 and the resulting synthetic SEDs are plotted in Figure 3. The fits are quite satisfactory for the high-energy part of the SED. One-zone models are appropriate only to reproduce this high-energy, variable part. They cannot account for the radio emission produced at much larger scales, where the whole jet contributes, possibly being the superposition of many successive flaring events. In some range of Lorentz factor, a steady state solution corresponding to a $\gamma^{-2}$ power law (see eq. [33] and case 3b, above) can be observed, giving a $\nu^{-1/2}$ synchrotron flux index. This corresponds to the main part of the April 16 spectrum and the low-energy part of the April 7 spectrum. Above this range, the spectrum is modified by a factor $t_{\text{cool}}/t_{\text{obs}}$ because the particles have cooled before the end.
of the observation. This produces a steepening of the spectrum by $\Delta \alpha = 0.5$ and explains the flat high-energy part of the synchrotron spectrum. The position of the spectral break is thus directly determined by the ratio between the cooling and the observation time. We obtain values of the magnetic field and of the transverse radius of the source ($R_{15}$) by the ratio between the cooling and the synchrotron spectrum. The position of the spectral break is thus determined by the ratio $B/R_{15}$. This produces a steepening of the spectrum due to the pair creation process in the spherical blob, while the dashed curves show the unabsorbed spectra. Open squares show data corrected from the DIRB attenuation.

### Table 1

| State                     | $B$  | $R_{15}$ | $\delta_B$ | $\gamma_{\text{max}}$ | $\tau_{\text{Th}}$ | $\lambda_{\text{inj}}$ | $\lambda_{\text{obs}}$ | $\Delta\lambda_{\text{obs}}$ | $\lambda_{\text{var, min}}$ |
|---------------------------|-----|----------|-------------|------------------------|---------------------|---------------------|---------------------|-----------------------------|-----------------------------|
| High state                | 0.077 | 0.65     | 25          | 1.26                   | 58.5                | 9.4                 | 3.29 ($= 0.35 \lambda_{\text{inj}}$) | 39.7                        | 15                          |
| Medium state              | 0.075 | 1.75     | 25          | 1.29                   | 8.95                | 40                  | 35 ($= 0.875 \lambda_{\text{inj}}$) | 37.8                        | 40                          |
| “GeV flaring state”       | 0.047 | 1.06     | 16          | 2                      | 144.1               | 24.4                | 8.54 ($= 0.35 \lambda_{\text{inj}}$) | 102                         | 38                          |

In common with $B$ and $\delta_B$, $\gamma_{\text{max}}$ is kept constant between the two different states. In light of our scenario, the “medium state” spectrum could be due to a previous ejection observed in a later stage (with respect to the injection time) than the April 16 one and in a much larger part of the jet.

We can also estimate the minimum variability timescale derived from standard causality arguments and given by $\lambda_{\text{var}} > \lambda_{\text{var, min}} = 555R_{15}(1+z_s)^{-8}\Delta t_B$ minutes. Its value is reported in the last column of Table 1 and is equal to 15 minutes for the high state and roughly 40 minutes for the medium state when the size of the source is much larger as mentioned above. Note, however, that the injection time is much larger than the above values.

Simulated light curves corresponding to April 16 parameters are shown in Figure 4. The curves are calculated for the three energy ranges of BeppoSAX and for VHE instruments above 250 GeV. Lags between various energy bands are clearly visible on the curves. Remarkably, the high-energy photons lead the soft-energy ones in the X-ray synchrotron component, but the high-energy gamma-ray curve lags behind the synchrotron component. This is due to the fact that when particle distribution cools, the photon density below the Klein-Nishina limit increases. This makes the IC emission level keep rising even after injection has stopped. Precise comparisons with observed light curves have not been made because the statistics are too poor to allow a meaningful analysis. However, we can see qualitatively that complex temporal effects can arise from a time-dependent simulation, leading to apparently contradictory behaviors. Note that whereas the presented simulations span the whole temporal interval of BeppoSAX observations, high-energy observations have been performed only during a limited time within this interval. If they have taken place around the maximum of the light curve, few variations are expected.

Of course, a more realistic model could include a succession of different flares as well as a variation of the injection characteristic energy. This is, however, beyond the scope of this paper, which only aims at showing that pileup particle
energy distribution can well explain the high-energy shape of TeV blazar spectra.

3.3.3. Solutions with Pair Creation

Interestingly, we also found possible solutions with a much lower bulk Doppler factor and a correspondingly lower magnetic field, a higher particle Lorentz factor, and a larger radius. Examples of such parameters are shown in Table 2 and the corresponding spectrum is displayed in Figure 5. In this regime, pair production can be important. Although the IC flux should be much larger, the $\gamma \gamma$ absorption reduces effectively the luminosity so that it can be compatible with the observations. The intense creation of new particles produces more synchrotron photons, which accelerates the cooling. This phenomenon also amplifies the effect of the increase of the Klein-Nishina threshold energy. This leads eventually to a catastrophic pair production/cooling process, producing a strong flare at GeV energy and in low-energy X-rays. This explains the bump occurring in the GeV/sub-TeV energy range (and to a lesser extent in the radio submillimeter range). The flare is also more clearly visible on the light curves (Fig. 6), appearing as a very sharp flare in some energy ranges. Although the relevance of such a scenario is not clear, it may be possible that such events could explain the most rapid variations in the TeV light curve.

### Table 2

| Coefficient | Value          |
|-------------|----------------|
| $a_0$       | $0.201447 \times 10$ |
| $a_1$       | $0.344606$      |
| $a_2$       | $-0.429682$     |
| $a_3$       | $0.273331 \times 10^{-2}$ |
| $a_4$       | $0.966844$      |
| $a_5$       | $0.964518$      |
| $\chi^2$   | $1.27 \times 10^{-5}$ |

Note.—The rms error is 0.05%.
curves. In principle, the GeV flare would have been observable by the EGRET instrument. However, because of the small number of simultaneous EGRET/TeV observations and the briefness of the event, it could have escaped detection. We note also that for different value of the parameters it is possible that the X-ray flare shifts toward lower energy, disappearing from the X-ray data. This could be related to the “orphan flare events” observed in some occasions (Krawczynski et al. 2004).

4. CONCLUSION

This paper shows that the high-energy spectra of TeV blazars can be well reproduced by a cooling pileup EDF. This offers an alternative to the narrow power-law injection terms often used in the literature, whose justification is unclear. Inclusion of time-dependent effects permits us to reproduce the main features of both the light curves and the time-averaged spectra. For a given SED shape, the parameters of the model are fully constrained. For Mrk 501, different states could be the result of the time variation of the transverse radius of the source $R$, the quantity of injected leptons via $\gamma_{\text{inj}}$, and the observational parameters with respect to the injection time. However, this one-zone model shares a common issue with other homogeneous models in that it requires high values of the bulk Lorentz factor to avoid a strong gamma-ray absorption, even in the case of the optically thick solution (see § 3.3.3). These high Lorentz factors appear to be inconsistent with those deduced from FR I/BL Lac object unification models (Urry & Padovani 1995; Chibberge et al. 2000). They are also difficult to reconcile with the absence of superluminal motion at TeV scales and the relatively low brightness temperature of TeV blazars (Edwards & Piner 2002; Piner & Edwards 2004). We will argue in a future work that inhomogeneous models could fix this issue.

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APPENDIX

AN ACCURATE ANALYTIC APPROXIMATION OF PITCH ANGLE AVERAGED SYNCHROTRON Emitted POWER

As mentioned in § 2.3.1, the $R_{\text{CS}}(x)$ function (see eq. [7]) results from the monochromatic emitted power averaged on a population of particles with randomly distributed pitch angle. It mathematically reads (Crusius & Schlickeiser 1986)

$$R_{\text{CS}}(x) = \frac{1}{2} \int_0^{\pi} d\theta \sin^2 \theta F_{\text{syn}} \left( \frac{x}{\sin \theta} \right).$$

(A1)
where $F_{\text{syn}}(x)$ is the usual synchrotron fundamental kernel (Blumenthal & Gould 1970),

$$F_{\text{syn}}(x) = x \int_{x}^{\infty} dz K_{5/3}(z).$$

(A2)

For $x \ll 1$, an approximate expression of $F_{\text{syn}}$ is

$$F_{\text{syn}}(x) \approx \frac{4\pi}{\sqrt{3\Gamma(1/3)}} \left(\frac{x}{2}\right)^{1/3}, \quad x \ll 1,$$

(A3)

and we immediately obtain the relevant expression for $R_{\text{CS}},$

$$R_{\text{CS}}(x) = \frac{1}{2} \int_{0}^{\pi} d\theta \sin^{2}\theta F_{\text{syn}}\left(\frac{x}{\sin \theta}\right) = \frac{2^{1/3}}{5} \Gamma^{2}(1/3)x^{1/3} \approx 1.808418x^{1/3}.$$  

(A4)

Conversely, for large argument $(x \gg 1)$ the asymptotic development of $F_{\text{syn}}$ is $(\pi x/2)^{1/2} e^{-x}$ and thus

$$R_{\gg}(x) = \frac{1}{2} \int_{0}^{\pi} d\theta \sin^{2}\theta F_{\text{syn}}\left(\frac{x}{\sin \theta}\right) = \frac{1}{2} \left(\frac{\pi x}{2}\right)^{1/2} \int_{0}^{\pi} d\theta \sin^{3/2}\theta e^{-\pi x/\sin \theta}.$$  

(A5)

The preceding integral can be rewritten as

$$I(x) = \int_{0}^{\pi} d\theta \exp \left[-f_{\theta}(\theta)\right], \quad f_{\theta}(\theta) = \frac{x}{\sin \theta} - \frac{3}{2} \ln(\sin \theta).$$

(A6)

Noting that the exponential argument is symmetric around $\theta = \pi/2$ and therefore on the integration range, we set $\theta = \pi/2 + \varphi$ and make a Taylor expansion around $\varphi,$

$$f_{\theta}(\varphi) = \frac{x}{\cos \varphi} - \frac{3}{2} \ln(\cos \varphi) \approx x + \frac{\varphi^{2}}{2} \left(x + \frac{3}{2}\right) + O(\varphi^{3}).$$

(A7)

Then integral (A6) can be rewritten as

$$I(x) \approx e^{-x} \int_{-\pi/2}^{\pi/2} 2d\varphi e^{-\varphi^{2}/2\sigma^{2}}, \quad \text{with} \quad 1/\sigma^{2} = x + \frac{3}{2},$$

(A8)

which is a standard Gaussian integral. Noting that the larger is $x$ the sharper is the integral argument, the lower and upper bounds can be extended to infinity as $x \gg 1$ (error is less than $10^{-10}$ for $x \geq 3$ and). We finally obtain

$$R_{\gg}(x) = \frac{\pi}{2} \sqrt{\frac{2x}{2x + 3}} e^{-x}.$$  

(A9)

For intermediate values, we perform a polynomial fit of the form

$$R_{\text{int}}(x) = a_{0}(x^{a_{1}} + a_{2}x^{2a_{1}} + a_{3}x^{3a_{1}}) \exp(-a_{4}x^{a_{5}}),$$

(A10)

where the coefficient $a_{i}$ is obtained from least-squares fitting and is given in Table 2. In this regime, rms error is on the order of 0.05%.

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