Lepto-hadronic Jet-disk Model for the Multiwavelength SED of M87

Margot Boughelilba1, Anita Reimer1, and Lukas Merten1,2

1 Institute for Astro and Particle Physics, University of Innsbruck, 6020 Innsbruck, Austria; margot.boughelilba@uibk.ac.at
2 Ruhr-Universität Bochum, Institut für Theoretische Physik IV, D-44801 Bochum, Germany

Received 2022 May 31; revised 2022 August 12; accepted 2022 August 30; published 2022 October 14

Abstract

The low-luminosity active galactic nuclei M87, archetype of Fanaroff–Riley I radio galaxies, was observed in a historically quiet state in 2017. While one-zone leptonic jet models alone cannot explain the core radio-to-gamma-ray spectrum, we explore a hybrid jet-disk scenario. In this work, we model the overall spectral energy distribution of M87’s core with a dominating one-zone lepto-hadronic jet component, coupled with the contribution from the accretion flow. We find close-to-equipartition parameter sets for which the jet component fits the radio-to-optical data as well as the gamma-ray band, while the accretion flow mainly contributes to the X-ray band. The effects of gamma-ray absorption by the extragalactic background light during the propagation toward Earth are probed and are found to be negligible for this model. The neutrino flux produced by such scenarios is also calculated, but remains below the current instruments’ sensitivity.

Unified Astronomy Thesaurus concepts: Relativistic jets (1390); Non-thermal radiation sources (1119); Gamma-ray sources (633); Cosmic ray sources (328); Astrophysical black holes (98); Low-luminosity active galactic nuclei (2033); High energy astrophysics (739); Active galactic nuclei (16); Particle astrophysics (96)

1. Introduction

M87 is one of the closest examples of low-luminosity active galactic nuclei (AGN), located at a distance of \( \sim 16.8 \) Mpc from Earth (corresponding to a redshift \( z \approx 0.004 \)) in the Virgo cluster. The mass of the supermassive black hole at its center was estimated around \( 6.5 \times 10^8 \) \( M_\odot \) (Event Horizon Telescope Collaboration et al. 2019a). In 2017, an extensive multiwavelength observation campaign was launched, taking quasi-simultaneous data from several telescopes over the entire electromagnetic band (EHT MWL Science Working Group et al. 2021). For nearly 2 months, M87’s core region was observed in a particularly low state. These observations allow studying the innermost radiation from the AGN, in particular, the launching region of the jet that M87 exhibits, as the broadband spectrum of these observations is dominated by the emission from the core and not from the jet’s knots such as HST-1. Furthermore, the closeness and the size of M87 make it a prime candidate accelerator of the observed high-energy cosmic rays (see, e.g., Biermann et al. 2000; Protheroe et al. 2003). To explain the multiwavelength spectral energy distribution (SED) of M87, different emission models are usually probed. Leptonic jet models are, for the case of M87, typically based on the synchrotron self-Compton mechanism: synchrotron photons produced by the interaction between the jet’s relativistic electrons and positrons with the ambient magnetic field are used as a target field for inverse-Compton scattering by the same particles, thereby producing high-energy radiation. In EHT MWL Science Working Group et al. (2021), two different one-zone leptonic models were applied, but failed to reproduce both the high and low energy parts of the SED of M87 at the same time.

On the other hand, lepto-hadronic models have been proposed to explain the SED of objects such as M87 (e.g., Reimer et al. 2004). In such models, accelerated protons are also present in the jet together with electrons, and the high-energy part of the SED is assumed to be the result of proton-initiated processes. A clear observational signature between these two kinds of models is the production of neutrinos in the case of lepto-hadronic models. In this paper, we explore a global model coupling the jet lepto-hadronic emission and the accretion flow, in order to explain the observed SED of M87. With this model, all the emission would originate from the core region of the AGN.

This paper is organized as follows: in Section 2 we describe the jet model, then in Section 3 we detail the accretion flow model component and its parameters. Results of the simulations are presented in Section 4 and we conclude discussing the global core emission in Section 5.

2. Jet Model Component

As of today, the one-sided jet launching from M87 has been well-studied in all different wavelengths. In the 2017 observation campaign, EHT MWL Science Working Group et al. (2021) did not infer any time variability in the flux above 350 GeV. The data collected focus on the core emission. The angular resolution of radio observations suggests that the radio emission region is close to the jet launching region. While launch mechanisms are still unclear, the total estimated jet power for M87 of \( \sim 10^{43} \text{--} 44 \) erg s\(^{-1}\) (Stawarz et al. 2006; de Gasperin et al. 2012; Prieto et al. 2016) can be provided through, e.g., the Blandford–Znajek mechanism (Blandford & Znajek 1977; Event Horizon Telescope Collaboration et al. 2019b).

In this paper, we explore models that can reproduce the quiet and steady state of M87’s core observed between 2017 March and May. There is evidence of sub- to superluminal motion of radiating jet components in M87’s inner jet (e.g., Walker et al. 2018; Snijs et al. 2019) that can support a jet model setting in which the emission region is viewed as a moving blob. On the other hand, the possibility that the jet is a continuous zone in...
which the particles flow is often considered (see, e.g., Blandford & Königl 1979; Massi 2011) in the case of quiet state emission, and cannot be ruled out. We investigate both scenarios here. First we consider the jet emission region as a spherical blob with a constant radius \( r'_\text{em} \) of magnetized plasma moving at a mildly relativistic speed along the axis of a non-expanding jet during the observation time, inclined by an angle \( \theta \) with respect to the line of sight. This defines a Doppler factor \( \delta = \Gamma_j^{-1}(1 - \beta_j \cos \theta)^{-1} \) where \( \Gamma_j \) and \( \beta_j \) are the bulk Lorentz factor and velocity, respectively. For the second scenario, we consider the jet as a continuous cylinder of radius \( r'_\text{em} \) and proper length \( l' = \Gamma_l l \) with \( l \) being the observed length.

The EHT observation provides a strong constraint on the size of the emission region, as the angular resolution allows probing the closest regions to the black hole. At 230 GHz, the radio flux was measured with an angular resolution \( \theta_{\text{obs}} \) of 0.0606, corresponding to 7.5 \( r_g \) (for M87, \( r_g = GM_{\odot}/c^2 \approx 9.8 \times 10^{14} \) cm) in radius. However, even for a mildly relativistic jet velocity, the blob travels farther than 7.5 \( r_g \) over the observation time. In the continuous jet scenario, we assume that the jet is launched around the innermost stable orbit of the black hole, i.e., 6 \( r_g \) for a static black hole. Hence, when observing the core region within 7.5 \( r_g \) at 230 GHz the jet component is likely not the dominant one. Since we choose to focus on the core emission, we take care that for an emission region of size \( \lesssim 7 \ r_g \), the predicted radio flux does not exceed this particular data point.

Furthermore, the SED of M87 indicates a self-absorbed, stratified jet below at least 86 GHz (Blandford & Königl 1979; EHT MWL Science Working Group et al. 2021). This lower limit on the self-absorption frequency \( \nu_{\text{SSA,obs}} \) and corresponding flux \( S_{\nu_{\text{SSA,obs}}} \) coupled with the estimate of the size of the emission region allows deriving a relation for the magnetic field strength \( B \) required. Following the treatment by Kino et al. (2014) for a moving blob

\[
B = b(p) \left( \frac{\nu_{\text{SSA,obs}}}{1 \text{GHz}} \right)^{5} \left( \frac{\theta_{\text{obs}}}{1 \text{mas}} \right)^{-1} G, \tag{1}
\]

with \( b(p) \) described in Appendix A. From this, we estimate an order of magnitude for the magnetic field strength and then adjust the primary electron injection parameter so that the synchrotron radiation produced is of the order of \( S_{\nu_{\text{SSA,obs}}} \) at the given frequency. Considering that the self-absorption frequency is \( \lesssim 230 \) GHz (around the EHT data point), with the observed flux being \( S_{\nu_{\text{SSA,obs}}} \approx 0.6 \) Jy, gives an estimate for the magnetic field strength between \( \sim 5 \) and \( 60 \) G.

We assume that the emission region contains primary relativistic electrons and protons that are isotropically and homogeneously distributed in the comoving jet frame, and following a power-law energy spectrum cutting off exponentially, such that the spectral number density \( n'_\text{e,p}(E') \propto E'^{-p} e^{-E'/E'_{\text{max}}(e,p)} \text{ cm}^{-3} \), for \( E' \geq E'_{\text{min}}(e,p) \) (where \( e,p \) denotes the electrons or the protons, respectively).

These primary particles are injected continuously into the emission region at a rate of \( q_i \) (\( \text{cm}^{-3} \text{ s}^{-1} \)), where they suffer from different interactions. These are photomeson production, Bethe–Heitler pair production, inverse-Compton scattering, \( \gamma-\gamma \) pair production, decay of all unstable particles, synchrotron radiation (from electrons and positrons, protons, and \( \pi^+ \), \( \mu^\pm \), and \( K^\pm \) before their respective decays) and particle escape. Positrons are treated the same way as electrons; hence, in the following we will use electrons to refer to the two populations irrespective of their type.

Primary particles can also interact with external target photon fields (i.e., produced outside the jet). However, no evidence of a dusty torus has been found (Perlman et al. 2007) and no Fe Kα line has been observed to support the existence of a strong broad-line region (BLR) component (Di Matteo et al. 2003). This is in line with the properties of true type 2 AGN (Laor 2003; or see Ho 2008 for a review). Hence, we do not consider the dusty torus nor the BLR as external target fields. On the other hand, the accretion flow could serve as an external photon field for the jet’s primary particles, a possibility we discuss in Section 4.

The maximum energy of the primary particles is determined by \( E_{\text{max}} = \max(E_{\text{Hillas,max}}', E_{\text{max,loss}}', E_{\text{max,acc}}') \), where \( E_{\text{Hillas,max}}' \) is the energy given by the Hillas criterion (Hillas 1984) and \( E_{\text{max,loss}}' \) is the energy obtained by balancing the particles’ acceleration and loss rates.

The Hillas criterion constrains the Larmor radius of the particles to be smaller or equal to the size of their acceleration region, leading to an estimate of the maximum particle energy \( E_{\text{Hillas,max}}' \approx 10^{21}Z^2/(\nu/\text{pc})(B/G)\text{eV} \).

Expressions for \( E_{\text{max,loss}}' \) and \( E_{\text{max,acc}}' \) are obtained by equating the acceleration timescale \( t_{\text{acc}}(E_{\text{max,acc}}') \) and the loss timescale \( t_{\text{cool}}(E_{\text{max,loss}}') \) respectively. We follow the work of Reimer et al. (2004) to verify that the ratio of the two maximum energies \( E_{\text{max,loss}}'/E_{\text{max,acc}}' = (m_p/m_e)^{4/3}(\beta-1) \) is obtainable with a realistic turbulence spectrum. For Kolmogorov diffusion, \( \beta = 5/3 \), we get \( E_{\text{max,loss}}'/E_{\text{max,acc}}' \approx 6 \times 10^9 \). Bohm diffusion, where the magnetic field is fully tangled, corresponds to \( \beta = 1 \), and in the case of strong magnetic fields, Kraichnan turbulence \( \beta = 3/2 \) can be present (Kraichnan 1965).

To compute the time-dependent direct emission and cascade component from the jet’s particles, we use a particle and radiation transport code (see, e.g., Reimer et al. 2019) that is based on the matrix multiplication method described in Protheroe & Staney (1993) and Protheroe & Johnson (1996). The interaction rates and secondary particles and photons yields are calculated by Monte Carlo event generator simulations (except for synchrotron radiation, for which they are calculated semi-analytically). These are then used to transfer matrices that describe how each particle spectrum will change after a given timestep \( \delta t \). To ensure numerical stability, we set \( \delta t \) equal to the smallest interaction time for any given simulation. In each timestep, energy conservation is verified. For steady-state spectra, we run the simulation until we reach convergence, which we define here as the ratio \( R_{\text{conv}} \) between the flux at a simulation time \( t \) and the flux at a simulation time \( t - \delta t \). Convergence is reached when \( R_{\text{conv}} = F_t/(t + \delta t)/F_t(t) < 1 \pm 10^{-3} \).

All the calculations listed above are done in the jet frame. The observed spectrum \( \nu F_\nu \) is then given by the frame transformation \( F_\nu = (1 + z)g_{\text{boost}}\nu_{\text{obs}}/4\pi d_L^2 \), where \( \nu_{\text{obs}} \) is the comoving luminosity from the jet with \( d_L = 16.8 \) Mpc the luminosity distance of the source and \( g_{\text{boost}} = \delta \nu'/(1+z) \). The Doppler enhancement factor is \( g_{\text{boost}} = \delta \nu'/\Gamma \) for a moving blob and \( g_{\text{boost}} = \delta \nu^2/\Gamma \) for a continuous jet (Sikora et al. 1997; Stawarz et al. 2003). For a given comoving energy density, we obtain the intrinsic luminosity through \( \nu' F_{\nu'} = (\nu_{\text{em}}/c)(L'_{\nu'}/V') \).
where $V'$ is the comoving volume of the emission region (i.e., depending on the geometry). We find that we can obtain the same observed flux for both jet configurations by setting the length of the continuous cylinder to \( l' = 2\delta_l \Gamma r \Gamma_{cm}/3 \). We apply this for the remaining part of this work; hence, the results that we show in Section 4 are identical for the moving blob and the continuous jet scenario, given this condition. The effect of gamma-ray absorption by the extragalactic background light (EBL) on the escaping photon beam traveling from the source to Earth is taken into account. Three different models, using different approaches to calculate the EBL SED as a function of the redshift are used here to compute the flux attenuation factor. We use the models of Franceschini et al. (2008), Domínguez et al. (2011), and Gilmore et al. (2012), which are based on existing galaxy populations and extrapolate them back in time, based on the evolution of galaxy populations directly observed over the range of redshifts that contribute the most significantly to the EBL, and based on forward evolution of galaxy populations starting with cosmological initial conditions, respectively. The high-energy flux of M87 can be used to probe these models, and constrain the EBL density, especially in the far-infrared band, where the differences are especially between the models. However, we find that due to the distance of M87, the effects of gamma-ray absorption are negligible for gamma-rays with an energy lower than 10 TeV (\( \sim 10^{27} \) Hz). As we predict the emitted flux to peak at \( \sim 10^{24-25} \) Hz with a strong flux decrease toward higher energies (see Section 4), we hence cannot discriminate between any of the three models.

3. Accretion Flow

Low-luminosity AGNs like M87 are expected to host accretion flows around their SMBH that are radiatively inefficient. This is characterized by the formation of geometrically thick, optically thin, very hot accretion flows, called advection-dominated accretion flows (ADAFs; introduced by Rees et al. 1982, Ichimaru 1977 and further developed by, e.g., Narayan & Yi 1995; Abramowicz et al. 1995). ADAFs exist only when the accretion rate is sufficiently low (\( M \lesssim 0.01 M_{\text{Edd}} \)), and consist of a plasma of thermal electrons and ions, where both components may have different temperatures, \( T_e \) and \( T_p \), respectively. In addition to the ADAF, we assume the existence of a truncated standard thin accretion disk (Shakura & Sunyaev disk, Shakura & Sunyaev 1973) extending the outer parts of the ADAF. Here, we investigate how far an ADAF/disk system can contribute to the X-ray component, while not overshooting the radio-to-optical part of the SED that is considered to be jet dominated.

In the following, we use the quantities \( X_\eta = \frac{x}{107} \) and the normalized quantities \( r = r/R_s \), with the Schwarzschild’s radius \( R_s = 2\Gamma/\gamma_s \approx 2.95 \times 10^7 M_{\text{BH}}/M_{\odot} \), and \( \eta_s = M/M_{\text{Edd}} = \eta_{\text{Edd}} M_c^2/L_{\text{Edd}} \), where \( \eta_{\text{Edd}} \) is the radiation efficiency of the standard thin disk (\( \eta_{\text{Edd}} \approx 0.1 \)) and the Eddington luminosity \( L_{\text{Edd}} = 1.3 \times 10^{47} M_{\text{BH}} \text{ erg s}^{-1} \). We make use of the one-zone, height-integrated, self-similar solutions of the slim disk equations derived by Narayan & Yi (1995) to describe (see Appendix B) the hot plasma.

To obtain the spectrum emitted by an ADAF, the balance between the heating and cooling of the thermalized electrons present in the plasma \( q^{++} = q^{--} \) is solved to determine the scaled electron temperature \( \theta_e = k_b T_e/m_e c^2 \). Here, \( q^{++} \) is the electrons’ heating rate, and \( q^{--} \) is their cooling rate. The emission mechanisms that we consider in the following are synchrotron radiation, bremsstrahlung, and Comptonization of the two previous components. The total cooling rate is the sum of the three individual cooling rates, detailed in Appendices C.1 and C.2. The heating mechanisms and rates are described in Appendix D and they consist of Coulomb collision between ions and electrons, and viscous energy dissipation.

The plasma is a two-temperature plasma where the ion temperature is related to the electron temperature through \( T_i + 1.087 T_e \approx 6.66 \times 10^{12} \) \( \beta \) (Narayan & Yi 1995), where \( \beta \) is the ratio between the gas \( p_g \) and the total pressure \( p = p_e^{1/2} + p_i^{1/2} \) with \( p_{\text{me}} = B^2/8\pi \), and \( \rho \) is the mass density and \( B \) is the isotropically tangled magnetic field.

We obtain the electron temperature by varying \( T_e \) using a bisection method to solve the balance equation for each radius. Furthermore, we take \( \dot{m} \) of the form \( \dot{m} = \dot{m}_{\text{out}} (r/r_{\text{out}}) \), where \( r_{\text{out}} \) is the outer radius of the ADAF and is associated with an accretion rate \( \dot{m}_{\text{out}} \), and \( s \) is a mass-loss parameter (introduced by Blandford & Begelman 1999) that is used to include the presence of outflows or winds from the ADAF.

Upon obtaining the electron temperature, the emitted spectrum from the ADAF is computed, integrating over the radius of the ADAF. In order to take into account absorption, we follow the method of Manmoto et al. (1997), and derive the flux from synchrotron and bremsstrahlung emission as

\[
F_{\nu,0} = \frac{2\pi}{\sqrt{3}} B_\nu [1 - \exp(-2\sqrt{3} \tau_\gamma)] \text{ erg cm}^{-2} \text{ s}^{-1} \text{ Hz}^{-1}, \tag{2}
\]

where

\[
B_\nu = \frac{2h\nu^3}{c^2} \frac{1}{\delta \nu \kappa_\nu - 1}
\]

is the Planck’s function, and \( \tau_\gamma \) is the optical depth for absorption defined such that \( \tau_\gamma = (\sqrt{2}/2)\kappa_\gamma H \), with \( \kappa_\gamma = (j_{\gamma,\text{syn}} + j_{\gamma,\text{br}})/(4\pi B_\nu) \) the absorption coefficient. The emissivities \( j_{\gamma,\text{syn}} \) and \( j_{\gamma,\text{br}} \) are given in Appendix C.1.

| Parameter | Minimum Value | Maximum Value | Best Choice | Reference |
|-----------|---------------|---------------|-------------|-----------|
| $\alpha$ | 0.01 | 1 | 0.1 | 1, 2 |
| $\beta$ | 0.5 | <1 | 0.9 | 3, 4 |
| $m_{\text{out}}$ | $1 \times 10^{-4}$ | $2.93 \times 10^{-3}$ | $1.6 \times 10^{-3}$ | 5, 6 |
| $s$ | 0 | 1 | 0.39 | 6, 7 |
| $\delta$ | $10^{-4}$ | $10^{-1}$ | $5 \times 10^{-3}$ | 5, 8 |

Note. Here, $\alpha$ is the viscosity parameter introduced by Shakura & Sunyaev (1973), $\beta$ is the ratio between the gas and the total pressure, $m_{\text{out}}$ is the accretion rate at the outermost part of the ADAF, $s$ is the mass-loss parameter characterizing the evolution of the accretion rate over the volume, and $\delta$ is the fraction of viscous energy directly transmitted to the plasma electrons.

The upper limit is the Bondi accretion rate, calculated with the mass estimate of Event Horizon Telescope Collaboration et al. (2019a).

References—(1) Chael et al. (2019), (2) Martin et al. (2019), (3) Chael et al. (2018), (4) Ressler et al. (2017), (5) Di Matteo et al. (2003), (6) Nemmen et al. (2014), (7) Blandford & Begelman (1999), (8) Mahadevan (1997).
Hence, the local luminosity from synchrotron and bremsstrahlung at a given radius is given by \( L_{\nu,0} = 2\pi R^2 F_{\nu,0} \).

Synchrotron radiation and bremsstrahlung further act as a photon field for inverse-Compton scattering by the thermal electrons. Following the work of Kimura et al. (2015), we compute the number density of photons after the \( i \)th scattering:

\[
N_{\gamma,i}(\epsilon) = \frac{R_o}{\epsilon} \int d\gamma \frac{3}{2\gamma} N_e(\gamma, \theta_e) \times N_{\gamma,i-1}\left(\frac{3\gamma}{4\gamma_0^2}, \gamma\right),
\]

where \( R_e(\epsilon, \gamma) \) is the scattering rate for electrons with Lorentz factor \( \gamma \) and photons with dimensionless energy \( \epsilon = h\nu/(m_e c^2) \).
which we take from Coppi & Blandford (1990). $N_\gamma(\gamma, \theta_\gamma)$ is the Maxwellian distribution of electrons, described in Equation (C1). The initial condition is given by $N_{\gamma,0}(\epsilon) = L_{\gamma,0}/(h\nu\pi c R^2)$ with $L_{\gamma,0} = (m_e c^2/h)L_{\gamma,0}$.

The self-similar solutions provide a good estimate for the ADAF emission for sufficiently large radii ($r \gg r_{\text{sonic}}$, where $r_{\text{sonic}}$ is the sonic radius; Narayan et al. 1997); however, the inner part of the ADAF ($r \sim 2.5 - 4$) is thought to be at the origin of the ring observed by Event Horizon Telescope Collaboration et al. (2019b) at 230 GHz. We cannot use the self-similar solutions to account for this inner part emission, but considering the ADAF framework, we expect that synchrotron radiation is the dominant process in this region. The synchrotron radiation is self-absorbed until the peak frequency corresponding to the emission radius (here, 230 GHz at $r \sim 2.5 - 4$ corresponds to $R \sim 5 - 8 r_g$); hence, we add a power-law component $F_\nu \propto \nu^{5/2}$ that we scale to the observed flux at 230 GHz to the existing ADAF spectrum (coming from regions $r \geq 5$). The maximum radius $r_{\text{max}}$ is poorly constrained. As there is no evidence for the presence of a truncated thin disk in the infrared data, we set $r_{\text{max}} = 2 \times 10^5$ so that any contribution from an outer disk truncated at this radius would be negligible (the computation of the outer disk spectrum is performed in Appendix E). This value is consistent with the Bondi radius derived by Asada & Nakamura (2012). For the remaining parameters, we explore a broad range of values, as summarized in Table 1.

For this work, we wish to probe whether an ADAF component could explain the X-ray data, without overestimating the radio-to-optical observations. In Figure 1, we present the spectrum obtained with the parameter values that represent the data best. With the accretion rate dependency on the radius, its value in the innermost regions $r \sim 2.5$ is set to the value inferred from the black hole ring observations (Event Horizon Telescope Collaboration et al. 2019b), where an accretion rate...
in the inner region of $m \sim 2 \times 10^{-5}$ was estimated. The values of $\beta$ and the electron density in the black hole vicinity are compatible with values derived for magnetically arrested disk (MAD; see, e.g., Bisnovatyi-Kogan & Ruzmaikin 1976; Narayan et al. 2003) simulations (Event Horizon Telescope Collaboration et al. 2021). The ADAF component alone is not entirely consistent with the X-ray data; however, its contribution is added to the jet emission to produce the overall SED (see Section 4).

4. Results

With the methods described above we probe whether the total joint model (jet component added to the ADAF component) can explain the global SED. We start by setting the fixed parameters of the jet. Since the synchrotron self-absorption frequency is a critical feature of the observed spectrum (see Section 2), we fix the size of the emission region in order to maximize the self-absorption frequency value, while being consistent with the measured Event Horizon Telescope Collaboration et al. (2019c) flux value, namely, we set $r_{\text{em}} = 5 \times 10^{15}$ cm $\approx 5 r_g$. This corresponds to the radius of the sphere in the moving blob scenario, while for the continuous jet this gives the transverse radius of the cylinder.

For this region, we explore a parameter space starting with varying the magnetic field strength between 10 and 50 G. For each magnetic field we adjust the electron maximum energy and spectral index in order to reproduce the observed cutoff in the optical band, while complementing the ADAF contribution around $10^{16}$ Hz. Once we find the combination between the jet magnetic field strength, the size of the emission region and the injection rate of electrons inside the jet region we determine the Doppler factor $\delta_j \approx 2.3$. This corresponds to a velocity $\beta c = 0.73c$ with an inclination of the jet $\theta = 17^\circ$. With this value of the Doppler factor and the size of the emission region we choose, the length of the cylinder, using the geometry described in Section 2, is $l' \approx 10^{16}$ cm. For the injected proton population, we explore cutoff energies between $10^9$ and $10^{10}$ GeV, and spectral indices between 1.7 and 2.0.
There are less observational constraints on the proton population than for the electrons. We check the ratio of the maximum proton-to-electron energy (see Section 2) and consider models for which the total energy density in particles is lower than or equal to the magnetic energy density. As mentioned in Section 2, the accretion flow could serve as an external target photon field for the jet’s interactions. To assess if the ADAF would make a relevant target field, we compare the energy density of the internal (jet) and external (flow) photon fields in the jet’s frame. To do so we transform the accretion flow radiation field into the jet’s frame, assuming for simplicity that the the flow is seen as a point source behind the jet. This is a rough approximation; however, we only want to estimate the dominant field here.

In Figure 2, we compare the photon spectral number density of the photon fields (for the jet model for which $B = 10$ G, $p_p = 1.7$, $E_{p,\text{max}} = 6 \times 10^9$ GeV, corresponding to the top panel of Figure 3, and the ADAF shown in Figure 1) at two frequencies $10^{11}$ and $10^{18}$ Hz, at which we expect the accretion flow to contribute (see Section 3). Obviously, the internal radio photon field is dominating the external radio photon field. Even at X-ray energies, after only a short time (10 days, over the 2 months of simulated observation) the internal target field contribution is larger than the external one. Therefore, we do not consider the accretion flow as an external target photon field for the jet particles.

The SED is obtained by averaging the light curves over a time corresponding to the observation campaign time, i.e., 2 months. The goodness of the fits (for both the light curve and the SED) is estimated by computing the $p$-value of the $\chi^2$ test for each model. We keep models that have a $p$-value $p < 0.01$.

In Figures 3 and 4, we present four models for which $B = 10$ G. The models have the lowest (highest) maximum proton energy and lowest (highest) proton index possible given the observations and the constraints listed in Section 2. With...
these models we obtain a jet power of $P_j = 2\times 10^{43}$ erg s$^{-1}$ and ratios of magnetic-to-particle energy density of $U_{\text{part}}/U_B = 0.6-1.3$. In Figures 5 and 6 we did the same exploration, and present four models for which $B = 50$ G. We find that with such a high value for the magnetic field strength, it is harder to fit the data, and one has to consider lower proton densities and higher maximum proton energies. For a proton injection spectrum of $p = 2$, we find a good fit only for maximum proton energies $E_{\text{p,max}} \approx 8 \times 10^9$ GeV (see top panel in Figure 6). With these models we obtain a jet power of $P_j \approx 3 \times 10^{44}$ erg s$^{-1}$ and ratios of magnetic energy density to particle energy density of $U_{\text{part}}/U_B \approx 10^{-2}$.

For both $B = 10$ and 50 G, it is easier to obtain a light curve above 350 GeV complying with the observations with a higher value of the maximum proton energy, but since proton synchrotron radiation represents the main contribution to the high-energy spectral bump, the higher the maximum proton energy, the higher the frequency the emission will peak at, and the SED fits get poorer.

The neutrino spectra (single-flavor flux) produced by the source from the models with $B = 10$ G are presented in Figure 7. The predicted flux is low, because the main gamma-ray emission contribution is due to proton synchrotron radiation, which does not produce neutrinos. We compare this value to the sensitivity to a point-like source of high-energy neutrinos with a neutrino flux $\propto E^{-2}$, of the Pierre Auger Observatory (Aab et al. 2019) and the IceCube observatory (Aartsen et al. 2017) at M87’s.

5. Conclusions

We have applied a lepto-hadronic, time-dependent jet model, complemented with an advection-dominated accretion flow to M87’s nuclear emission in a low flux state. We found a range of parameter values that allow to reproduce the multiwavelength data taken in 2017 during the EHT MWL Science Working Group et al. (2021) observation campaign. We investigated two types of jet configuration, namely the moving blob and the continuous jet scenario. For a given geometry, we
Figure 7. Predicted electron and muon neutrino spectra, for the models with $B = 10$ G.

were able to find identical results for both geometries. We focused on a jet emission region of the size similar to the EHT angular resolution at M87’s distance, namely 5 $r_g$. Within this region we estimated a magnetic field strength in the range 5–60 G. The level of flux around the synchrotron self-absorption frequency ($86 < f_{SSA} < 230$ GHz) constrains further the injection parameters of the relativistic electrons in the jet. For a range $10^9 G \leq B \leq 50 G$ we found that the electrons spectral index is limited to $p_e \approx 1.80$–1.85, in order to reproduce the radio-to-optical part of the SED. For the same reason, the maximum energy for the electron distribution is found to be $E_{\text{max},e} \lesssim 5$ GeV. Concerning the high-energy emission, we have found parameter values that fit the data for the whole range of magnetic field strengths considered. However, it is worth pointing out that the proton maximum energy and spectral index ranges are dependent on the value for the magnetic field strength. For $B = 10$ G we found that when $p_p \approx 1.7$ (minimum proton spectral index) the proton maximum energy is in the range $6 \times 10^9$ GeV $\lesssim E_{\text{max},p} \lesssim 1 \times 10^{10}$ GeV while for $p_p \approx 1.85$ (maximum proton spectral index for these parameter values) it is in the range $7 \times 10^9$ GeV $\lesssim E_{\text{max},p} \lesssim 1 \times 10^{10}$ GeV. For $B = 50$ G when $p_p \approx 1.7$ (minimum proton spectral index) the proton maximum energy is in the range $7 \times 10^9$ GeV $\lesssim E_{\text{max},p} \lesssim 1 \times 10^{10}$ GeV while for $p_p \approx 2.00$ (maximum proton spectral index for these parameter values) it is in the range $8 \times 10^9$ GeV $\lesssim E_{\text{max},p} \lesssim 1 \times 10^{10}$ GeV. This required increase in the maximum proton energy makes it harder to fit the gamma-ray part of the SED above $10^{25}$ Hz. Combining the jet’s emission with the ADAF allows us to reproduce at the same time the apparent cutoff in the optical band, and the power-law-like flux component at X-rays energies. Unlike previous works (e.g., Nemmen et al. 2014; Feng & Wu 2017), we found a configuration where both the jet and the ADAF have a distinct contribution in the SED. Beyond the scope of this paper, as mentioned in the introduction, would be the estimation of the contribution to the cosmic-ray flux from M87. Indeed, with protons accelerated up to $10^{10}$ GeV and a jet power of $10^{43–44}$ erg s$^{-1}$, M87 could contribute to the detected high-energy cosmic-ray flux on Earth (Protheroe et al. 2003, see Gaissler 2013 concerning the power requirements of cosmic-ray sources).

In this work we have considered only one-zone models for the jet emission. In the framework of structured jet models, multi-zone scenarios have been invoked to explain M87’s SED (e.g., Sol et al. 1989; Georganopoulos & Kazanas 2003). In particular, Ghisellini et al. (2005) developed a leptonic scenario in which a fast inner jet is embedded in a slower outer sheath. Here, the beaming pattern related to the boosting of one layer into the other could explain the high-energy part of the SED. They applied it to M87’s SED (Tavecchio & Ghisellini 2008). A transverse structure in jets is further supported by observations of limb brightening (for M87 see Kovalev 2008, for radio-galaxies and blazars see Giroletti et al. 2004; Giovannini et al. 1999). However, with twice as many parameters as for one-zone jet models, such as the one we considered, it is difficult to constrain two-zone scenarios to date.

M.B. received funding from the European Union’s Horizon 2020 research and innovation program under the Marie Sklodowska-Curie grant agreement No. 847476. The views and opinions expressed herein do not necessarily reflect those of the European Commission. M.B. wishes to thank Paolo Da Vela for the fruitful discussions and insightful comments on this paper.

L.M. acknowledges support from the DFG within the Collaborative Research Center SFB1491 “Cosmic Interacting Matters—From Source to Signal”.

This research was funded in part by the Austrian Science Fund (FWF; grant No. I 4144-N27). For the purpose of open access, the author has applied a CC BY public copyright licence to any Author Accepted Manuscript version arising from this submission.
**Software:** This work benefited from the following software: NumPy (van der Walt et al. 2011), Matplotlib (Hunter 2007), pandas (Wes McKinney 2010; Jeff et al. 2022), jupyter notebooks (Pérez & Granger 2007).

### Appendix A

**Jet’s Magnetic Field: b(p) Coefficient**

Kino et al. (2014) derived a relation between the jet’s magnetic field and the observable quantities:

\[
B = b(p) \left( \frac{\nu_{SSA, obs}}{1 \text{GHz}} \right)^{5/4} \left( \frac{\theta_{obs}}{1 \text{mas}} \right)^{3} \left( \frac{S_{SSA, obs}}{1 \text{Jy}} \right)^{-2} \left( \frac{\delta}{1 + z} \right)^{4} \text{G}.
\]

Here, \( b(p) \) is defined as:

\[
b(p) = 5.52 \times 10^{5/7} \left( 3X_{c}c_{2}(p) / (2\pi X_{c}(p)) \right)^{2},
\]

where

\[
X_{1} = \frac{\sqrt{3}}{8\pi m_{e}} \left( \frac{3e}{2\pi m_{e} c^{2}} \right)^{1/2}, \quad c_{1}(p) = \frac{3p + 2}{12} \Gamma \left( \frac{3p + 22}{12} \right),
\]

\[
X_{2} = \frac{\sqrt{3}}{8\pi m_{e}} \left( \frac{3e}{2\pi m_{e} c^{2}} \right)^{(p-1)/2}, \quad c_{2}(p) = \frac{3p + 19}{12} \Gamma \left( \frac{3p + 1}{12} \right)
\]

\[\times \Gamma \left( \frac{p + 5}{4} \right) / \Gamma \left( \frac{p + 7}{4} \right) / (p + 1).
\]

### Appendix B

**ADAF Self-similar Solutions**

The one-zone, height-integrated, self-similar solutions of the thin disk equations were derived by Narayan & Yi (1995) to describe the hot plasma. The solutions and their expression using the relevant scaled quantities are given by

\[
v_{R} \approx \frac{\alpha}{2} v_{K} \approx 1.06 \times 10^{9} \alpha_{-1} r^{-1/2} \text{ cm s}^{-1},
\]

\[
c_{s} \approx \frac{1}{2} v_{K} \approx 1.06 \times 10^{10} r^{-1/2} \text{ cm s}^{-1},
\]

\[
\rho \approx \frac{p_{m}}{4\pi R H v_{K}} \approx 2.66 \times 10^{-15} m_{BH,9}^{1} m_{-3}^{-1} \alpha_{-1}^{-1} r^{-3/2} \text{ g cm}^{-3},
\]

\[
B \approx \sqrt{\frac{8\pi \rho \epsilon_{\gamma}^{2}(1 - \beta)}{c_{s}^{2}}}
\]

\[
\approx 2.75 \times 10^{3} m_{BH,9}^{1/2} m_{-3}^{1/2} \alpha_{-1}^{-1/2} (1 - \beta)^{1/2} r^{-5/4} \text{ G},
\]

\[
\tau_{\gamma} \approx \rho_{m} \sigma_{T} R \approx 0.313 m_{-3}^{-1} \alpha_{-1}^{-1} r^{-1/2},
\]

where \( v_{K} = \sqrt{GM_{BH}/R} \) is the Keplerian velocity, \( v_{K} \) is the radial velocity, and \( c_{s} \) is the isothermal sound’s speed. Here, \( \alpha \) is the viscosity parameter introduced by Shakura & Sunyaev (1973), \( \beta \) is the ratio between the gas \( p_{g} \) and the total pressure \( p = \rho c_{s}^{2} = \rho_{m} + p_{g} \) with \( p_{m} = B^{2}/8\pi \), where \( B \) is the isotropically tangled magnetic field. The Thomson optical depth is denoted by \( \tau_{\gamma} \).

### Appendix C

**ADAF Cooling Mechanisms**

#### C.1. Synchrotron Radiation and Bremsstrahlung

We assume that the plasma electrons follow a relativistic Maxwellian distribution

\[
N_{e}(\gamma_{e}, \theta_{e}) = n_{e} \left( \frac{\beta_{e} \gamma_{e}}{\theta_{e}} \right)^{3} \frac{\exp(-\gamma_{e}/\theta_{e})}{\theta_{e} K_{2}(1/\theta_{e})},
\]

where \( n_{e} \approx n_{p} \) is the electrons number density, \( \beta_{e} \) and \( \gamma_{e} \) are the relative velocity and the Lorentz factor of the thermal electrons, respectively, and \( K_{2}(x) \) is the 2nd order modified Bessel function.

For synchrotron radiation from thermal electrons and bremsstrahlung, we use the fitting formula derived by Narayan & Yi (1995). The synchrotron emissivity is given by

\[
j_{\nu, syn} = 4.43 \times 10^{-30} \frac{4\pi n_{e} \nu}{K_{2}(1/\theta_{e})} \left( \frac{4\pi m_{e} c \nu}{3 e B \theta_{e}^{2}} \right) \text{ erg cm}^{-3} \text{ s}^{-1} \text{ Hz}^{-1},
\]

where \( I'(x) \) is defined in Narayan & Yi (1995):

\[
I'(x) = \frac{4.0505}{x^{1/6}} \left[ 1 + \frac{0.4}{x^{1/4}} + \frac{0.5316}{x^{1/2}} \right] \exp(-1.8899 x^{1/3}).
\]

The Bremsstrahlung cooling rate is given by the sum of the rates from electron–electron and ion–electron interactions (Svensson 1982; Stepney & Guilbert 1983):

\[
q_{br} = q_{ee} + q_{ie}.
\]

The ion–electron and electron–electron bremsstrahlung cooling rates are respectively given by Svensson (1982); Stepney & Guilbert (1983):

\[
q_{ee} = 1.48 \times 10^{-22} n_{e}^{2} F_{e}(\theta_{e}) \text{ erg cm}^{-3} \text{ s}^{-1},
\]

\[
q_{ie} = \begin{cases} 
2.56 \times 10^{-22} n_{e}^{2} \theta_{e}^{1/2} (1 + 1.16 \theta_{e}^{1/2} - 1.25 \theta_{e}^{1/2}) & \text{if } \theta_{e} < 1 \\
3.40 \times 10^{-22} n_{e}^{2} \theta_{e} [\ln(1.123 \theta_{e}) + 1.28] & \text{if } \theta_{e} > 1
\end{cases}
\]

where

\[
F_{e}(\theta_{e}) = \begin{cases} 
4 \left( \frac{2\theta_{e}}{\pi} \right)^{0.5} (1 + 1.1781 \theta_{e}^{3/4}) & \text{if } \theta_{e} < 1 \\
9 \theta_{e}^{5/2} [\ln(1.123 \theta_{e}) + 1.5] & \text{if } \theta_{e} > 1
\end{cases}
\]

Assuming a Gaunt factor equal to unity, we approximate the emissivity due to bremsstrahlung to be

\[
j_{\nu, br} \approx q_{br} \frac{h}{k_{B} T_{e}} \exp \left(-\frac{h\nu}{k_{B} T_{e}} \right).
\]

#### C.2. Compton Cooling

In order to take into account the inverse-Compton scattering in the cooling mechanism we use the formulation derived by Esin et al. (1996) to compute the cooling rate. Assuming the Comptonization is enhancing the initial energy of the seed photons, we can define the energy enhancement factor \( \eta(\nu) \)
such that
\[ \eta = \exp[s(A - 1)][1 - P(j_m + 1, As)] + \eta_{\text{max}} P(j_m + 1, s), \]  
where \( P(a, x) \) is the regularized lower incomplete gamma function and
\[ A = 1 + 4\theta_c + 16\theta_c^2, \quad s = \Gamma + \tau^2. \]
\[ \eta_{\text{max}} = \frac{3k_B T_e}{h\nu}, \quad j_m = \frac{\ln \eta_{\text{max}}}{\ln A}. \]
Finally, the total cooling rate of the electrons is given by
\[ q^e^- = \frac{1}{H} \int \text{d}n(\nu)F_{\nu,0}, \]  
where \( F_{\nu,0} \) is described in the main text, Equation (2).

**Appendix D**

**ADAF Heating Rates**

The electrons are heated in two different ways in the plasma, namely, they can be directly heated by a fraction \( \delta_c \) of the viscous dissipated energy, and they can also be heated through Coulomb collisions with the ions.

The viscous energy dissipation rate per unit volume \( q^{\text{visc}} \) is given in (Narayan & Yi 1995) as
\[ q^{\text{visc}} = \frac{3\epsilon^e \rho v^2 e_i^2}{2R} = 0.08e \rho m_{\text{BH}}^2 m_{\text{e}} \epsilon^e^{-4} \text{ erg cm}^{-3} \text{ s}^{-1}, \]  
where \( \epsilon^e = (5/3 - \gamma)/\gamma - 1, \) with \( \gamma = (32 - 24\beta - 3\beta^2)/(24 - 21\beta). \)

For the Coulomb interaction heating rate per unit volume \( q^{\text{Coul}} \) from Stepney & Guilbert (1983), assuming \( n_e \approx n_p, \) we use the approximation from Mahadevan (1997):
\[ q^{\text{Coul}} = 5.61 \times 10^{-32} \frac{\rho^2 (T_e - T_i)}{K_e (1/\theta_e)} \left( \frac{\theta_e \theta_i}{\theta_e + \theta_i} \right)^{1/2} \times \frac{2(\theta_e + \theta_i)^2 + 1 + 2(\theta_e + \theta_i)}{(\theta_e + \theta_i)} e^{-1/\theta_e} \text{ erg cm}^{-3} \text{ s}^{-1}. \]  

The total heating rate is then given by
\[ q^e = q^{\text{visc}} + \delta_c q^{\text{Coul}}. \]

**Appendix E**

**Truncated Thin Disk**

For completeness, we compute the spectrum from an outer disk, such that the inner radius of the disk is equal to the outer radius of the accretion flow \( r_\text{in} = r_{\text{max}}. \) The emission is characterized by the sum of blackbody spectra with temperature
\[ T_{\text{disk}}(R) = \left( \frac{G M_{\text{BH}} \dot{M}}{8\pi R^3 \sigma} \right) \left[ 1 - \left( \frac{R_{\text{tr}}}{R} \right)^{1/2} \right]^{1/4}. \]  

The emission is then given by
\[ F_{\nu,\text{disk}} = \frac{4\pi h \cos \theta \nu^3}{c^2 D^2} \int_{R_{\text{tr}}}^{R_{\text{max, disk}}} \frac{R dR}{e^{h\nu/kT_{\text{disk}}(R)} - 1}, \]  
where we have set an outer radius of \( \sim 3 \times 10^6 \) but this parameter has a poor influence on the spectrum, given the large truncation radius.

**ORCID iDs**

Margot Bougheliba @ https://orcid.org/0000-0003-1046-1647
Anita Reimer @ https://orcid.org/0000-0001-8604-7077
Lucas Merten @ https://orcid.org/0000-0003-1332-9895

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