THE BULK LORENTZ FACTOR CRISIS OF TeV BLAZARS: EVIDENCE FOR AN INHOMOGENEOUS PILEUP ENERGY DISTRIBUTION?

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ABSTRACT

There is growing evidence that the estimations of the beaming Doppler factor in TeV BL Lac objects based on the synchrotron self-Compton (SSC) models are in strong disagreement with those deduced from the unification models between blazars and radio galaxies. When corrected from extragalactic absorption by the diffuse infrared background (DIRB), the SSC one-zone models require a very high Lorentz factor (around 50) to avoid strong \( \gamma \-\gamma \) absorption. However, the statistics on beamed versus unbeamed objects, as well as the luminosity contrast, favors a much lower Lorentz factor, on the order of 3. In this paper, we show that for the special case of Markarian 501, the need for a very high Lorentz factor is unavoidable for all one-zone models in which all photons are assumed to be produced at the same location at the same time. Models assuming a double structure with two different beaming patterns can partially solve the problem of luminosity contrast, but we point out that they are inconsistent with the statistics on the number of detected TeV sources. The only way to solve the issue is to consider inhomogeneous models, in which low-energy and high-energy photons are not produced at the same place, allowing for much smaller Lorentz factors. This approach implies that the jet is stratified, but also that the particle energy distribution is close to a monoenergetic one and that pair production is likely to be significant. The implications for relativistic jet physics and the particle acceleration mechanism are discussed.

Subject headings: BL Lacertae objects: general — galaxies: active — galaxies: jets — gamma rays: theory — radiation mechanisms: nonthermal

1. INTRODUCTION

It is now acknowledged that the blazar phenomenon is due to the relativistic Doppler beaming of the nonthermal emission taking place in radio-loud active galactic nuclei (AGNs) in which the jet axis is closely aligned with the observer’s line of sight. These objects exhibit a significant level of optical polarization, a flat radio spectrum, a strong variability in all frequency bands, and a very broad spectral energy distribution (SED), ranging from the radio to the extreme gamma-ray bands. The SED consists typically of two broad components. In the so-called synchrotron self-Compton process (SSC) model, the lowest energy hump is attributed to the synchrotron emission from relativistic electrons and/or positrons and the highest one is attributed to the effect of the inverse Compton (IC) process of the same charged particles on the synchrotron photons and/or external photons. The blazar class of objects includes both the flat spectrum radio quasars (FSRQs) and the BL Lac sources, depending, respectively, on the existence or the lack of detectable emission lines in their spectra. Following Chiaberge et al. (2000), one can define two classes of BL Lac objects (which are most probably two extreme cases in a continuous distribution): the LBL, or “red” BL Lac objects, for which the synchrotron component peaks in the far-IR to optical and the IC component peaks in the MeV–GeV range, and the HBL, or “blue” BL Lac objects, for which the synchrotron component peaks in the UV–X-ray range and the IC component peaks above 10 GeV. The most extreme objects observed up to now are those whose nonthermal emission extends up to the TeV range, the so-called TeV blazars. The two main prototypes are Mrk 421 (Punch et al. 1992) and Mrk 501 (Quinn et al. 1996), two radio-loud AGNs relatively close to us and roughly at the same distance, \( z_s \approx 0.031 \) and 0.034, respectively. Five other TeV detections have been repeatedly detected (1ES 1959+650 and PKS 2155-304; Aharonian et al. 2005a; 1ES 1426+428 and PKS 2005-489; Aharonian et al. 2005b; and 1ES 2344+514; Aharonian et al. 2004). All of them are BL Lac objects, although it is not yet clear whether only BL Lac objects emit TeV radiation or if this is due to a selection effect. As a matter of fact, BL Lac objects appear to be much more numerous than quasars, and the closest blazars all belong to this class. A high-sensitivity threshold would strongly bias the detection toward the closest sources. Furthermore, it is well known that TeV photons are absorbed by the diffuse infrared background (DIRB) to create electron-positron pairs, and it is not obvious whether even the closest quasar, 3C 273 (\( z_s \approx 0.158 \)), is detectable in the TeV range.

One-zone SSC models assume that highly relativistic particles are injected in a spherical zone, where they cool by emitting synchrotron radiation and by the inverse Compton process. The models require that the source radius, the magnetic field, as well as the density and the energy distribution of the emitting relativistic particles, be specified. The latter is most often assumed to be a power law or a broken power law (e.g., Marscher 1983; Tavecchio et al. 1998). It turns out, however, that the computation of emitted radiation is not compatible with the hypothesis of a static source, because in most cases the photon density would be so high that all TeV photons should be absorbed to form electron-positron pairs. Furthermore, the time variability is so short (down to 15 minutes in some cases; Gaidos et al. 1996) that it must be considered incompatible with a spherical static source by virtue of the causality argument. This leads one to
assume that the source is moving with a relativistic bulk velocity \( v = \beta c \). The effect of relativistic bulk motion is entirely described by the Doppler beaming factor \( \delta = 1/\Gamma(1 - \beta \mu) \), where \( \Gamma = (1 - \beta^2)^{-1/2} \) is the usual Lorentz factor and \( \mu = \cos \theta \) is the cosine angle of the jet according to the observer’s line of sight. The Doppler effect shifts all frequencies by a factor \( \delta \) and all specific intensities by a factor \( \delta^2 \). So the actual photon density in the jet frame is much lower than what would be deduced for a static source. The relativistic motion has been invoked for a long time (Rees 1966) to solve a similar issue for radio emission of quasars. Namely, the brightness temperature is so high that for a static source, the relativistic lepton-emitting synchrotron radiation should have cooled immediately through the so-called inverse Compton catastrophe (Rees & Simon 1968). Again, the relativistic motion can account for this phenomenon, because the actual photon density in the jet frame is much lower when the Doppler amplification is taken into account. This beautiful theoretical explanation was later confirmed by the discovery of superluminal motion, which requires Lorentz factors at least as great as the observed apparent reduced velocity \( \beta_{\text{app}} = v_{\text{app}} / c \) (for a review, see Zensus 1997). For \( \mu \approx 1 \), corresponding to \( \theta \approx 1/\Gamma \), one has \( 1 \leq \delta \leq 2 \Gamma \), whereas \( \delta \approx 1/\Gamma \) outside this interval. This means that for a few beamed Doppler samples of boosted sources, one expects a large number of unbeamed and unamplified counterparts. It is natural to think that the unboosted counterparts of bright quasars are the weaker radio galaxies, whose jet is thought to make a larger angle with the line of sight. In particular, it has been proposed that the unboosted counterparts of BL Lac objects could be a subclass of radio galaxies, the so-called Fanaroff-Riley I (FR I) radio galaxies (Urry & Padovani 1995). These are characterized by a rather faint, weakly beamed, and core-brightened radio jet. Statistical studies of radio and X-ray AGN samples have confirmed the possibility of such an association. The inferred beaming factors seem to imply a rather modest value of the bulk Lorentz factor, of about 3. However, the modeling of SSC radiation by one-zone models requires much higher values: according to some studies, they range from 10 to 50 (Tavecchio et al. 1998; Konopelko et al. 2003; Saugé & Henri 2004; Krawczynski et al. 2001). The highest value seems to be needed when one takes properly into account the extragalactic absorption. The problem is further complicated by the absence of clear superluminal motion in TeV blazars, together with a rather modest brightness temperature, which also implies a low Lorentz/Doppler factor (Edwards & Piner 2002; Piner & Edwards 2004). All these contradictory facts lead to what we call here the “bulk Lorentz factor crisis of TeV blazars.”

The aim of this paper is first to ascertain this crisis. We first show that all one-zone SSC models imply high Lorentz factors, based solely on the argument of \( \gamma \gamma \) absorption and disregarding any variability argument. Then we review arguments for the low Lorentz factor, based on general geometric properties of the Doppler boosting. We show that the explanations, based on two different structures and with a possible deceleration of a fast spine responsible for TeV emission, cannot satisfactorily account for the statistics on TeV blazars. We argue that the best solution is to admit the low Lorentz factor constraint, abandoning the one-zone assumption. We show that this conclusion has important consequences regarding the jet physics and the particle acceleration mechanism.

2. THE CASE FOR A HIGH LORENTZ FACTOR

In the following, we develop the need for high Lorentz factors for one-zone models, based on the fewest theoretical assumptions and relying only on observational data. We assume only that the SSC process is at work, in addition to the usual assumptions of one-zone models: that the relativistic particles are injected in a spherical homogeneous “blob” of radius \( R \), moving at a relativistic velocity characterized by the Lorentz factor \( \Gamma \) and a corresponding Doppler factor \( \delta \). The blob is filled with a tangled magnetic field of constant strength \( B \). We label all quantities expressed in source rest frame with a star and leave quantities in observer’s frame unlabeled. All energies are expressed in reduced units of \( m_{\text{ec}}^2 \). Throughout this paper, we express the Hubble parameter by \( H_0 = 100 \ h \ \text{km} \ \text{s}^{-1} \ \text{Mpc}^{-1} \), assuming \( h \) to be equal to \( h = 0.65 \).

2.1. The Synchrotron and IC Differential Luminosity

Inspection of the TeV blazars’ spectra shows that the IC spectra reach their maximum luminosity at some peak energy \( \varepsilon_c \), which is on the order of \( 10^9 \) for TeV photons. This energy corresponds to an energy \( \varepsilon_c^* = \varepsilon_c \delta^{-1} \) in the blob frame. We consider only the particles emitting this typical energy via the IC process, which have a typical individual Lorentz factor (in the blob frame) \( \gamma_c \), which must be greater than \( \varepsilon_c^* \). We then define another typical energy \( \varepsilon_s \), which is emitted by synchrotron process by the same particles. It can be expressed in the blob frame as \( \varepsilon_s^* = (B/B_0)\gamma_s^2 \), where \( B_0 = 3B/2 \) and \( B = 2\pi m_{\text{ec}}^2 c^3 \varepsilon_c / 4.41 \times 10^{13} \ \text{G} \) is the usual “QED critical magnetic field strength.” One has then \( \varepsilon_s^* = \varepsilon_c^* \delta (B/B_0) \gamma_s^2 \). Synchrotron spectra of TeV blazars typically peak in the 1–100 keV range, so that \( \varepsilon_s^* \sim 10^{-2} \) to \( 10^{-1} \). Synchrotron photons are up-scattered at high energy via the IC process. It has been stressed by various authors that the collisions between the most energetic particles and the peak synchrotron photons take place in the Klein-Nishina regime, that is, \( \gamma_s^* \gamma_c^* \gg 1 \) (Fig. 1). In this condition, the particle (electron or positron) gives up all of his energy in a single interaction. It follows that \( \gamma_c^* \sim \varepsilon_s^* \). This gives an estimate of the magnetic field strength,

\[
B = B_0 \left( \frac{\varepsilon_s^*}{\varepsilon_c^*} \right)^{1/2} \delta B_0 \frac{\varepsilon_s^*}{\varepsilon_c},
\]

which is only valid in the Klein-Nishina scattering regime.

The differential synchrotron luminosity \( L_{\text{s},s} \) per unit reduced energy emitted by a population of particles of energy \( \gamma \) at the energy \( \varepsilon_s \), with differential energy number of particles \( dN/d\gamma \) reads

\[
L_{\text{s},s}(\varepsilon_s) = \frac{dL_s}{d\varepsilon} (\varepsilon_s) = \delta^3 \frac{dL_s^*}{d\varepsilon^*} (\varepsilon_s^*) = \delta^3 \frac{dL_s^*}{d\gamma^*} \frac{d\gamma}{d\gamma^*} \frac{d\gamma^*}{d\varepsilon^*}.
\]

The total power lost per particle of energy \( \gamma \) is given by the well-known relation \( dL_s^*/dN = (4/3) c \sigma_{\text{th}} \gamma^2 W_B \), and we obtain

\[
L_{\text{s},s}(\varepsilon_s) = \delta^4 \frac{4}{3} c \sigma_{\text{th}} W_B \frac{\varepsilon_s^3}{2\varepsilon_s} \frac{dN}{d\gamma},
\]

where \( W_B = B^2/8\pi \) is the usual magnetic energy density. Combining with equation (1), we can write

\[
L_{\text{s},s}(\varepsilon_s) = \delta^4 \frac{1}{6\pi} c \sigma_{\text{th}} B_0 \frac{\varepsilon_s^3}{2\varepsilon_s} \frac{dN}{d\gamma}.
\]
density. However we have to take into account that the Klein-Nishina cutoff reduces the effective energy density available for IC scattering. We thus define a new characteristic energy, corresponding to the synchrotron photon energy at the limit between the Thomson and Klein-Nishina regime for particles with an energy $\gamma_c$. This energy is $\epsilon_i^* = 1/\gamma_c$; i.e., in the observer’s frame

$$\epsilon_i = \delta^2 / \epsilon_c.$$  

(5)

Photons of this energy will also be the main contributors to absorbing photons of energy $\epsilon_c$ to create electron/positron pairs. If we neglect the Klein-Nishina contribution above $\epsilon_i^*$, the total power lost per particle of energy $\gamma$ can be written $dL_{ic}/dN = (4/3)c\sigma_{th}\gamma^2 W_{\text{ph}}^{\text{eff}}$, where

$$W_{\text{ph}}^{\text{eff}} = \frac{L_{\text{ph}}^{\text{eff}}}{4\pi R^2 c} = \frac{1}{4\pi R^2 c} \int_{\epsilon_i}^{\epsilon_c} d\epsilon^* \frac{dL_s}{d\epsilon^*}.$$  

(6)

For a power-law spectrum $\nu F_\nu \propto \nu^\beta$ with $\beta > 0$, a simple calculation gives

$$W_{\text{ph}}^{\text{eff}} \approx \frac{1}{4\pi R^2 c} \frac{\delta^2}{\beta_{\text{ph}}} L_{\text{ph}}(\epsilon_i).$$

The coefficient $\beta$ can be replaced by another numerical coefficient close to 1 as long as the $\nu F_\nu$ spectrum is growing with energy. The differential IC luminosity reads then

$$L_{\text{ic}}(\epsilon_c) = \delta^{-1} \beta^{-1} \frac{\sigma_{\text{th}}}{3\pi R^2} \epsilon_c L_{\text{ph}}(\epsilon_i) \frac{dN}{d\gamma},$$  

(7)

Comparing equations (4) and (7) we can now estimate the radius of source $R$ as a function of observed luminosities and the unknown Doppler factor:

$$R = \delta^{-2} \frac{3 \epsilon_c}{2 \pi m_e c^2 \epsilon_i} \frac{\epsilon_c}{\beta_{\text{ph}}} \left[ \frac{L_{\text{ph}}(\epsilon_i) L_{\text{ph}}(\epsilon_c)}{L_{\text{ic}}(\epsilon_c)} \right]^{1/2}.$$  

(8)

We now use this radius estimate to compute the $\gamma\gamma$ optical depth for the photons of energy $\epsilon_c$.

2.2. The $\gamma\gamma$ Photon Opacity

As we mentioned above, gamma-ray photons of energy $\epsilon_i^*$ are mainly absorbed by photons of energy $\epsilon_i^* - 1 = \epsilon_i^*$, creating pairs. So the same soft photons control both the amount of IC process and the absorption of IC photons. The absorption probability (or opacity) per unit path length of a photon of energy $\epsilon_i^* \gg 1$ due to pair production in the case of a power-law SED is given approximately by

$$\tau_{\gamma\gamma}(\epsilon_i^*) = \frac{d}{dz} \Gamma(\epsilon_i^*) = \alpha_{\gamma\gamma} \sigma_{\text{th}} \epsilon_i^* n(\epsilon_i^*),$$

where $\Gamma(\epsilon_i^*)$ is the free mean path of the photon and $n(\epsilon_i^*)$ the differential photon density per unit of reduced photon energy $\epsilon_i^*$. In the framework of one-zone model, the typical interaction scale is of the order of the size of source $R$. It follows that the typical $\gamma\gamma$ optical depth can be written

$$\tau_{\gamma\gamma}(\epsilon_i^*) = \frac{R}{\Gamma_{\gamma\gamma}(\epsilon_i^*)} = \alpha_{\gamma\gamma} R \sigma_{\text{th}} \epsilon_i^* n(\epsilon_i^*).$$

The function $\alpha_{\gamma\gamma}$ (Svensson 1987; Coppi & Blandford 1990) depends on the index $\beta$ of the power law of the spectral soft photon density expressed in $\nu F_\nu \propto \nu^\beta$ form. A commonly used value of $\alpha_{\gamma\gamma}(\beta)$ is 0.2 or 0.25. More precisely, we have (Svensson 1987)

$$\alpha_{\gamma\gamma}(\beta) = \frac{4}{6} \frac{\Gamma^2(2 - \beta) 44 - \beta(41 - \beta)(12 - \beta)}{(4 - \beta)(3 - \beta)}.$$  

(9)

The differential energy density number of particle is given as a function of the differential luminosity

$$n(\epsilon_i^*) = \frac{L_{\text{ph}}(\epsilon_i^*)}{4\pi m_e c^3 R^2 \epsilon_i^*},$$  

(10)

so we get finally the optical depth as a soft compactness at the energy $\epsilon_i^*$:

$$\tau_{\gamma\gamma}(\epsilon_i^*) = \alpha_{\gamma\gamma} \frac{\sigma_{\text{th}} L_{\text{ph}}(\epsilon_i^*)}{4\pi m_e c R} = \delta^{-3} \alpha_{\gamma\gamma}(\beta) \frac{\sigma_{\text{th}} L_{\text{ph}}(\epsilon_i^*)}{4\pi m_e c R}.$$  

(11)

Using our estimate on the source radius $R$ from equation (8), we obtain

$$\tau_{\gamma\gamma}(\epsilon_c) = \delta^{-1} \alpha_{\gamma\gamma}(\beta) \frac{\sigma_{\text{th}} m_e c^{1/2}}{6 \epsilon_c} \frac{\epsilon_c}{\beta_{\text{ph}}} \left[ L_{\text{ph}}(\epsilon_i^*) L_{\text{ph}}(\epsilon_c) \right]^{1/2},$$  

(12)
where we introduce the modified function $\tilde{\alpha}_{\gamma\gamma}(\beta)$ as $\tilde{\alpha}_{\gamma\gamma}(\beta) = \alpha_{\gamma\gamma}\sqrt{\beta}$. Values of $\tilde{\alpha}_{\gamma\gamma}$ and $\alpha_{\gamma\gamma}$ for some $\beta$ are tabulated in Table 1.

### Table 1

| $\beta$  | 0   | 1/2 | 1   | 4/3 |
|----------|-----|-----|-----|-----|
| $\alpha_{\gamma\gamma}(\beta)$ | 0.122 | 0.236 | 0.583 | 1.397 |
| $\tilde{\alpha}_{\gamma\gamma}(\beta)$ | ... | 0.043 | 0.583 | 1.613 |

#### 2.3. Constraints on the Local Synchrotron Spectral Shape

Equation (2) shows that if we are able to measure the position in frequency and flux of both the synchrotron and the IC peak, then we can evaluate the optical depth to $\gamma-\gamma$ absorption at the IC peak as a function of the assumed Doppler factor value. This optical depth is controlled by the synchrotron luminosity at the frequency $\varepsilon_\gamma = \delta^2/\varepsilon_c$. We can use this relation either by assuming some Doppler factor and evaluating the optical depth, or by limiting the value of $\tau_{\gamma\gamma}$ and hence the value of $\delta$. We can define $r_{\text{max}}$, the Compton dominance parameter, as the ratio of the peak of the IC luminosity’s peak to that of the synchrotron:

$$r_{\text{max}} = \frac{\nu_c F_c(\nu_c)}{\nu_\gamma F_\gamma(\nu_\gamma)} = \frac{\varepsilon_\gamma L_{\gamma\gamma}(\varepsilon_\gamma)}{\varepsilon_c L_{\gamma\gamma}(\varepsilon_c)}.$$  \hspace{1cm} (13)

We can rewrite equation (12) to express the luminosity at $\varepsilon_\gamma = \delta^2/\varepsilon_c$ as a function of the optical depth and the $r_{\text{max}}$ parameter. We finally obtain

$$\varepsilon_t L_{\gamma\gamma}(\varepsilon_t) = \frac{\delta^4}{r_{\text{max}}} \left[ \frac{\tau_{\gamma\gamma}(\varepsilon_c)}{\tilde{\alpha}_{\gamma\gamma}(\beta)} \frac{6eh}{\sigma_{\text{Th}} m_e c^{1/2} \varepsilon_c^{\gamma \gamma}} \varepsilon_c^{\gamma \gamma} \right]^2.$$  \hspace{1cm} (14)

Equations (5) and (14) can be considered as a system of two parametric equations of the curve, giving $\varepsilon_t L_{\gamma\gamma}(\varepsilon_t)$ as a function of $\varepsilon_t$. Eliminating $\delta$ between the two previously cited equations, one gets the following expression:

$$\varepsilon_t L_{\gamma\gamma}(\varepsilon_t) = \frac{\varepsilon_t^2}{e^{-\varepsilon_t}} \left[ \frac{\tau_{\gamma\gamma}(\varepsilon_c)}{\tilde{\alpha}_{\gamma\gamma}(\beta)} \frac{6eh}{\sigma_{\text{Th}} m_e c^{1/2} \varepsilon_c^{\gamma \gamma}} \varepsilon_c^{\gamma \gamma} \right]^2.$$  \hspace{1cm} (15)

For nearby sources, the luminosity distance is written $d_L(z) \approx cz/H_0$, and the previous expression can be given in terms of flux $F$, instead of luminosity $L$, using the well-known relation $F = L/4\pi d_L^2$:

$$\left( \frac{\nu_c d F_c}{d \nu_c} \right)_{\nu_c = 4\varepsilon_t} = 3.4 \times 10^{-36} \delta^4 \left[ \frac{\tau_{\gamma\gamma}(\varepsilon_c)}{\tilde{\alpha}_{\gamma\gamma}(\beta)} \frac{6eh}{\sigma_{\text{Th}} m_e c^{1/2} \varepsilon_c^{\gamma \gamma}} \varepsilon_c^{\gamma \gamma} \right]^2 r_{\text{max}}^{-1}.$$  \hspace{1cm} (16)

For an observed SED and a given value of the opacity parameter $\tau_{\gamma\gamma}$, the only remaining unknown quantity in the previous equation is the beaming Doppler factor $\delta$. Each value of $\delta$ gives a point in the $\log \nu$-$\log(\nu F_\gamma)$ plane lying on straight line of slope 2, the level of the curve depending only on the value of $\tau_{\gamma\gamma}$. The intersection of the synchrotron spectrum with the straight line directly constrains the minimum value $\delta_{\text{min}}(\tau_{\gamma\gamma})$ of the beaming Doppler factor required to avoid the $\gamma-\gamma$ absorption with an opacity value of $\tau_{\gamma\gamma}$ of the IC bump (at the peak frequency).

#### 2.4. Application to Mrk 501

We apply this calculation to the case of the Mrk 501 object from 1997 April 16 when the BeppoSAX satellite (Pian et al. 1998) and the CAT imaging atmospheric Cerenkov telescope (Djannati-Ataï et al. 1999; Barrau et al. 1998) recorded simultaneous data (see Fig. 2). All observational parameters we need in the equation (16) are reported Table 2. We consider the two cases in which we take into account, or not, the attenuation of the high-energy component by the DrBl. This effect consists in the interaction of emitted gamma-rays during their travel through the universe with the photon field of the DrBl to create pairs (Gould & Schréder 1967a, 1967b; Stecker et al. 1992; Vassiliev 2000). The tail of the high-energy spectra is then dereddened using the method described in Sauge & Henri (2004). This situation changes the position of IC peak and the Compton dominance parameter $r_{\text{max}}$. In this case, for $\tau_{\gamma\gamma} = 1$, we obtain both in the reddened and the dereddened case $\delta_{\text{min}}(1) \approx 50$ (see Fig. 3).

Given equation (16), the position of the line constraining $\delta$ depends on the value of $(\varepsilon_c^{\gamma \gamma})^4/L_{\gamma\gamma}$. It turns out that the IR unfolding of the spectrum also changes both quantities, and the previous ratio depends only slightly on the level of assumed absorption. The values of $\delta_{\text{min}}(1)$ are thus quite similar in the two cases because when we correct the IC bump, the position of the maximum moves both in luminosity and in frequency. This effect can be clearly seen on Figure 3, in which the difference between the two panels is hardly perceptible.

Note that in fact the level of the curve depends implicitly also on the value of modified power-law index of the spectrum $\beta$. 

![Fig. 2.—SED of Mrk 501 during the flaring period in 1997 April, showing the simultaneous data taken by BeppoSAX (Pian et al. 1998) and by the CAT imaging atmospheric Cerenkov telescope (Djannati-Ataï et al. 1999; Barrau et al. 1998). Regarding high-energy data points, filled gray circles are the ones observed with CAT, while open squares are unabsorbed ones, corrected from our estimation of the DrBl attenuation.](image-url)
In our case, we choose a value $\beta = 0.5$ directly measured on the SED.

### 2.5. Constraint from the Variability Timescale

Another constraint can be derived from the observation of short variability timescale. The classic argument is that a spherical static source cannot be variable on a timescale smaller than $R/c$. So one gets an upper bound on the radius of the source,

$$R_{\var} \leq \frac{\delta c}{t_{\var}}.$$

Combining previous inequality to equation (8) and expressing all the quantities in their fiducial units, we finally get a constraint similar to the one obtained in the previous section for the local synchrotron shape (see eq. [16]):

$$\frac{dF_s}{dv_s} \left|_{v_s = \beta v_{s,\max}} \right. \leq 8.3 \times 10^{-26} \text{erg s}^{-1} \text{cm}^{-2} \left[ \frac{h}{\beta v_{s,\max}^{\max}} \right] 10^{\gamma/\max} \gamma_{\max}^{2} t_{\var}^{-1} r_{\max}^{2}.$$

(17)

Taking a characteristic variability timescale of roughly 15 minutes, we obtain the left solid thick line displayed on Figure 3. It appears that this constraint is less restrictive than the previous one. In context of homogeneous modeling, it gives a minimum value for Doppler factor of 6–8 and 8–10 for the reddened and the dereddened cases, respectively.

### 3. THE CASE FOR A LOW LORENTZ FACTOR

In this section we briefly review all the arguments and evidence in favor of moderate or low values of the bulk Lorentz factor.

#### 3.1. Absence of Superluminal Motion at the Parsec Scale

Observations at the VLBI scale ($\approx$ mas) show that blazars often display superluminal apparent velocities. This phenomenon, predicted by Rees (1966), is expected for relativistic moving sources that are highly beamed and closely aligned with the observer’s line of sight. For a component moving along the jet axis at a reduced speed $\beta = \beta/c$ and making an angle $\theta$ from the line of sight, the apparent transverse velocity measured by the observer is

$$\beta_{\text{app}} = \beta \sin \theta \frac{1}{1 - \beta \cos \theta} \leq \beta \Gamma.$$

(18)

![Fig. 3.—Constraints on the local shape of the synchrotron spectrum of Mrk 501 during the 1997 April 16 high state. The gray polygon is obtained considering the gamma-ray transparency argument. It is defined by the zone in which opacity extends into the interval $r_{\var} \times A_{\beta} A_{\gamma} \in [0.1, 1]$, where $A_\beta$ is the characteristic opening angle of the jet. A constraint coming from the typical variability timescale leads to the leftmost straight thick line. Also represented (dot-dashed line) is a spectrum with a spectral index equal to 4/3 in $\nu F_\nu$ resulting from the emission of a (quasi-) monoenergetic distribution of electrons and/or positrons. Left (respectively, right): The dereddened (reddened) case.](image-url)
If $\beta > \beta_{\text{crit}} = \sqrt{2}/2$ and $\theta$ is such that $\sin 2\theta > (1^2 - 1)^{-1}$, the motion will appear to be superluminal, i.e., $\beta_{\text{app}} > 1$. Expressed in degrees, the latter condition can be written $\theta > \theta_{\text{crit}} = 0.28 (\Gamma/10) - 2^\circ \text{deg}$.

As a matter of fact, VLBI/VLBA campaigns have not clearly succeeded in finding superluminal motion at the parsec scale for any TeV blazars (Edwards & Piner 2002; Piner & Edwards 2004). Observed radio components seem to be stationary or subluminal, requiring low or moderate values of the Lorentz factor ($\Gamma \approx 2-4$). The absence of superluminal motion could be explained by a very close alignment of the jet with the line of sight. Indeed, following the previous expressions, if $\theta \leq 1/\Gamma^2$, the apparent velocity is always smaller than $c$, and the object appears to be subluminal despite the large value of $\Gamma$. But in this case, a simple statistical argument based on the density number of unbeamed counterparts rules out this possibility, as we see in the next section.

Moreover, derived values of the brightness temperature of the VLBI core is on the order of $10^{10}-10^{11}$ K and lie well below the usual IC limit of $\approx 10^{11}-10^{12}$ K necessary to avoid the “inverse Compton catastrophe,” i.e., a situation in which ultrarelativistic particles undergo dramatical Compton cooling in a very short time. Piner & Edwards (2004) have concluded that the jet should be only mildly relativistic at the parsec scale. They propose that the TeV-emitting inner jet is strongly decelerated before reaching the parsec scale. However, we see in the following that the existence of the highly relativistic motion is challenged by other observational facts concerning the statistics of beamed versus unbeamed sources.

3.2. Number of Beamed Sources in the BL Lac/FR I Unification Paradigm

As we said in the introduction, the blazar phenomenon arises from a close alignment of the jet axis with the observer’s line of sight. Following this scheme, one expects the existence of sources sharing the same physical properties (i.e., intrinsically the same objects), but which are viewed at larger angle. It has been proposed that Fanaroff-Riley radio galaxies can be the unbeamed parent population of blazars and, in particular, that FR I galaxies can be the counterparts of BL Lac objects (Urry & Padovani 1995). The unification hypothesis can be tested on samples of objects both by their luminosity ratio (i.e., intrinsically the same objects), but are viewed at different angles, the bolometric luminosity contrast between the two parent populations (BL Lac and FR I radio galaxies) is given by

$$\varpi = \frac{L_{\text{Lac}}}{L_{\text{FR I}}} = \left( \frac{\delta_{\text{Lac}}}{\delta_{\text{FR I}}} \right)^4.$$  (22)

In the case of BL Lac objects, relativistic beaming requires $0 < \theta < 1/\Gamma$, or equivalently $2\Gamma > \delta_{\text{Lac}} > \Gamma$. On the other hand, we suppose that off-axis counterparts verify $\delta = \approx 2/\Gamma$ (corresponding to an average angle value of $\theta \approx 60^\circ$ for $\Gamma > 1$). Then equation (22) can be rewritten

$$\Gamma^4 \gg \varpi \gg \frac{\Gamma^8}{16}.$$  (23)

This estimate can of course be complicated by an intrinsic luminosity distribution. It may be also that we cannot detect the unbeamed sources due to limited sensitivity of the instrument. However, some other indicators, such as the power of the extended radio lobes or the galaxy luminosity itself, are not highly beamed and can serve as an unbiased criterion to select samples.

Capetti & Celotti (1999) studied a sample of 12 BL Lac objects and 5 FR I sources with the Hubble Space Telescope (HST) and compared the core luminosity ratio between objects sharing similar radiative properties. It clearly appears that the entire emission of the BL Lac cores is roughly $10^2-10^5$ times brighter than that from the corresponding radio galaxies. Moreover, Chiaberge et al. (2000) performed similar work on a larger and more complete sample. They obtained roughly the same
conclusions: the luminosity ratio between BL Lac objects and radio galaxies falls in the interval $10^{2.5} - 10^{5.5}$. Applying relation (23), we obtain typical values of $\Gamma \approx 2 - 5$ for the bulk Lorentz factor. Chiaberge et al. also compare the broadband spectra of both classes of objects and find that the spectra could be deduced by a simple Doppler boosting, although once again with modest values of the Doppler factor.

3.4. Detection of a TeV Unbeamed Source: The Case of M87

The nearby giant elliptical radio galaxy M87 (NGC 4486, $z_t \approx 0.00436$) is the first (and for the time being unique) detected unbeamed radio-loud source in the TeV energy range. The first detection was reported by the HEGRA collaboration with an integral flux above 250 GeV at about 3.3% of the flux of the Crab Nebula (with a significance of 4.7 $\sigma$) during a high state (Aharonian et al. 2003; Beilicke et al. 2004b). Such TeV events are confirmed by recent measurements by the High Energy Stereoscopic System (HESS; Beilicke et al. 2004a). The powerful radio jet of M87 has been thoroughly studied in various wavelengths from radio to X-ray, with the results showing that the jet axis is positioned at an angle between 30° and 40° to the line of sight. This angle is clearly large enough to ensure that the emission is unboosted. Previous works based on the study of the proper motion of the VLBI knots (Biretta et al. 1995) or on the detailed analysis of HST and VLA observations (Lobanov et al. 2003) converge toward a value for $\Gamma$ of 3–5 at the kiloparsec scale. The jet differential flux of a source expressed in the observer’s frame can be written as a function of the intrinsic differential luminosity as

$$F'_\nu (\nu; \theta, z) \approx (1 + z)^3 \frac{L'_\nu (\nu)}{4\pi d_L^2},$$ (24)

with $\nu = \nu' \delta s/(1 + z)$ and where $dl(z) \approx zc/H_0$ is the usual luminosity distance. We now consider two different versions of the same intrinsic object, a beamed one corresponding to a blazar and an unboosted one corresponding to a radio galaxy. In this case, the ratio $\mathcal{R}$ of the observed photon fluxes above some threshold frequency $\nu_{th}$ is written

$$\mathcal{R} = \int_{\nu_{th}}^{\nu_{max}} \frac{dF'_{\nu}(\nu)/\nu}{F'_\nu (\nu)/\nu},$$

$$= \left( \frac{z_b}{z_u} \right)^2 \left( \frac{1 + z_b}{1 + z_u} \right)^{\alpha - 2} \left( \frac{\delta_u}{\delta_b} \right)^{2 + \alpha},$$

$$= k(z_b, z_u; \alpha) \left( \frac{\delta_u}{\delta_b} \right)^{2 + \alpha},$$ (25)

where the index $u$ (respectively, $b$) refers to the observed unboosted (beamed) quantities and where we suppose that the high-energy spectrum can be expressed as a simple power law with a photon index $\alpha$.

For beamed sources, the Doppler factor can be written as $\delta_b \approx 2 \Gamma$, while for the unbeamed case one has $\delta_u = 1/(\Gamma (1 - \beta \cos \theta)) < \Gamma$ with $\beta \approx 1 - 1/2 \Gamma^2$. Finally, we can express the bulk Lorentz factor as a function of $\theta$ and the observational parameters as

$$\Gamma(\theta) \approx \frac{1}{\delta} \left[ \frac{k(z_b, z_u; \alpha)}{R} \right]^{1/2} \left[ \frac{1/(2 + \alpha) - \cos \theta}{2(1 - \cos \theta)} \right]^{1/2},$$ (26)

A raw approximation of the previous expression is

$$\Gamma(\theta) \approx \frac{1}{\delta} \left[ \frac{k(z_b, z_u; \alpha)}{R} \right]^{1/2} \left[ \frac{1/(2 + \alpha) - \cos \theta}{2(1 - \cos \theta)} \right]^{1/2},$$ (27)

showing the $1/\delta$ functional dependence of $\Gamma$ and its slow power-law variation with $R$ (or $k$). For instance, for a typical value of $\alpha = 2.5$, a factor of 10 for $R$ implies only a factor $10^{1.5} \approx 1.29$ for $\Gamma$.

In the 1997 April flaring period, the TeV blazar Mrk 501 ($z \approx 0.034$) became roughly 8 times as bright as the Crab Nebula, as reported by the French CAT collaboration (Djannati-Atai et al. 1999). Assuming that Mrk 501 is an unboosted counterpart of M87, with an angle $30' < \theta_{Mrk87} < 40'$, we obtain $4 \leq \Gamma \leq 5.3$. Again we find that the luminosity ratio is compatible with modest values of the Lorentz factor. Due to the increasing sensitivity of the present and the next generation of the imaging atmospheric Cerenkov telescope arrays, the detections of more and more TeV radio galaxies should help us to constrain the dynamics of the emitting plasma at the subparsec scale in a more reliable statistical way.

3.5. Summary

All the above considerations show that observational data are compatible with the beaming model only if the bulk Lorentz factor for the X-ray– and TeV–emitting part of the object is relatively low, between 3 and 5. This value reproduces correctly the luminosity ratio and the statistical number of sources (which are a priori independent factors). Conversely, a value of $\Gamma = 12.5$, which is the minimum typical value derived from the one-zone modeling approach, would lead to a luminosity contrast of $\sim 10^6$. This latter estimation is clearly not compatible with the previous observations ascertaining the “bulk Lorentz factor crisis of TeV blazars. In the following, we examine some suggestions made by various authors to solve the crisis.

4. HOW TO SOLVE THE CRISIS

4.1. Two Pattern Model

Chiaberge et al. (2000) and Trussoni et al. (2003) argue that a jet velocity structure can solve the problem of the unification scheme for BL Lac and FR I objects. They consider a (mildly) relativistic external layer and a fast internal spine that dominates the emission in the case of a favorable alignment along the observer’s line of sight, i.e., in the blazar case. Although similar in appearance to the two-flow model of Pelletier & Sol (1992; see below for details), it differs by the fact that both flows are relativistic, one with a “low” Lorentz factor (around 3) and one with a high Lorentz factor (at least 10). In the following, we consider the same approach in regard to two-component modeling of the velocity structure, in which a fast inner structure is supposed to be surrounded by a slow one (Fig. 4). Each of these components, respectively, is characterized by a bulk Lorentz factor $\Gamma_r$ and $\Gamma_r$. As we saw, the radiative emission of the moving source with a bulk Lorentz factor $\Gamma$ is beamed in a cone sustained by a solid angle $\Omega_c = \pi \Gamma^2$ along the motion. Therefore, the emitted radiation appears to be Doppler boosted when the jet axis points into $\Omega_c$ around the observer’s line of sight. In this case, the luminosity contrast between the two parent populations (BL Lac objects and FR I radio galaxies) is written

$$\Gamma_s \leq \omega \leq (\Gamma_r \Gamma_r)^4,$$ (28)
where the right and left bounds correspond to the case in which, respectively, the fast velocity component either dominates or does not dominate the emission. Unification models are sensitive to the slow component only, so $(\Gamma_s > \Gamma_f \Leftrightarrow$ BL Lac object and, more precisely, if $\theta \ll 1/\Gamma_s$ the fast inner component dominates the emission. In this latter case we are dealing with a TeV blazar.

where $n$ TeV emitters among the same sample of $N$ BL Lac objects, which is written

$$P(n \gg n_0 \gg n) = \sum_{k=n}^{\infty} P(k/N).$$

4.1.2. Applications

For a given value of $(n, N)$, requiring that $P(n \gg n_0 \gg n)$ is larger than an a priori probability $P_0$ constrains the space parameters $(\Gamma_s, \Gamma_f)$. The latter inequality leads to the elimination of the parameter region lying above a straight line, which corresponds equivalently to a constant value of $P_0$ or of the ratio $\Gamma_s/\Gamma_f$ (see eqs. [32], [31], and [30]). Further restrictions come from the gamma-ray transparency argument and unification models of FR I and BL Lac objects, as developed above.

1. First, the gamma-ray transparency argument developed in the first part of this work directly constrains the value of the fast component, as it requires a minimum value of the Doppler factor $\delta_{\min}$ and therefore $\Gamma_f > \delta_{\min}/2$. We have shown that $\delta_{\min} \approx 50$, excluding all of the $\Gamma_s$, $\Gamma_f$ lying above $\Gamma_f = 25$.

2. Second, a basic statistical argument regarding on the number FR I radio versus regarding BL Lac galaxies and a comparison of the luminosity distribution of the previous populations constrain the value of the slow component to reasonable values less than $\Gamma_s \approx 7$ (Urry & Padovani 1995; Chiaberge et al. 2000).

We test this result on the catalog of BL Lac objects from Padovani & Giommi (1995). At $z < 0.13$ they report 29 BL Lac objects with known redshift. Setting $n = 7$, and $N = 29$, $P_0 = 1\%$, and recalling that $\Gamma_{s,\max} = 7$ and $\Gamma_{f,\min} = 25$, the intersection of all listed previous constraints reduces to a null region (see Fig. 5). Even with the hypothesis of a structured flow, the large value of the Lorentz factor required by one-zone homogeneous models is clearly untenable (excluded with a confidence level of 99\%).

We demonstrate that even if the two-component velocity structure can give a satisfactory answer to the luminosity problem of the BL Lac objects even with large value of the Doppler factor required by high-energy emission models, it fails to explain the
detection statistics of the TeV emitters within the BL Lac object population supposed to be off-axis FR I sources.

5. DISCUSSION

5.1. Inhomogeneous Models

Altogether, the previous considerations lead to a serious paradox, in which a high Lorentz factor larger than 20 seems mandatory to avoid strong $\gamma-\gamma$ absorption, whereas all other indications tend to favor modest values around 3. The only way to solve the discrepancy seems to be to give up the implicit assumption of all one-zone models, i.e., that all photons are produced cospatially and simultaneously in some characteristic region of size $R$. Alternatives to one-zone models have already been discussed in the literature. For instance, in the “blob-in-jet” model (Katarzyński et al. 2001, 2003), low-energy photons are produced in a continuous jet, and only the high-energy ones are produced in a spherical blob. This makes it possible to fit the overall spectrum with a smaller Doppler factor of around 15. Another possibility is to treat the variability explicitly and use a time-dependent model to reproduce the data (Krawczynski et al. 2002). Again, the constraints arising from $\gamma-\gamma$ opacity can be somewhat relaxed because soft photons are emitted at a later stage than high-energy ones. As has been remarked by Ghisellini et al. (1985), a time-dependent model will produce effects comparable to inhomogeneous ones. If the evolving source is moving at a relativistic velocity, and many flares are contributing to the emission, the overall system will be in fact a stratified jet composed of many “one-zone” regions in different evolutionary stages. However, none of these models uses bulk Lorentz factors as low as 3.

We are thus led to consider models in which photons are distributed along a jet in a continuous structure, instead of filling a spherical source. In this case, the luminosity is proportional to the photon density times the lateral surface of the jet, which is $2\pi R_j h_j = 2\pi AR_j^2$, where $R_j$ and $h_j$ are the typical jet radius and length at the emission region and $A = h_j/R_j$ is an aspect ratio of the source. For a self-similar jet for which all quantities (radius, magnetic field, particle density, etc.) are described by a power law as a function of the distance $z$, one expects $h_j \sim z$, where $z$ is the distance of the emitting region from the center. It follows that $A \sim z_j/R_j \sim \theta_j^{-1}$, where $\theta_j$ is the typical opening angle of the jet. One can see that for a given synchrotron luminosity and photon density (implying the same IC luminosity), one must conserve the quantity $AR_j^2$, so the typical radius of the jet, and hence the $\gamma-\gamma$ optical depth, will be reduced by a factor of $A^{1/2}$ with respect to a spherical source. This simple geometrical modification helps thus to increase luminosity without increasing optical depth. Furthermore, the particle distribution need not be the same all along the jet. Rather, one expects a gradual cooling of the particles, the overall spectrum being the envelope of all slices of the jet. The local photon spectrum can thus be different from the observed one, and in particular the local soft photon density can be much lower, helping again to reduce the $\gamma-\gamma$ optical depth. As we shall see, all these factors can offer a clue to the bulk Lorentz factor crisis, but they imply strong constraints on the physical picture of relativistic jets.

5.2. Theoretical Implications

5.2.1. Local Photon Density

In view of the above constraints, we take the opposite attitude, considering that the value of the bulk Lorentz factor is constrained by the unification models and the detection of unbeamed sources is around 3. The typical high-energy emission zone as defined above is equivalent to the superposition of $A$
spherical sources with individual luminosities $L_v/\gamma \sim \theta L_v$. Therefore, all previous equations in §2 are still valid, provided we replace the observed luminosity $L_v$ by $\theta L_v$. Using equation (12), we conclude that all opacity constraints remain unchanged if we replace the optical depth $\tau_{\gamma\gamma}$ by $\theta^{1/2}\tau_{\gamma\gamma}$. This is of course in accordance with the estimate made in the previous paragraph. So we can use Figure 3 with slightly different values of $\tau_{\gamma\gamma}$. The typical angle $\theta$ must be of the order of $10^{-2}$ to $10^{-4}$, so the optical depth will be reduced by a factor between 3 and 10. In the following, we still use the same line, $\theta^{1/2}\tau_{\gamma\gamma} \approx 1$, to constrain the optical depth, meaning that $\tau_{\gamma\gamma} \leq 0.1$ to 0.3.

5.2.2. Local Photon Spectrum

We can thus put an upper limit on the soft photon luminosity corresponding to this value and a Doppler factor of 3, which constrains the soft photon luminosity at an energy of $\epsilon_\gamma = \delta^2 \epsilon_c \sim 10$ eV. As the spectrum is by definition approximately the same throughout the characteristic emission region, we can thus estimate the local photon spectrum by interpolating between the peak synchrotron luminosity and the upper limit given above. Inspection of Figure 3 shows that the spectral index between 10 and 100 keV is very close to $1/2$, which is characteristic of a quasi-monoenergetic distribution, an example of which is provided by the quasi-Maxwellian or “pileup” distribution (Henri & Pelletier 1991; Schlücker 1985; Saugé & Henri 2004), which is a natural outcome of some acceleration processes such as second-order Fermi acceleration or magnetic reconnection. This distribution is not the usual power law that is often claimed to exist in AGNs and is naturally produced in MHD shocks. Rather than localized shocks, the assumption of a low Lorentz factor leads to a picture of a continuous jet filled by relativistic particles, continuously reheat by a diffuse acceleration mechanism.

5.2.3. Pair Production

The local synchrotron spectrum cannot be harder than the monoenergetic one, so Figure 3 proves that this implies a lower limit to the quantity $\theta^{1/2}\tau_{\gamma\gamma} \geq 1$. Thus, the limit on $\gamma$-photon optical depth cannot be very low, unless we have an extremely well collimated jet that is not supported by the general FR I morphology. Modest collimation factors imply that $\tau_{\gamma\gamma} \geq 0.1$. This supports the formation of an electron-positron pair plasma in the acceleration site. If the acceleration is not localized, which is suggested by the picture of a continuous jet filled by a pileup distribution, the pairs created by $\gamma$-$\gamma$ interaction cannot avoid being reaccelerated and will trigger a pair cascade (Henri & Pelletier 1991). These constraints are thus suggestive of a inner continuous pair-dominated, jetlike emission zone, maintained at a relativistic temperature.

5.3. Compatibility with the Two-Flow Model

All the previous considerations find a natural explanation in the context of the two-flow model, which was proposed to account for the formation of relativistic jets in AGNs. In this model, extragalactic jets are in fact the results of a double structure: a first jet, not highly but only mildly relativistic ($v \approx 0.5c$), is emitted by an MHD mechanism by a large-scale magnetic field anchored in an accretion disk (Blandford & Payne 1982; Ferreira & Pelletier 1993a, 1993b, 1995); this powerful, but weakly dissipative jet can sustain MHD turbulence able to accelerate nonthermal particles. These particles will produce synchrotron and gamma-ray photons, and if the optical depth becomes large enough, these photons will trigger an intense pair cascade, leading to a dense pair plasma in the empty “throat” of the jet. We have shown in previous works that this pair plasma will be spontaneously accelerated to relativistic velocities even if the surrounding jet is not highly relativistic by itself, by the so-called Compton rocket effect, which is a recoil effect associated with the anisotropic IC process originally introduced by Odell (1981). The Compton rocket effect has been shown to be inefficient in accelerating an isolated relativistic plasma because the cooling time is always shorter than the bulk acceleration time (Phinney 1982). In the two-flow model, however, the heating by the surrounding jets compensates for the cooling and the pair plasma remains relativistic over large distances (Marcowith et al. 1995).

Detailed calculations of the Compton rocket effect in this configuration (Renaud & Henri 1989) show that the pair plasma accelerates gradually in the vicinity of an accretion disk, being maintained to a quasi-equilibrium Lorentz factor $\Gamma_{\text{eq}} \approx (z/r_i)^{1/4}$, where $r_i$ is the inner radius of the accretion disk (3 gravitational radii for a Schwarzschild black hole). The equilibrium Lorentz factor is defined by the fact that the photon field of the accretion disk, seen in the comoving frame, appears to be nearly isotropic due to relativistic aberration. It grows slowly with the distance, because the field becomes more and more anisotropic. The acceleration continues until the photon density becomes too low to efficiently accelerate the plasma. Then the plasma decouples from the ambient radiation field and ends up with an asymptotic ballistic motion at constant $\Gamma_b \rightarrow \Gamma_{\text{b,eq}}$, which depends on the disk luminosity and the particle energy distribution. For a relativistic energy distribution function $n(\gamma) \propto \gamma^2 \exp(-\gamma^2/2)$, the asymptotic bulk Lorentz factor is approximately $\Gamma_{\text{b,eq}} \approx (\ell_\gamma)^{1/2}$, where $\ell_\gamma = L_\gamma \sigma_T / 4\pi m_e c^3 r_i$ is the soft photon compactness and $\ell_\gamma$ is the characteristic energy of the pileup depending on the details of the acceleration/cooling processes (Renaud & Henri 1998).

The first interesting feature in this model is that it predicts naturally a gradual acceleration from the core. The value of $\Gamma_b \approx 3$ is naturally obtained at $\approx 100 r_g$, which is a typical distance at which gamma-ray emission seems to occur, based on variability arguments. Thus, low Lorentz factors are not surprising in this model but are explained naturally. As a matter of fact, very high values of 20 near the core would be difficult to explain in this frame!

The second feature is that the asymptotic bulk Lorentz factor is controlled by the density of the photon field emitted by the accretion disk. For BL Lac objects and FR I galaxies, the disk luminosity is known to be much lower than luminous FSRQ and FR II galaxies, by a factor around $10^{-3}$. One would thus expect a lower asymptotic Lorentz factor for BL Lac objects on average, which would help in understanding the absence of superluminal motion in TeV blazars. As a matter of fact, numerical estimates show that the expected asymptotic Lorentz factors are between 10 and 20 for near-Eddington accreting supermassive black holes, whereas, in contrast, they are between 5 and 10 for low-luminosity AGNs. We note that bulk Lorentz factors around 5 are indeed observed in M87, which would mean that the decoupling occurs at some thousands of Schwarzschild radii from the core. Unification models are compatible with slowly accelerating jets, the inner (X-ray emitting) jets having bulk Lorentz factors around 3 and the outer radio jet having a larger Lorentz factor around 7 (Urry & Padovani 1995). Again this is perfectly compatible with the predictions of the two-flow model, with an inner jet emitting X-ray and TeV radiation with a modest bulk Lorentz factor and an outer jet responsible for radio emission with a higher one. Also we note that there is no need for deceleration to explain FR I mildly relativistic jets: even if the large-scale jet has only a moderately relativistic velocity $v \approx 0.5c$, this can be attributed to the “slow” MHD component surrounding the relativistic beam, the latter being dissipated on a kiloparsec scale.
An inhomogeneous model offers also a convenient explanation for the lack of obvious correlation between X-ray and gamma-ray variability. If the magnetic field is varying along the jet, the photons with a given energy could be produced by electrons with different energies and locations. If several flares contribute to the observed spectrum—which is necessary to account for the global spectral shape in case of a monoenergetic distribution—a complicated variability pattern could emerge. This is much less easy to understand in the homogeneous steady state models. Thus we think that inhomogeneous models, although more complicated to compute, seem to be unavoidable to explain the spectral and temporal features of TeV blazar emission.

6. CONCLUSION

We have investigated in detail the so-called bulk Lorentz factor crisis of TeV blazars, which seem to imply an incompatibility between a high Lorentz factor required to ensure gamma-ray transparency and a low Lorentz factor deduced from statistical arguments and luminosity contrast, including the detection of the nonblazar TeV source M87. We show that the transparency argument is common to all one-zone models, and that the only way of solving the paradox is to consider inhomogeneous jet models, where all photons are not produced cospatially. The spectrum is then the spatial convolution of different jet slices, and the opacity problem can be avoided by invoking geometrical arguments and harder local photon spectrum. We show, however, that for modest values of geometrical beaming of the jet, which seem natural considering the morphology of FR I galaxies, the optical depth for $\gamma-\gamma$ absorption cannot be very low, even for a local quasi-monoenergetic particle distribution. This has profound implications for the physics of the jet: the acceleration mechanism must be distributed all along the jet and is more probably ensured by a second-order Fermi mechanism or reconnection sites than by localized shocks. A moderately high value of $\gamma-\gamma$ optical depth implies a fair production rate of electron-positron pairs, which are likely to be reaccelerated by the acceleration process to trigger a pair cascade. All these features are natural consequences of the two-flow model, which attributes the relativistic phenomena (high energy emission and superluminal motion) to the formation of such a pair plasma inside a powerful but mildly relativistic jet ensuring the confinement and the heating of the relativistic beam. The bulk Lorentz factor is also fully in accordance with a continuous acceleration along the jet by the Compton rocket effect, which predicts naturally $\Gamma_b \approx 3$ at 100 Schwarzschild radii from the core. We conclude that all observational facts are more in accordance with light, moderately relativistic leptonic beams than with highly relativistic baryonic jets.

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REFERENCES

Aharonian, F., et al. 2005a, A&A, 430, 865
——. 2005b, A&A, 436, L17
——. 2004, A&A, 421, 529
——. 2003, A&A, 403, L1

Barrau, A., et al. 1998, Nucl. Instrum. Methods Phys. Res. A, 416, 278

Beilicke, M., Cornils, R., Heinzelmann, G., Raue, M., Ripken, J., Benbow, W., Horns, D. & Tluczykont, M. 2004a, in Proc. 22nd Texas Symposium on Relativistic Astrophysics, http://www.slac.stanford.edu/econf/C041213/papers/2403.PDF

Beilicke, M., Götting, N., & Tluczykont, M. 2004b, NewA Rev., 48, 407

Biretta, J. A., Zhou, F., & Owen, F. N. 1995, ApJ, 447, 582

Blandford, R. D., & Payne, D. G. 1982, MNRAS, 199, 883

Capetti, A., & Celotti, A. 1999, MNRAS, 304, 434

Chiaberge, M., Celotti, A., Capetti, A., & Ghisellini, G. 2000, A&A, 358, 104

Coppi, P. S., & Blandford, R. D. 1990, MNRAS, 245, 453

Djannati-Ataï, A., et al. 1999, A&A, 350, 17

Edwards, P. G., & Piner, B. G. 2002, ApJ, 579, L67

Ferreira, J., & Pelletier, G. 1993a, A&A, 276, 625
——. 1993b, A&A, 276, 637
——. 1995, A&A, 295, 807

Gaidos, J. A., et al. 1996, Nature, 383, 319

Ghisellini, G., Maraschi, L., & Treves, A. 1985, A&A, 146, 204

Gould, R. J., & Schreder, G. P. 1967a, Phys. Rev., 155, 1404
——. 1967b, Phys. Rev., 155, 1408

Hardcastle, M. J., Worrall, D. M., Birkinshaw, M., & Canosa, C. M. 2003, MNRAS, 338, 176

Henri P., & Pelletier G. 1991, ApJ, 383, L7

Katarzyński, K., Sol, H., & Kus, A. 2001, A&A, 367, 809
——. 2003, A&A, 410, 101

Krawczynski, H., Coppi, P. S., & Aharonian, F. 2001, ApJ, 559, 187
Krawczynski, H., et al. 2002, MNRAS, 336, 721
Lebanov, A., Hardee, P., & Eilek, J. 2003, NewA Rev., 47, 629
Marcowith, A., Henri, G., & Pelletier, G. 1995, MNRAS, 277, 681
Marscher, A. P. 1983, ApJ, 264, 296

Odell, S. L. 1981, ApJ, 243, L147

Padovani, P., & Gioni, M. 1995, MNRAS, 277, 1477

Pelletier, G., & Sol H. 1992, MNRAS, 254, 635

Phinney, E. S. 1982, MNRAS, 198, 1109

Pian, E., et al. 1998, ApJ, 492, L17

Piner, B. G., & Edwards, P. G. 2004, ApJ, 600, 115

Punch, M., et al. 1992, Nature, 358, 477

Quinn, J., et al. 1996, ApJ, 456, L83

Rees, M. J. 1996, Nature, 211, 468

Rees, M. J., & Simon, M. 1968, ApJ, 152, L145

Renaud, N., & Henri, G. 2004, ApJ, 616, 136

Schlickeiser, R. 1985, A&A, 143, 431

Stecker, F. W., de Jager, O. C., & Salamon, M. H. 1992, ApJ, 390, L49

Svensson, R. 1987, MNRAS, 227, 403

Tavecchio, F., Maraschi, L., & Ghisellini, G. 1998, ApJ, 509, 608

Trussoni, E., et al. 2003, A&A, 403, 889

Urry, C. M., & Padovani, P. 1991, ApJ, 371, 60
——. 1995, PASP, 107, 803

Vassiliev, V. V. 2000, Astropart. Phys., 12, 217

Zensus, J. A. 1997, ARA&A, 35, 607