One-zone SSC model for the core emission of Centaurus A revisited

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ABSTRACT

Aims. We investigate the role of the second synchrotron self-Compton (SSC) photon generation to the multiwavelength emission from the compact regions of sources that are characterized as misaligned blazars. For this, we focus on the nearest high-energy emitting radio galaxy Centaurus A and we revisit the one-zone SSC model for its core emission. Methods. We have calculated analytically the peak luminosities of the first and second SSC components by, first, deriving the steady-state electron distribution in the presence of synchrotron and SSC cooling and, then, by using appropriate expressions for the positions of the spectral peaks. We have also tested our analytical results against those derived from a numerical code where the full emissivities and cross-sections were used. Results. We show that the one-zone SSC model cannot account for the core emission of Centaurus A above a few GeV, where the peak of the second SSC component appears. We, thus, propose an alternative explanation for the origin of the high energy (∼0.4 GeV) and TeV emission, where these are attributed to the radiation emitted by a relativistic proton component through photohadronic interactions with the photons produced by the primary leptonic component. We show that the required proton luminosities are not extremely high, e.g. ∼10^{43} erg/s, provided that the injection spectra are modelled by a power-law with a high value of the lower energy cutoff. Finally, we find that the contribution of the core emitting region of Cent A to the observed neutrino and ultra-high energy cosmic-ray fluxes is negligible.

Key words. radiation mechanisms: non-thermal – gamma-rays: general – Radio galaxies: Centaurus A

1. Introduction

Centaurus A (Cen A) is the nearest radio galaxy to earth with a luminosity distance D_L ∼ 3.7 Mpc and, therefore, one of the best laboratories for studying the physics of radio lobes, relativistic outflows, shock formation, thermal and non-thermal emission mechanisms. Due to its proximity, emission from the extended lobes and jet as well as from its nucleus has been detected across the electromagnetic spectrum (see e.g. Israel (1998) for a review). In radio wavelengths it has an FRI morphology (Fanaroff & Riley 1974), while in higher energies (X-rays) is regarded as a misaligned BL Lac object (Morganti et al. 1992; Chiaberge et al. 2001) in agreement with the unification scheme of Active Galactic Nuclei (AGN) (Padovani & Urry 1990, 1991). Although the angle between the jet axis and our line of sight is large, it is still not well constrained mainly due to the assumptions used in its derivation (see e.g. Hardcastle et al. 2003); it ranges between 15° (Hardcastle et al. 2003) up to 50° − 80° (Tingay et al. 1998).

Gamma-ray emission (∼0.1 − 10 GeV) from Cen A has been detected by EGRET (Hartman et al. 1999) but the identification of the γ-ray source with the core was rather uncertain due to large positional uncertainties. The recent detection of very high energy (VHE) emission (∼TeV) from the core of Cen A by H.E.S.S. (Aharonian et al. 2004) along with the Fermi satellite observations above 100 MeV from the core (Abdo et al. 2010a) and X-ray data from various telescopes make now possible the construction of a well sampled Spectral Energy Distribution (SED) for its nuclear emission which requires physical explanation. Whether the HE/VHE core emission originates from very compact or extended regions is still unclear because of lacking information regarding the variability in the GeV/TeV energy ranges and of the current resolution of γ-ray instruments. This complicates further the attempts of fitting the multiwavelength (MW) core emission.

The one-zone synchrotron self-Compton (SSC) model is one of the most popular emission models due to its simplicity and to the small number of free parameters. In the past it has been successfully applied to the SEDs of various blazars – see e.g. Ghisellini et al. (1998); Celotti & Ghisellini (2008) for steady-state models and Mastichiadis & Kirk (1997); Bottcher & Chiang (2002) for time-dependent ones. Note, however, that rapid flaring events and recent contemporaneous MW observations of blazars pose problems to homogeneous SSC models (Begelman et al. 2008; Bottcher et al. 2009; Costamante et al. 2009). If FRIs are indeed misaligned BL Lac objects, then one expects that the one-zone SSC model applies also successfully to their MW emission (see e.g. Abdo et al. 2009f) for M87 and Abdo et al. 2009d) for NGC 1275). We note, however, that alternative emission

1 Although there is still considerable debate on its distance, see e.g. Ferrarese et al. 2007; Maiolati 2010, Harris et al. 2010, we will adopt this value as a representative one.

2 We note that high-energy (HE) emission was also detected from the radio lobes of Cent A (Abdo et al. 2010b), while a recent analysis by Yang et al. (2012) shows that this emission extends beyond the radio lobes. However, we will not deal with the lobe emission in the present work.
models have also been proposed (e.g. Giannios et al. (2010) for M87). In the case of Centaurus A it is still the leading interpreting scenario for the core emission, at least below the TeV energy range (Chiang et al. 2008; Abdo et al. 2010a; Ronayette & Böttcher 2011).

However, there is a subtle point that must be taken into account when applying one-zone models to FRIs: due to the large viewing angle the Doppler factor \( \delta \) cannot take large values (in most cases \( \delta < 5 \)) in contrast to blazars where typical values are \( \delta \sim 20 - 30 \), while even higher values \( (40 < \delta < 80) \) appear in the literature (Konopelko et al. 2003; Aleksic et al. 2012). Thus, unless the observed \( \gamma \)-ray luminosity of FRIs is by a few orders of magnitude lower than the one of blazars, the injection power of relativistic radiating electrons must be high enough to account for it. The above imply that in cases where the radius of the emitting source is not very large, higher order SSC photon generations may, in general, contribute to the total SED and are not negligible as in the case of blazars.

In the present work we focus on Cen A as a typical example of a misaligned blazar. We show that the simple homogeneous SSC model cannot fully account for its MW core emission due to the emergence of the second SSC photon generation. We, therefore, present an alternative scenario where the SED up to the GeV energy range is attributed to SSC emission of primary electrons, while the GeV–TeV emission itself is attributed to photohadronic processes.

The present work is structured as follows: in Section 2 we calculate analytically using certain approximations the peak luminosities of the synchrotron and SSC (first and second) components for parameters that are relevant to Cen A; in the same section we test our results against those obtained from a numerical code that employs the full expressions for the cross sections and emissivities of all processes. In Section 3 we show the effects that the second SSC component has on the overall one-zone SSC fit of the MW core emission of Cen A. In Section 4 we introduce a relativistic proton distribution in addition to the primary electron one, and present the resulting leptohadronic fits to the emitted MW spectrum; we also discuss the resulting neutrino and ultra-high energy cosmic ray emission. Finally, we conclude in Section 5 with a discussion of our results.

2. Analytical arguments

The calculation of the steady-state electron distribution in the case of a constant in time power-law injection under the influence of synchrotron and SSC (in the Thomson regime) cooling can be found in Lefa & Mastichiadis (2013) – hereafter LM13. However, in the present section and for reasons of completeness, we derive the analogous solution for monoenergetic electron injection. On the one hand, this choice significantly simplifies our analytical calculations. On the other hand, it is justified since the power-law photon spectrum in the range \( 10^{13} - 10^{15} \) Hz is very steep and it can be therefore approximated by the synchrotron cutoff emission of a monoenergetic electron distribution.

### 2.1. Steady state solution for the electron distribution

We assume that electrons are being injected at \( \gamma_0 \) and subsequently cool down due to synchrotron and SSC losses. Here we assume that all scatterings between electrons and synchrotron photons occur in the Thomson regime, which is true for parameter values related to the spectral fitting of Cen A (see Sections 2.2 and 3). The electron distribution cools down to a characteristic Lorentz factor \( \gamma_c \), where the escape timescale \( (t_{e,\text{esc}}) \) equals the energy loss timescale and it is given by

\[
\gamma_c = \frac{3m_ec}{4\pi\sigma_T}\frac{u_B + u_s}{u_B + u_s},
\]

where \( t_{e,\text{esc}} = R/c, R \) is the size of the emitting region, \( u_p \) is the magnetic energy density and

\[
u_s = mc^2\int_{\gamma_{\text{min}}}^{\gamma_{\text{max}}} d\gamma \epsilon_n(\gamma),
\]

is the energy density of synchrotron photons. The integration limits in eq. (2) are \( \epsilon_{\text{max}} = b\gamma_{\text{c}}^2 \) and \( \epsilon_{\text{min}} = b\gamma_{\text{c}}^2 \), where \( b = B/B_{\text{c}} \) and \( B_{\text{c}} = 4.4 \times 10^{13} \) G. In what follows, all energies that appear in the relations will be normalized with respect to \( mc^2 \), unless stated otherwise. Here we assume that \( \gamma_c \) is much smaller than \( \gamma_0 \), which further implies that the particle escape is less significant than the energy losses in shaping the electron distribution at the particular energy range. Thus, the electron distribution \( \nu_s \) at the steady state is described by the kinetic equation below

\[
Q(\gamma) = \frac{4}{3m_ec^2\sigma_Tc} \frac{\partial}{\partial \gamma} \left[ \gamma^2 \epsilon_n(\gamma) (u_B + u_s) \right],
\]

where and \( Q(\gamma) = Q_0\delta(\gamma - \gamma_0) \) is the injection rate per unit volume of electrons having Lorentz factors in the range \((\gamma - \gamma_0, \gamma + \gamma_0)\). Under the \( \delta \)-function approximation for the synchrotron emissivity, the differential number density of synchrotron photons is given by – see e.g. Mastichiadis & Kirk (1995),

\[
n_\nu(\epsilon) = A_2 e^{-1/2 \epsilon/(\epsilon_0/\epsilon)},
\]

where

\[
A_2 = \frac{2}{3} R \sigma_T u_0 b^{-3/2},
\]

and \( u_0 \) is the dimensionless magnetic energy density, i.e. \( u_0 = u_B/m_ec^2 \). Plugging eqs. (2) and (4) into eq. (3) we find

\[
Q_0\delta(\gamma - \gamma_0) = \frac{4}{3} \sigma_T c \frac{\partial}{\partial \gamma} \left[ \gamma^2 \epsilon_n(\gamma) G_c \right],
\]

where \( G_c \) depends on the electron distribution as

\[
G_c = \left( u_0 + \frac{4}{3} \sigma_T R u_0 \int_{x_c}^{\gamma_0} dx \frac{x^2}{2} \epsilon_n(x) \right).
\]

An ansatz for the solution \( \epsilon_n \) of the above integro-differential equation is \( \epsilon_n(\gamma) = k_c \gamma^{-p} \) with \( k_c \) and \( p \) being the parameters to be determined. By substituting the above solution into eq. (6) we find that \( p = 2 \) and that \( k_c \) satisfies the following quadratic equation

\[
\left( \frac{4}{3} \sigma_T R u_0 \gamma_0 \right) k_c^2 + u_0 k_c - \frac{3Q_0}{4\pi\sigma_Tc} = 0,
\]
with \( k_e \) being the positive root
\[
  k_e = \frac{3}{8 \sigma_T R^2 \gamma_0} \left( 1 + \frac{1}{1 + \frac{4 Q \gamma_0}{c u_b}} \right)^{1/2} \tag{9}
\]

It is more convenient to express \( k_e \) in terms of the electron injection compactness \( \ell_{\text{inj}} \), which is defined as
\[
  \ell_{\text{inj}} = \frac{\sigma_T L_{\text{inj}}}{4 \pi R u_b e^3}, \tag{10}
\]
where \( L_{\text{inj}} \) is the total injection luminosity of electrons. Using the relation between \( Q_0 \) and \( \ell_{\text{inj}} \) for monoenergetic injection, i.e.
\[
  Q_0 = \frac{3 \ell_{\text{inj}}}{8 \sigma_T R^2 \gamma_0} \left( 1 + \frac{1}{1 + \frac{12 \ell_{\text{inj}}}{\ell_B}} \right), \tag{11}
\]
we find
\[
  k_e = \frac{3}{8 \sigma_T R^2 \gamma_0} \left( 1 + \frac{1}{1 + \frac{12 \ell_{\text{inj}}}{\ell_B}} \right)^{1/2}, \tag{12}
\]
where the ‘magnetic compactness’ \( \ell_B = \sigma_T R u_b \) was introduced. There are two limiting cases that can be studied depending on the ratio \( \ell_{\text{inj}} / \ell_B \).

- Synchrotron dominated cooling or \( \ell_{\text{inj}} << \ell_B / 12 \) where we find
  \[
  k_e \approx \frac{9 \ell_{\text{inj}}}{4 \sigma_T^2 R^2 u_b \gamma_0} + O \left( \frac{\ell_{\text{inj}}}{\ell_B} \right)^2 \tag{13}
  \]
- Compton dominated cooling or \( \ell_{\text{inj}} \gg \ell_B / 12 \) where we find
  \[
  k_e \approx \frac{3}{4} \left( \frac{3 \ell_{\text{inj}}}{R^2 \sigma_T^2 u_b \gamma_0} \right)^{1/2}. \tag{14}
  \]

### 2.2. Peak luminosities

The relation between the electron injection rate and the normalization of the distribution at the steady-state (eqs. \( \| \) or \( \| \) ) is crucial for the correct calculation of the peak luminosities. The calculation is complete when the proper expressions of the emissivities and of the energies where the peaks appear are taken into account. Our results, for each emission component, are presented below.

#### Synchrotron component

In the optically thin to synchrotron self-absorption regime, which is the case considered here, the differential synchrotron luminosity per unit volume is given by \( J_{\text{sync}}(\epsilon) = (c/R) u_b(\epsilon) \); we note that the units of \( J_{\text{sync}} \) are erg cm\(^{-3}\) s\(^{-1}\) per dimensionless energy \( \epsilon \). Under the \( \delta \)-function approximation for the synchrotron emissivity, the peak luminosity (per unit volume) of the corresponding component (\( L_{\text{peak}}^{\text{sync}} \)) emerges at \( \epsilon_{\text{peak}}^{\text{sync}} \) and it is given by
\[
  L_{\text{peak}}^{\text{sync}} = \epsilon J_{\text{sync}}(\epsilon) |_{\epsilon = \epsilon_{\text{peak}}^{\text{sync}}} = \frac{2}{3} \sigma_T m e^3 u_b \gamma_0 k_e, \tag{15}
\]
or using eq. \( \| \)
\[
  L_{\text{peak}}^{\text{sync}} = \frac{u_b m e^3}{4 R} \left( -1 + \frac{1}{1 + \frac{12 \ell_{\text{inj}}}{\ell_B}} \right). \tag{16}
\]
We note that if we were to use the full expression for the synchrotron emissivity (e.g. Rybicki & Lightman 1979), the peak in a \( n F_{\epsilon} \) plot would appear at a slightly different energy than \( \gamma_0 \).

#### First SSC component

For parameter values related to the spectral fitting of Cen A, e.g. for \( \gamma_0 = 10^3 \) and \( b \sim 10^{-13} \) we find \( \gamma_0^{\text{sync}} = b \gamma_0^3 << 1 \), i.e. scatterings between the maximum energy electrons with the whole synchrotron photon distribution occur in the Thomson regime. Under the above assumption the peak luminosity (per unit volume) of the first SSC component (\( L_{\text{peak}}^{\text{ssc}}(\epsilon_{\text{peak}}) \)) emerges at
\[
  \epsilon_{\text{peak}}^{\text{ssc}} = \begin{cases} 
    \frac{3}{4} b \gamma_0^3 e^{-\alpha} & \text{for } p < 3 \\ 
    \frac{3}{4} b \gamma_0^2 e^{\alpha} & \text{for } p > 3 
  \end{cases} \tag{17}
\]
where \( \alpha = (p-1)/2 \) is the synchrotron spectral index and \( p \) is the power-law index of the electron distribution at the steady state – see e.g. Gould (1974). In our case the energy of the peak is given by the first branch of the above equation since \( p = 2 \). The peak luminosity is then given by
\[
  L_{\text{peak}}^{\text{ssc}}(\epsilon_{\text{peak}}) = \epsilon_1 J_{\text{ssc}}(\epsilon_1) |_{\epsilon_1 = \epsilon_{\text{peak}}^{\text{ssc}}} = \frac{c}{4 \pi R} m e^2 \epsilon_1 n_{\text{ssc}}^{\text{1p}}(\epsilon_1), \tag{18}
\]
where \( n_{\text{ssc}}^{\text{1p}}(\epsilon_1) \) is the differential number density of SSC photons (1st generation) that is given by
\[
  n_{\text{ssc}}^{\text{1p}}(\epsilon_1) = \frac{4 \pi R 3 \sigma_T e}{4} \int_{\epsilon_{\text{min}}}^{\epsilon_{\text{max}}} \text{d} \epsilon \frac{\epsilon}{\epsilon} I_{\text{c}}(\epsilon_1, \epsilon), \tag{19}
\]
where
\[
  I_{\text{c}}(\epsilon_1, \epsilon) = \int_{\epsilon_{\text{min}}}^{\epsilon_{\text{max}}} \frac{n_{\text{c}}(\gamma)}{\gamma^2} F_{\text{c}}(q, \Gamma_\epsilon). \tag{20}
\]
Here \( F_{\text{c}}(q, \Gamma_\epsilon) \) is the Compton kernel
\[
  F_{\text{c}} = 2q \ln q + (1 + 2q)(1 - q) + \frac{1}{2} (\Gamma_\epsilon q)^2 (1 - q) \tag{21}
\]
and
\[
  \Gamma_\epsilon = 4\epsilon \gamma q + q = \frac{\epsilon_1 - \epsilon}{\epsilon_1 (1 - \epsilon / \epsilon_1)}. \tag{22}
\]
In the Thomson limit, which therefore applies in our case, \( \Gamma_\epsilon << 1 \) and \( \epsilon_1 / \epsilon << 1 \); the Compton kernel takes then the simplified form
\[
  F_{\text{c}} \approx \left( 2 \frac{\epsilon_1}{4 \epsilon \gamma q} \ln \left( \frac{\epsilon_1}{4 \epsilon \gamma q} \right) + \frac{\epsilon_1}{4 \epsilon \gamma q} + 1 - 2 \left( \frac{\epsilon_1}{4 \epsilon \gamma q} \right)^2 \right). \tag{23}
\]
Following Blumenthal & Gould (1970) – henceforth BG70 – we assume that the energies of the scattered photons lie away from the high- and low-energy cutoffs. Since the integrand of \( I_{\text{c}} \) is a steep function of \( \gamma \), the upper cutoff does not contribute to the integration, and \( I_{\text{c}} \) is written as
\[
  I_{\text{c}} = \frac{1}{2} k_e \frac{\epsilon_1}{\epsilon_1} - 3/2 \int_0^1 d y y^{1/2} (2 y \ln y + y + 1 - 2 y^2) = \frac{4 k_e}{3} \frac{\epsilon_1}{\epsilon_1} - 3/2 C_1 \tag{24}
\]
where \( y = \frac{\gamma_0}{\gamma} \) and \( C_1 = 0.975 \approx 1 \). The above expression is then inserted in eq. (19) and we find
\[
u^{(1)}_{\text{ssc}}(\epsilon_1) = 8\pi R^2 \sigma_{1}^{2} K_{2} \gamma_{b} b^{-1/2} C_1 \ln \Sigma_{1} e^{-\epsilon_1/3},
\]
for \( 4\gamma_{0}^{4}/b < \epsilon_1 < 4\gamma_{0}^{4}/b \). In the above, \( \Sigma_{1} \) is the Compton logarithm which also depends on \( \epsilon_1 \). In reality, \( \Sigma_{1} \) changes functional form at \( \epsilon_{c} = \frac{1}{3} \beta_{c}^{2} \gamma_{0}^{2} \) but for the case studied here \((p = 2)\) the departure of \( \nu^{(1)}_{\text{ssc}} \) from a pure power-law with index \(-3/2\), at least away from the cutoffs, is not significant – see also eqs. (27)-(28) in Gould [1979]. Inserting the above expression into eq. (18) we find
\[
L^{\text{peak}}_{\text{ssc},1} = \frac{3\sqrt{3} u_{b} m_{e} e^{3}}{16e R} \left( 1 + \sqrt{1 + \frac{12 \epsilon_{\text{inj}}}{\epsilon_{0}} \frac{1}{r_{B}}} \right)^{2}
\]
where we have neglected the factor \( C_1 \ln \Sigma_{1} \). Whether our choice is justified or not it will be tested later on, by comparing eq. (26) against the results obtained with the numerical code.

**Second SSC component**

As already mentioned in the introduction, in the case of blazars, higher order scatterings, i.e. between electrons and SSC photons of the first generation, are negligible (e.g. see Bloom & Marscher 1996). On the other hand, SSC modelling of SEDs from radio galaxies requires, in general, high electron compactness (\( \epsilon_{\text{inj}}^{(1)} \)) due to the deamplified radiation; of course, this is a rather qualitative argument since the determination of \( \epsilon_{\text{inj}}^{(1)} \) depends also on the absolute value of the observed flux, the ratio of the peak luminosities of the low- and high-energy humps and the size of the emitting region. Here we proceed to calculate analytically the peak luminosity of the second SSC component, which will be then compared to the synchrotron and first SSC peak luminosities.

An analogous calculation to that of eq. (19) for the second generation of SSC photons is, in principle, more complicated because of the Klein-Nishina effects, which for the parameters considered here, become unavoidable. In fact, the scatterings of electrons with SSC photons from the first generation occur only partially in the Thomson and Klein-Nishina regimes. Thus, one must use the full expression of the Compton kernel (e.g. eq. (2.48) in BG70), which hinders any further analytical calculations. In order to proceed, however, we used a simplified version of the single electron Compton emissivity
\[
J_{\text{ssc},2}(\epsilon_{2}) = j_{0}^{(1)}(\epsilon_{2} - \frac{4}{3} \gamma_{2} \gamma_{1}) \frac{2}{3} (\frac{3}{4} - \gamma_{1}),
\]
where the step-function introduces an abrupt cutoff in order to approximate the Klein-Nishina supression and \( j_{0} = 4/3\pi e^{2} \gamma_{b} u_{\text{ssc},1} \). Here \( u_{\text{ssc},1} = m_{e} c^{2} \int \delta\epsilon_{1} n_{\text{ssc}}^{(1)} \) and \( n_{\text{ssc}} \) is approximated by a single power-law, i.e. it is given by eq. (19) without the logarithmic term. The differential luminosity of the second SSC component (per unit volume) is then simply
\[
J_{\text{ssc},2} = 4\pi c \int_{\gamma_{0}}^{\gamma_{c}} \nu_{\text{ssc},1} \int_{\epsilon_{\text{inj}}^{(1)} \text{min}}^{\epsilon_{\text{inj}}^{(1)} \text{max}} d\epsilon_{1} \int_{\epsilon_{1}}^{\epsilon_{\epsilon_{1}} \epsilon_{1}} \frac{d\epsilon_{2}}{d\epsilon_{1}} I_{1}(\epsilon_{2}, \epsilon_{1}, \gamma),
\]
where
\[
I_{1} = \gamma_{2}^{2} \epsilon_{2}(\gamma_{2}) u_{\text{ssc},1}(\epsilon_{1}) \delta(\epsilon_{2} - \frac{4}{3} \gamma_{2} \gamma_{1}) \frac{1}{H \left( \frac{3}{4} - \gamma_{1} \right)}.
\]
After making the integration over \( \gamma \) we find
\[
J_{\text{ssc},2} = \frac{\sigma_{T} c}{\sqrt{3}} u_{\text{ssc},1}^{0} \epsilon_{2}^{-1/2} I_{2}(\epsilon_{2}),
\]
where
\[
I_{2} = \int_{\epsilon_{\min}^{\text{ssc}}}^{\epsilon_{\max}^{\text{ssc}}} d\epsilon_{1} \frac{1}{\epsilon_{1}} H \left( \epsilon_{2} - 4/3 \gamma_{2}^{2} \epsilon_{1} \right) \frac{H \left( \epsilon_{2} - \epsilon_{\min}^{2} \right)}{H \left( \epsilon_{2} - \epsilon_{\min}^{2} \right)}.
\]
Here \( \epsilon_{\min}^{2} = \min[3/4 \epsilon_{1}, 4/3 \gamma_{0}^{2} \epsilon_{1}], \epsilon_{\min}^{\text{ssc}} = 4/3b_{c}^{4}, \epsilon_{\max}^{\text{ssc}} = 4/3b_{c}^{4} \) and
\[
u_{\text{ssc},1}^{0} = 8 \pi m_{e} c^{2} R_{2}^{2} \gamma_{b}^{4} u_{b}^{-1/2} k_{e}^{3}.
\]

The integral of eq. (31) results in the logarithmic term \( \ln 2_{2} \), where \( 2_{2} \) is the ratio of the effective upper and lower limits of the first SSC photon distribution, which do not, in principle, coincide with the actual cutoffs. For the purposes of the present study, however, we will neglect the contribution of the logarithmic term. In most cases, the scatterings that result in the second SSC photon generation are only partially in the Klein-Nishina regime and the quantity \( \epsilon_{2} J_{\text{ssc},2} \) peaks at \( \epsilon_{\text{peak}}^{\text{ssc}} = \gamma_{0} e^{-\gamma_{c}^{2}}, \) where \( \epsilon_{1} \) is the spectral index of the first SSC component and equals \( 1/2 \) in our work – details about the calculation of the SSC peak in different scattering regimes can be found in LM13. Thus, the peak luminosity \( L_{\text{peak}}^{\text{ssc}} \) is given by
\[
L_{\text{peak}}^{\text{ssc}} = \frac{8\pi}{\sqrt{3}} m_{e} c^{2} R_{2}^{2} \gamma_{b}^{4} u_{b}^{-1/2} k_{e}^{3}.
\]

Finally, using eqs. (26) and (34) we define the ratio \( \zeta \) as
\[
\zeta \equiv \frac{L_{\text{peak}}^{\text{ssc}}}{L_{\text{peak}}^{\text{ssc},1}} = \frac{3}{4b_{c}^{2} \gamma_{o}^{2}} \left( 1 + \frac{12 \epsilon_{\text{inj}}}{\epsilon_{0}} \frac{1}{r_{B}} \right)^{3}
\]
In general, if \( \zeta > 1 \) the system can be led to the so-called ‘Compton catastrophe’, where the peak luminosity of the \( n^{th} \) SSC generation is larger than that of the previous one. This succession ceases, however, due to Klein-Nishina effects, as in our case. If the electron cooling is synchrotron dominated (\( \epsilon_{\text{inj}}^{(n)} < 8.3 \times 10^{-2} \)), we find \( \zeta > 1 \) if \( \epsilon_{\text{inj}}^{(n)} > 8.3 \gamma_{o}^{3} R_{15}^{5/2} b_{c}^{4} \), where we used the notation \( Q_{x} \equiv Q/10^{x} \) in cgs units. In this regime, both constraints on \( \epsilon_{\text{inj}}^{(n)} \) cannot be satisfied simultaneously for typical values of \( \gamma_{o} \) and \( R \), thus, making the Compton catastrophe not relevant. On the other hand, in the Compton cooling regime (\( \epsilon_{\text{inj}}^{(n)} > 8.3 \times 10^{-2} \)), \( \zeta \) becomes larger than unity if \( \epsilon_{\text{inj}}^{(n)} > 570 \gamma_{o}^{3} R_{15}^{3/2} b_{c}^{4} \).
Table 1. Peak luminosities (in logarithm) of the synchrotron, first and second SSC components along with the ratio ζ of the two SSC peak luminosities. In each row the numerical (N) and analytical (A) values are shown as the first and second values respectively.

| $\ell_{\text{peak}}^{\text{syn}}$ | $\log \ell_{\text{peak}}^{\text{ssc},1}$ | $\log \ell_{\text{peak}}^{\text{ssc},2}$ | $\log \ell_{\text{peak}}^{\text{ssc},1} - \log \ell_{\text{peak}}^{\text{ssc},2}$ | ζ |
|-----------------------------|--------------------------------------|-------------------------------------|------------------------------------|--------|
| $10^{-4}$                  | -4.16 (N)                            | -6.63                               | -9.64                              | 9.1 $\times 10^{-2}$ |
|                             | -3.85 (A)                            | -6.10                               |                                    | -9.20  |
| $10^{-3}$                  | -3.16                                 | -4.65                               | -6.66                              | 8.6 $\times 10^{-3}$ |
|                             | -2.85                                 | -4.13                               |                                    | -6.20  |
| $10^{-2}$                  | -2.22                                 | -2.77                               | -3.84                              | 6.5 $\times 10^{-2}$ |
|                             | -1.97                                 | -2.40                               |                                    | -3.60  |

2.3. Tests

In this paragraph we will compare the analytical expressions given by eqs. (14), (20) and (25) with those obtained from the numerical code described in Mastichiadis & Kirk (1995, 1997), where we have used the full expression for the synchrotron and Compton emissivities (c.f. eqs. (6.33) and (2.48) in Rybicki & Lightman (1979) and BG70 respectively).

For the comparison we used $B = 4$ G, $R = 10^{17}$ cm, $\gamma_0 = 10^3$ and three indicative values of the electron injection compactness, i.e. $\ell_{\text{peak}}^{\text{syn}} = 10^{-4}, 10^{-3}$ and $10^{-2}$. Our results are summarized in Table 1 where the first and second value in each row correspond to the numerical and analytical ones respectively; the ratio ζ given by eq. (35) is also shown. The magnetic compactness for the parameters used here is $\ell_B = 0.052$. The first two examples fall into the ‘synchrotron dominated’ regime since $12\ell_{\text{peak}}^{\text{syn}}/\ell_B = 2.3 \times 10^{-2}$ and $2.3 \times 10^{-1}$ for $\ell_{\text{peak}}^{\text{ssc}} = 10^{-4}$ and $10^{-3}$ respectively. Although, for the highest $\ell_{\text{peak}}^{\text{ssc}}$ considered here, electrons cool preferentially through the ICS of synchrotron photons, we still find ζ < 1.

In particular, our approximation for the position of the synchrotron peak (see Section 2.2) is the main cause for the differences appearing in the first column of Table 1. In general however, our approximations used for the derivation of eqs. (14), (25) and (35) are reasonable, even in the case of $\ell_{\text{peak}}^{\text{ssc}} = 10^{-2}$, where $\nu_{\text{peak}}^\text{ssc,1} \approx 4(u_0 + u_e)$; we remind that our analysis neglects the energy density of SSC photons in the electron cooling.

3. One-zone SSC fit to the core emission of Cen A

The emission from the core of Cen A has the double-peaked shape observed in many blazars with the low- and high-energy humps peaking at the infrared and sub-MeV energy ranges respectively (Jourdain et al. 1993; Chiaberge et al. 2001). The one-zone SSC model, where relativistic electrons are responsible for the radiation observed in low and high energies has been successfully applied over the years to various blazars and recently to FRI galaxies such as M87 (Abdo et al. 2009). Although it is also the dominant interpretative scenario for the core emission of Cen A it cannot explain the observed SED up to the TeV energy range (Abdo et al. 2010a, Roustazadeh & Böttcher 2011). Since the emitting region is compact enough for significant absorption of TeV gamma-rays on the infrared photons produced inside the source Abdo et al (2010a, Sahakyan et al. 2013). Note also that before the detection of Cen A at VHE gamma-rays, its whole SED was successfully reproduced by single zone SSC models (Chiaberge et al. 2001).

In this paragraph we attempt a similar application to the MW emission of Cen A, having in mind though, that the second SSC photon generation emerges in the SED for (i) high enough electron injection compactnesses, (ii) small size of the emitting region and (iii) relatively low Lorentz factor of electrons 4 – see also eqs. (10), (25) and (34). We note also that the combination of the low electron Lorentz factor with weak magnetic fields, as often used in SSC models, implies that the second generation Compton scatterings occur only partially in the Thomson regime. For this reason, the second SSC component is expected to be much steeper than the first one.

Under the assumption of monoenergetic electron injection the parameters that must be determined in the context of a one-zone SSC model are five: $B$, $R$, $\delta$, $\gamma_0$ and $\ell_{\text{peak}}^{\text{syn}}$; for power-law and broken power-law injection the unknown parameters increase to seven and nine respectively – see e.g. Mastichiadis & Kirk (1997); Aleksić et al. (2012). Because of no detections of variability in the HE/VHE regimes, the available observational constraints are only four: (i) the ratio of the observed peak frequencies $\nu_{\text{peak}}^{\text{ssc,1}}/\nu_{\text{peak}}^{\text{syn}}$; (ii) the peak synchrotron frequency $\nu_{\text{peak}}^{\text{syn}} = 3.2 \times 10^{13}$ Hz; (iii) the ratio of the observed peak fluxes ($\nu_{\text{peak}}^{\text{ssc,1}}, \nu_{\text{peak}}^{\text{syn}}$); (iv) the synchrotron peak flux ($\nu_{\text{peak}}^{\text{syn}}$) $\approx 4 \times 10^{-10}$ erg cm$^{-2}$ s$^{-1}$. From constraints (i) and (ii) we can determine the injection Lorentz factor of electrons $\gamma_0$ and find a relation between the magnetic field strength $B$ and the Doppler factor $\delta$ respectively:

$$\gamma_0 = \sqrt{\frac{3\nu_{\text{peak}}^{\text{ssc,1}}}{\xi_{\text{peak}}^{\text{ssc,2}}}} = 1.1 \times 10^3$$

(36)

and

$$B = B_{\text{cr}} \frac{h \nu_{\text{peak}}^{\text{syn}}}{\delta \gamma_0 m_e c^2} = 8 \delta^{-1} \text{ G},$$

(37)

where we neglected the factor $1 + z$ due to the small value of the redshift ($z = 0.00183$). The ratio $1 + z$ of the electron to magnetic compactness is determined by constraint (iii) and eqs. (16) and (25)

$$\ell_{\text{peak}}^{\text{ssc}} = \frac{1}{\delta} \left[ 1 + \left( \frac{4c}{3\sqrt{3}} \frac{(\nu_{\text{peak}}^{\text{ssc,1}})}{(\nu_{\text{peak}}^{\text{syn}})} \right) \right] \approx 5$$

(38)

Combining constraint (iv) with eqs. (16), (37) and (38) leads to a relation between $R$ and $\delta$

$$R \approx 10^{15} \delta^{-1} \text{ cm}.$$  

(39)

Finally, using eqs. (38) and (39) we find

$$\ell_{\text{peak}}^{\text{ssc}} \approx 10^{-3} \delta^{-3}.$$  

(40)

Since the viewing angle of the jet is in the range $15^{\circ} - 80^{\circ}$ the Doppler factor cannot exceed the value 3.7, whereas 4 Here we imply monoenergetic injection at $\gamma_0$. In the case of steep power-law injection between $\gamma_{\text{min}}$ and $\gamma_{\text{max}}$, the minimum Lorentz factor of electrons plays the role of $\gamma_0$. 5

5
values as low as 0.52 have been used in the literature (Roustazadeh & Böttcher 2011). From this point on we will adopt the representative value $\delta = 1$, which for an angle $30^\circ$ implies a bulk Lorentz factor $\Gamma = 7$. The derived values ($\gamma_0 = 1.1 \times 10^3$, $B = 8$ G, $R = 10^{15}$ cm, $\epsilon_{\text{inj}} = 10^{-3}$ and $\delta = 1$) were then used as a stepping stone for a more detailed fit to the SED, where we assumed the injection of a steep power-law electron distribution for better reproducing the photon spectrum above 10$^{13}$ Hz. The parameter values, which are only slightly different from the analytical estimates, are listed in Table 2. In the same table are also listed for comparison reasons the values of the SSC fit by Abdo et al. (2010a). We note that the parameter values that we derived at the beginning of this section, the analytical expressions given by eqs. (26) and (34) predict a ratio $\sim 0.08$, which is in good agreement with the numerical one.

4. An attempt to fit the SED using the maximum possible Doppler factor ($\delta = 3.7$) would result in smaller values of $R$, $B$ and $\epsilon_{\text{inj}}$ than those listed in Table 2. This would suppress electron cooling, i.e. near/mid-infrared and X-ray observations could not be modelled unless one would assume the injection of a broken power-law electron distribution.

**4. Addition of a relativistic proton component**

In the previous section we showed that the one-zone SSC model fails to reproduce the core emission of Cen A for energies above a few GeV. A recent analysis of Fermi data from four years of observations resulted in the detection of HE emission up to $\sim 50$ GeV (Sahakyan et al. 2013). It is now believed that this part of the spectrum along with the TeV data is produced by a second component that originates either from a compact (sub-pc) or from the black hole magnetosphere (Rieger & Aharonian 2009), photoion and photopair production on background (UV or IR) (Kachelrieß et al. 2010) or SSC photons (Sahu et al. 2012), $\gamma$-ray induced cascades in dust torus surrounding the high-energy emitting source (Roustazadeh & Böttcher 2011), non-thermal emission from relativistic protons and electrons that are being injected and accelerated at the base of the jet and cool as they propagate along it (Reynoso et al. 2011), are proposed scenarios that fall into the first category, whereas scenarios such as inverse

**Table 2.** Parameter values for the one-zone SSC model fit to the SED of Cen A shown in Fig. 1. For comparison reasons, the respective values of the SSC fit by Abdo et al. (2010a) are also shown.

| Parameter | Model | Model (Abdo et al. 2010a) |
|-----------|-------|---------------------------|
| $R$ (cm)  | $4 \times 10^{15}$ | $3 \times 10^{15}$ |
| $B$ (G)   | 6     | 6.2                       |
| $\delta$  | 1     | 1                         |
| $\gamma_{\text{min}}$ | $1.3 \times 10^4$ | 300                      |
| $\gamma_{\text{br}}$ | $- $ | 800                       |
| $\gamma_{\text{max}}$ | $10^6$ | $10^8$                   |
| $p_{e,1}$ | $- $ | 1.8                       |
| $p_{e,2}$ | 4.3  | 4.3                       |
| $\epsilon_{\text{inj}}$ | $6.3 \times 10^{-3}$ | $8 \times 10^{-3}$       |
| $\epsilon_{\text{th}}$ | $4.6 \times 10^{-3}$ | $3.7 \times 10^{-3}$     |

**Fig. 1.** SED of the core emission from Cen A with an one-zone SSC fit. This includes non-simultaneous observations from low-to-high frequencies: filled triangles (TANAMI VLBI), open triangles (Swift-XRT/Swift -BAT), grey circles (1-year Fermi-LAT by Abdo et al. 2010a), black circles (4-year Fermi-LAT by Sahakyan et al. 2013), black filled squares (H.E.S.S. by Aharonian et al. 2009) and black open squares are archival data from Marconi et al. (2000). The solid line is our one-zone SSC model fit with same slightly different parameters than those used in Abdo et al. (2010a). For the parameters used see Table 2.
Compton scattering of background photons in the kpc-scale jet (Hardcastle & Croston 2011) belong to the second one.

Here we propose an alternative explanation for the TeV and the HE emission in the Fermi energy range, which may as well be labeled as a ‘compact origin’ scenario. We assume that inside the compact emission region (e.g. \( R = 4 \times 10^{15} \) cm) relativistic protons, that have been co-accelerated to high-energies along with the electrons, are being injected in the source. In a co-acceleration scenario the ratio of the maximum Lorentz factors achieved by electrons and protons is \( \gamma_{\text{e, max}} / \gamma_{\text{p, max}} \), as predicted for example by first order Fermi and stochastic acceleration models (see e.g. [Rieger et al. 2007]). For this reason and given that \( \gamma_{\text{e, max}} = 10^8 \) we assume that \( \gamma_{\text{p, max}} = 1.8 \times 10^9 \), which furthermore does not violate the Hillas criterion since the corresponding gyroradius is \( r_g = 4.5 \times 10^{14} \) cm. To reduce the number of free parameters in our model we further assume that the accelerated distributions of protons and electrons have the same power-law index (\( p_e = p_p \)), although the resulting photon spectrum is insensitive to the exact value \( p_p \).

In order to follow the evolution of a system where both relativistic electrons and protons are being injected with a constant rate in the emitting region we used the time-dependent numerical code as presented in Dimitrakoudis et al. (2012) – hereafter DMPR12. The various energy loss mechanisms for the different particle species that are included in our code are

- Electrons: synchrotron radiation; inverse Compton scattering
- Protons: synchrotron radiation; photo-pair (Bethe-Heitler pair production) and photo-pion interactions
- Neutrons: photo-pion interactions; decay into protons
- Photons: photon-photon absorption; synchrotron self-absorption
- Neutrinos: no interactions.

Photohadronic interactions are modelled using the results of Monte Carlo simulations. In particular, for Bethe-Heitler pair production the Monte Carlo results by Protheroe & Johnson (1996) were used – see also Mastichiadis et al. (2005). Photo-pion interactions were incorporated in the time-dependent code by using the results of the Monte Carlo event generator SOPHIA (Mücke et al. 2004).

### 4.1. Photon emission

As a starting template for the parameters describing the primary leptonic component, we first used the one presented in Table 2. Then, we added five more parameters that describe the relativistic proton component in order to fit the HE/VHE emission; we refer to this as Model 1. The main difference between Models 1 and 2 is the value of Doppler factor, which is assumed to be higher in the second model. Subsequently, this affects, as already stated in point (4) of the previous section, the values of other parameters such as the electron injection luminosity. The parameters we used for our model SEDs shown in Fig. 2 are listed in Table 3. In general, the addition of a relativistic proton component successfully explains the HE emission detected by the Fermi satellite by both of our models. However, the TeV emission detected by H.E.S.S. can be satisfactorily explained only by Model 2. A zoom in the \( \gamma \)-ray energy range of the SED along with the model spectra is shown in Fig. 3. In what follows, we will first discuss the common features of Models 1 and 2 and, then, we will comment on their differences.

In both models, gamma-ray emission is attributed to the synchrotron radiation from secondary electrons produced via Bethe-Heitler pair production and photopion interactions as well as to the \( \pi^0 \) decay. The hardening of the spectrum at \( E \sim 0.4 \) GeV, in both cases, is caused by photon-photon absorption. This also explains the weak dependence of the resulting model fit on the slope of the proton distribution.

In the present treatment we consider only the internally produced photons (synchrotron and SSC) as targets for photohadronic interactions.

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Fig. 2. Leptohadronic fit of the MW core emission of Cen A using the parameter sets shown in Table 3. Models 1 and 2 are shown with dotted and dashed-dotted lines, respectively. For comparison reasons, the one-zone SSC fit shown in Fig. 1 is overplotted with solid line. All other symbols are the same as in Fig. 1.

Fig. 3. Zoom in the \( \gamma \)-ray energy range of the MW core spectrum of Cen A. The model spectra are overplotted with different line types marked on the plot.
Table 3. Parameter values used for our model SED shown in Fig. 2

| Parameter          | Model 1       | Model 2       |
|--------------------|---------------|---------------|
| $R$ (cm)           | $4 \times 10^{15}$ | $2.2 \times 10^{15}$ |
| $B$ (G)            | 6             | 3.5           |
| $t_{\text{tr}}$    | $1.3 \times 10^5$ | $7.3 \times 10^4$ |
| $\delta$           | 1             | 2             |
| $\Gamma$           | 7             | 7             |
| $\theta$           | $35^\circ$    | $20^\circ$    |
| $\nu_{\text{esc}}/t_{\text{tr}}$ | 1             | 4             |
| $\gamma_{\text{p, min}}$ | $1.3 \times 10^3$ | $1.3 \times 10^3$ |
| $\gamma_{\text{p, max}}$ | $10^6$        | $10^6$        |
| $\gamma_{\text{e, min}}$ | 4.3           | 4.5           |
| $\gamma_{\text{e, max}}$ | 7            | 7             |
| $\epsilon_{\text{p}}^{(0)}$ | $6.3 \times 10^{-3}$ | $7.9 \times 10^{-4}$ |
| $\epsilon_{\text{p, esc}}/t_{\text{tr}}$ | 1             | 5             |
| $\gamma_{\text{p, min}}$ | $2 \times 10^7$ | $2 \times 10^7$ |
| $\gamma_{\text{p, max}}$ | $1.8 \times 10^9$ | $1.8 \times 10^9$ |
| $p_{\text{p}}$     | 4.3           | 4.5           |
| $\epsilon_{\text{e}}^{(0)}$ | $4 \times 10^{-6}$ | $7.9 \times 10^{-7}$ |
| $u_{\text{e}}$ (erg/cm$^3$) | 12.3          | 2.6           |
| $u_{\text{e}}$ (erg/cm$^3$) | 1.9           | 2.3           |
| $u_{\text{p}}$ (erg/cm$^3$) | 6.8           | 15.4          |
| $u_{\text{p}}$ (erg/cm$^3$) | 1.4           | 0.5           |
| $L_{\text{cool}}^{(0)}$ (erg/s) | $1.2 \times 10^{44}$ | $1.3 \times 10^{44}$ |
| $L_{\text{p}}^{(0)}$ (erg/s) | $1.4 \times 10^{43}$ | $2.4 \times 10^{43}$ |
| $L_{\text{p}}^{(0)}$ (erg/s) | $2.5 \times 10^{43}$ | $2.5 \times 10^{43}$ |

$^a$ Here $\gamma_{\text{p, max}} \approx (m_{p}/m_{e}) \gamma_{\text{e, max}}$.

$^b$ The energy densities refer to the steady state of the system as measured in the comoving frame.

$^c$ The values refer to observed luminosities.

Table 3. Parameter values used for our model SED shown in Fig. 2

Fig. 4. Contribution of the photohadronic processes to the high energy part of the spectrum. Our model spectra when all processes are included are shown with solid lines, whereas when photopair and photopion processes are separately neglected are shown with dashed and dotted lines, respectively. The dash-dotted curve corresponds to the proton synchrotron emission. For the parameters used see Model 1 in Table 3.

Model 1

No Bethe-Heitler

No p

for photopair and photopion interactions with the relativistic protons, although external photon fields, such as the radiation from the accretion disk and/or the scattered emission from the Broad Line Region, could also be important (Atoyan & Dermer 2003a). The number density of synchrotron photons scales as $n_{\text{syn}}(\epsilon) \propto \epsilon^{-3/2}$ for $\epsilon_{\text{cool}} \ll \epsilon < \epsilon_{\text{peak}}$ where $\epsilon_{\text{peak}} \approx 2.4 \times 10^{-5}$ and $\epsilon_{\text{peak}} = 2.4 \times 10^{-7}$. Thus, protons with Lorentz factors down to $\gamma_{\text{p}} \gtrsim 2/\epsilon_{\text{peak}}$ or $\approx 8 \times 10^6$ can interact with this photon field through Bethe-Heitler pair production. Synchrotron protons cannot, however, serve as targets for photopion interactions, since this would require $\gamma_{\text{p, peak}} \gtrsim m_{p}/m_{e}$ or equivalently $\gamma_{\text{p}} \gtrsim \gamma_{\text{p, max}}$. Thus, pion production is solely attributed to interactions of protons with the SSC photon field (see also Sahu et al. 2012). For example, protons with Lorentz factors $\gamma_{\text{p}} \gtrsim 1.4 \times 10^4$ and $1.4 \times 10^7$ can interact with the upper ($\epsilon_{\text{esc,1}} \approx 0.2$) and lower ($\epsilon_{\text{esc,1}} \approx 2 \times 10^{-5}$) cutoff of the SSC photon distribution, respectively. For a fixed proton energy, the efficiency of both photopair and photopion interactions depends on the number density of the target field. For the particular set of parameters, one expects that interactions between the high-energy part of the proton distribution and the low-energy part of the photon fields are more efficient in the production of $\gamma$-rays. This is illustrated in Fig. 4 where the emitted spectra of Model 1 are shown when (i) all processes are included (solid line), (ii) Bethe-Heitler pair-production (dashed line) and (iii) photopion production (dotted line) are omitted. It becomes evident that the main contribution to the high-energy part of the spectrum comes from the Bethe-Heitler pair production process. Moreover, the proton synchrotron emission is by many orders of magnitude lower than the emission from the other components of hadronic origin – see dash-dotted line in the same figure.

For the values of $\gamma_{\text{p, min}}$ and $p_{\text{p}}$ used in the fit, the required injection compactness for obtaining an observable high-energy emission signature is $L_{\text{p}}^{(0)} = 4 \times 10^{-6}$ and $7.9 \times 10^{-7}$ for Models 1 and 2, respectively. This corresponds to observed injection luminosities $L_{\text{p, inj}}^{(0)} \approx 1.4 \times 10^{43}$ erg/s and $2.4 \times 10^{43}$ erg/s for the two models, respectively.5

For a black hole mass $M_{\text{BH}} = 5 \times 10^{7} M_{\odot}$ (Marconi et al. 2006; Neumayer 2010) the Eddington luminosity is $L_{\text{Edd}} = 6.5 \times 10^{45} (M_{\text{BH}}/M_{\odot})$ erg/s and, therefore, the proton injection luminosity in both models is only a fraction of it, i.e. $L_{\text{p, inj}}^{(0)} \approx \xi L_{\text{Edd}}$ with $\xi \approx 10^{-3}$. We note also that the required luminosity of the relativistic proton component is comparable to that of the leptonic one and, therefore, low compared to the values $10^{17} - 10^{18}$ erg/s that are inferred from typical hadronic modelling of blazars (see e.g. Bottcher et al. 2013). For the chosen parameters the emitting region is particle dominated with $u_{\text{p}} + u_{\text{e}} \approx \eta_{\text{p}} u_{\text{e}}$, where $\eta_{1} = 6$ and $\eta_{2} = 36$ for Models 1 and 2, respectively. We note also that the radiative efficiency $\eta_{\gamma}$, which we define as $\eta_{\gamma} = L_{\gamma}/(L_{\text{e, inj}}^{(0)} + L_{\text{p, inj}}^{(0)})$, is high for both models; specifically, the values listed in Table 3 indicate $\eta_{\gamma,1} = 0.98$ and $\eta_{\gamma,2} = 0.68$.

In both models we have used a high value for the minimum proton Lorentz factor, which cannot be explained by

5 For the calculation we used the definition of the proton injection compactness $L_{\text{p, inj}}^{(0)} = L_{\text{p, inj}}^{(0)} \delta (\pi R \delta m_{p} c^{3})$, where the factor $\delta$ takes into account Doppler boosting effects for radiation emitted from a spherical volume.
any theoretical model of particle injection and acceleration. However, any effort to extend such a steep power-law distribution \( (p_p = 4.3 - 4.5) \) down to \( \gamma_p = 1 \) is excluded from the energetics. As an indicative example, we used the parameter values of Model 1 listed in Table 3 except for a lower value of the minimum Lorentz factor. In order to obtain a good fit to the SED for \( \gamma_{p,\text{min}} = 2 \times 10^3 \), the required proton injection luminosity increases by almost three orders of magnitude, i.e. \( L_p^{inj} = 6 \times 10^{45} \) erg/s. Since there is no physical reason for such high values of the minimum proton energy, one can interpret it as the break energy of a broken power-law distribution. In such case, the power-law below the break must be rather flat, e.g. \( p_p = 1.5 - 2.0 \), in order to avoid too high proton luminosities. A detailed fit using broken power-law energy spectra lies, however, outside the scope of this work. In any case, since there is no known plausible physical scenario that predicts either high values of \( \gamma_{p,\text{min}} \) or broken power-law energy spectra with \( \Delta p_p \geq 2.5 \), the sub-Eddington proton luminosities listed in Table 3 can be considered as a lower limit of those retrieved using a more realistic proton distribution.

The key difference between Models 1 and 2 is the assumed value of the Doppler factor. In Model 1, where we did not allow any Doppler boosting of the emitted radiation \( (\delta = 1) \), we cannot explain the VHE emission. However, by assuming a slightly higher value for the Doppler factor the intrinsic absorbed spectrum is boosted by a factor \( \sim \delta \) in frequency and of \( \sim \delta^3 \) in flux, respectively. The boosting effect when combined with the fact that all other parameter values are of the same order of magnitude as those of Model 1, results in a model spectrum that satisfactorily goes through the H.E.S.S. data points. In the light of the recent analysis of the four-year Fermi-LAT data (Sahakyan et al. 2013) that implies a common origin of the HE and VHE emission, we believe that Model 2 describes better the emitting region of the core. Note that the connection between the GeV and TeV emission could not be suggested by the previously available one-year Fermi-LAT observations (Abdo et al. 2010a) – see grey circles in Fig. 3.

### 4.2. Neutrino and UHECR emission

The detailed neutrino spectra (of all flavors) obtained using the numerical code of DMPR12 for both models listed in Table 1 are shown in Fig. 3. The neutrino spectra from both models peak at \( \sim 10^8 \) GeV, while above that energy they can be approximated as power-laws with slopes \( p_\nu \approx 1.5 \) and \( \sim 1.6 \), respectively. This is in agreement with the approximate relation \( p_\nu \approx (p_p = 0.5)/2.5 \) derived in DMPR12. The steepening of the spectra above \( 3 \times 10^7 \) GeV (Model 1) and \( 10^8 \) GeV (Model 2) is due to the cutoff of the proton injection distribution. Although photohadronic processes are significant in modelling the photon spectra above a few GeV, the peak fluxes of neutrinos emitted through the charged pion and muon decay are far below the upper limit of the IceCube 40-string (IC-40) configuration (Abdasi et al. 2011) – see grey line in the same figure. The neutrino production efficiency that is defined as \( \eta_\nu = L_\nu/(L_p^{inj} + L_p^{up}) \), is approximately \( 2 \times 10^{-5} \) and \( 2 \times 10^{-7} \) for Models 1 and 2, respectively. Thus, we find that \( \eta_\nu \ll \eta_p \), where the radiative efficiency was found to be \( \sim 0.8 \). This differentiates the leptohadronic models presented here from others applied to blazar emission, where neutrino efficiencies as high as 0.1 can be obtained (see e.g. Dimitrakoudis et al. 2013 for the case of Mrk 421). In general, there is no case where \( \eta_\nu \approx \eta_p \) (e.g. Reimer 2011) and such low values are to be expected in cases of strong magnetic fields, weak target photon fields and/or low proton injection compactness; the latter applies to our case.

Cen A has been under consideration as a potential source of ultra high energy cosmic rays (UHECR) from as early as 1978 (Cavallo 1978), and its proximity to our galaxy compared to all other AGN has even inspired models where it is the sole originator of UHECR (Biermann & de Souza 2012). Recently, the Pierre Auger Observatory (PAO) has shown an excess in UHECR within 18° of Cen A (Pierre Auger Collaboration et al. 2007) and, although that region contains a high density of nearby galaxies, further analysis has shown that some of those UHECR may have originated from Cen A itself (Farrar et al. 2013; Kim 2013). For our two models we have obtained distributions for both the escaping protons and neutrinos. While the former are susceptible to adiabatic energy losses, and thus any calculation of their flux would constitute an optimistic upper limit, the latter can escape unimpeded and decay into protons well away from the core (Kirk & Mastichiadis 1985; Begelman et al. 1990; Giovannini & Kazanas 1990; Atwood & Dermer 2003). In Fig. 6 we have plotted the flux of protons resulting from the decay of neutrons that escape from the emitting region. Since we have not treated cosmic-ray (CR) diffusion in the intergalactic magnetic field, which generally decreases the CR flux that arrives at earth, our model spectra should be considered only as an upper limit. For both models, the peak fluxes are far lower than the observational limit of PAO. Although that makes Cen A as core an unlikely source of UHECR, those could potentially originate from its lobes instead (e.g. Gopal-Krishna et al. 2010).
by a low value of the Doppler factor (searchers agree that the emitting source is characterized been observed both at GeV and TeV energies. Most re-

sources. One could, by reversing the above arguments, find

to fit the SED of the nearby radio galaxy Cen A, that has

to fit, with varying degrees of success, the SED of

blazars. The discovery of high energy emission from an-

other class of AGN, i.e. that of radio galaxies, poses new

challenges to these models: if radio galaxies are misaligned

blazars, then the observed emission should come from a re-
gion moving with a relatively large angle with respect to our

line-of-sight. This implies a rather small value for the

Doppler factor that, for a given flux level of the source, can

be compensated only by a large value of the so-called

electron compactness parameter.

It is well known that sources with high electron, and

sequently high photon compactnesses, are subject to strong

Compton scattering. This usually leads to higher order genera-
tions of SSC, while, in extreme conditions it might lead to the ‘Compton catastrophe’. As clearly these conditions are not apparent in the MW spectra of radio galaxies, one could, by reversing the above arguments, find limits on the parameters used to model the SED of these sources.

As an example, in the present paper we have attempted to fit the SED of the nearby radio galaxy Cen A, that has been observed both at GeV and TeV energies. Most re-

searchers agree that the emitting source is characterized

by a low value of the Doppler factor ($\delta \approx 1 - 3$). In or-
der to show the relevance of the first and the second SSC components, we have calculated analytically in Section 2 the spectral luminosities at the peaks of these components. Under the assumption that all scatterings producing the

first SSC component occur in the Thomson regime, i.e. a

condition that can be easily satisfied in most of the rele-

vant cases, we found that the SSC dominates synchrotron cooling whenever $\gamma^{\text{inj}}_e \geq \ell_e / 12$, where $\gamma^{\text{inj}}_e$ and $\ell_e$ are the electron and magnetic compactnesses respectively. The cal-

ulation of the luminosity of the second SSC component is

more complicated as scatterings occur in both the Thomson and Klein-Nishina regimes. However, adopting the oft used cut-off approximation for the latter, we were able to find a closed expression for the luminosity which, in addition, agrees well with numerical calculations – the same can be said for the other two components (i.e. synchrotron and first SSC) as evidenced by Table [I].

Using the relations described above as a stepping stone, we have obtained in Section 3 a fit to the SED of Cen A. Limiting the Doppler factors by necessity to small values, we found that the one-zone SSC model can successfully fit the SED up to $10^{23}$ Hz. At that frequency the peak of the second SSC component appears, which is then followed by a steep power-law segment due to Klein-Nishina effects. This causes, typical one-zone SSC modelling to fail at fitting the high energy observations of Cen A.

In order to fit the emission at frequencies above $10^{23}$ Hz, we have introduced, in Section 4.1, a hadronic component which, we assume, is co-accelerated to high energies along with the leptonic one. Assuming that the two populations share the same characteristics, i.e. their injection power-
laws have the same slope and their maximum cutoffs are related to each other through a simple relation stemming from the Fermi acceleration processes, we found that ac-
teptable fits to the SED of Cen A can be obtained for proton injection luminosities of the same order of magnitude as the electron one (see Table [3]). Interestingly enough, fits using $\delta = 2$ can attribute the TeV observations to hadronic emission, while fits with $\delta = 1$ fail to do so due to strong photon-photon attenuation.

In Section 4.1 we have also showed that $\gamma^{\text{inj}}_p \gg \gamma^{\text{inj}}_e$ in order to obtain the required radiative efficiency of the pho-

tohadronic interactions under the assumption of a steep power-law distribution for protons and the requirement of a sub-Eddington proton injection luminosity. On the one hand, such high values of $\gamma^{\text{inj}}_p$ may be interpreted as the break energy of a broken-power law at injection. On the other hand, one could, in principle, reconcile the hypothet-
ical low values of $\gamma^{\text{inj}}_p$ and the high values of $L^{\text{inj}}_p$ by considering also as targets for photohadronic interactions external photon fields, such as diffuse and/or line emission from the Broad Line Region (BLR). In the case of Cen A, however, the lack of strong broad emission lines implies that these photon fields are negligible (Alexander et al. 1999, Chiaberge et al. 2001). Another possible photon target field could be the mid-IR radiation that is believed to be associ-
ated with cool dust in the nuclear region of Cen A (e.g. Karovska et al. 2003). For the observed fluxes, which range from 1 to 100 Jy (Israel 1998, Karovska et al. 2003), the number density of mid-IR photons as measured in the rest frame of the high-energy emitting region is by many orders of magnitude lower than that of the internally produced synchrotron photons. Thus, incorporating the IR photon field in the calculations presented here would not lower the requirement of high proton luminosities.

The consideration of relativistic protons in the emitting region is inevitably related to the neutrino emission, since proton interactions with the photon fields present in the
source result in charged meson production. In Section 4.2 we have presented the neutrino spectra calculated for both our models. For the employed parameters the efficiency of pion production is very low and this can also be seen at the low peak neutrino fluxes which are by many orders of magnitude below the IceCube upper limit.

Furthermore, high energy neutrons resulting from photopion interactions are an effective means of facilitating proton escape from the system, as they are unaffected by its magnetic field and their decay time is long enough to allow them to escape freely before reverting to protons (e.g. Kirk & Mestiriakos 1998). Those effects make them advantageous as they are unaffected by adiabatic energy losses that the protons may sustain in the system as it expands (Rachen & Mestiriakos 1998). Those effects make them excellent candidates of UHE protons. For our model parameters, i.e. steep injection proton spectra and small values of the Doppler factor, the obtained proton distributions peak in the range $10^{16} - 10^{17}$ eV, where the effects of CR diffusion in the intergalactic magnetic field cannot be neglected. Since in the present work we have not treated CR diffusion, our results should be considered as an upper limit. Still, these are well below the observed CR flux at such energies. In the light of recent results suggesting Cen A to be the origin of some UHECR events observed by PAO (Farrar et al. 2013; Kim 2013) and our model results, the core of Cen A cannot be the production site of UHECR.

Our analysis has shown that Cen A can be explained by means of a lepto-hadronic model as was the case of Mrk 421 (Mestiriakos et al. 2013). However, contrary to that source, a one zone SSC model fails to reproduce theSED of Cen A mainly due to complications arising from the appearance of the second SSC component. Although this feature has been overlooked by many researchers it may play a crucial role in fitting the SEDs of radio galaxies, as these require high electron luminosities, making the conditions very favourable for its appearance.

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