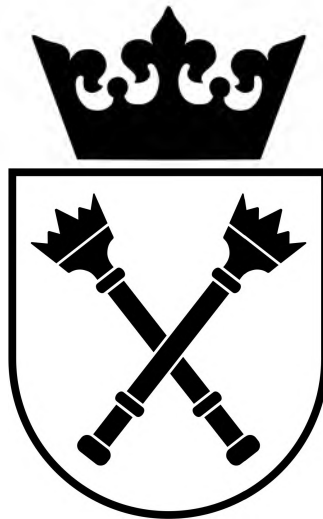


Physics of Evolving Black Holes and Their Relativistic Jets: From Collapsing Massive Stars to Active Galactic Nuclei

Ph.D. Dissertation
by
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Black holes (BHs) are associated with some of the most spectacular astrophysical phenomena. Stellar-mass black holes are the central engines of γ -ray bursts and X-ray binaries, whereas supermassive black holes (SMBHs), with masses exceeding millions of solar masses, drive the activity in galactic nuclei (AGN). The energy of the matter accreting onto a BH is dissipated through viscous interactions, which heat the gas and amplify the magnetic field. This process results in intense high-energy emission and drives the launch of disk winds and outflows. In some cases, spectacular relativistic jets are launched by the disk-supported magnetic field, at the expense of the black hole's rotational energy.

Interestingly, a BH can be described by only three parameters: mass, spin, and electric charge. In the astrophysical context, the charge is expected to quickly neutralize, leaving mass and spin as the only parameters describing black holes, regardless of their origin, evolution, or accretion history. Despite this apparent simplicity, BHs present a multifaceted appearance as astrophysical sources, playing an active role in shaping their surroundings. Still, many questions regarding this role remain unanswered. This Ph.D. dissertation consists of four scientific articles that delve into the physics of black holes and their relativistic jets across different astrophysical systems.

In the first project, we studied the spin and mass evolution of a black hole just after its formation in the collapse of a hyper-massive star. We aimed to determine how the final values of these parameters depend on the initial BH's spin, the angular momentum in the accreting matter, and the strength of the magnetic field. Collapsars are considered to be progenitors of long gamma-ray bursts. To conduct this research, we employed a modified version of the High Accuracy Relativistic Magnetohydrodynamics numerical code, which operates within the framework of general relativistic magnetohydrodynamics.

In the following project, we focused on the relativistic jets of AGN, in particular blazars – a class of active galaxies, whose broadband spectra are dominated by the Doppler-boosted jet emission. We were interested in the monochromatic variability features appearing in the S5 07165+714 optical light curve over the last ~ 20 years. We proposed that these features may originate from the gravitational lensing events, in which the emission of consecutive

jet segments is lensed by a heavy binary system, located within the blazar's host galaxy. Depending on the source-lens distance, the binary may contain an intermediate-mass black hole and a much less massive companion. This offers a fresh perspective on modeling of blazars variability and presents an exciting opportunity for the search and identification of intermediate-mass black hole systems.

The third project revolved around the broadband emission of Compact Symmetric Objects, a subset of AGN known for their young radio structures, which may provide valuable insights into the initial evolutionary stages of radio galaxies. In our analysis we were interested in the source of a high energy emission in these systems. Specifically, for three objects distinguished by the fluorescence iron K_{α} lines in their heavily absorbed X-ray spectra (J1407+2827, J1511+0518, and J2022+6137), we demonstrated that the high-energy emission can originate from compact radio lobes still embedded within dusty tori.

The last project included in the thesis focused on the magnetization of relativistic current-carrying jets with radial velocity shears at large distances from their launching sites. Through analytical and numerical calculations, we demonstrated that jets at these distances – assuming magneto-hydrostatic equilibrium and smooth pressure profiles – are inevitably dominated by the bulk kinetic energy of matter, rather than by the Poynting flux from the jet's toroidal magnetic field. Although this finding might be somewhat unexpected, it complements earlier studies on jet evolution. These studies indicate that collimating and accelerating magnetohydrodynamic jets, confined by ambient pressure, remain dominated by Poynting flux until they reach logarithmically large distances from the central black hole.

The dissertation begins with a brief introduction and concludes with a discussion of future prospects.

**Fizyka ewoluujących czarnych dziur i ich relatywistycznych dżetów:
od zapadających się masywnych gwiazd do aktywnych jąder galaktyk.**

Dominika Ł. Król

Czarne dziury powiązane są z najbardziej spektakularnymi zjawiskami obserwowanymi we Wszechświecie. Te o masach gwiazdowych stoją za rozbłyskami gamma oraz aktywnością rentgenowskich układów podwójnych, podczas gdy supermasywne czarne dziury, o masach przekraczających miliony mas Słońca, napędzają aktywność jąder galaktyk (ang: Active Galactic Nuclei, w skrócie AGN). Z uwagi na lepkość, energia materii akreującej na czarną dziurę jest dysypowana, co zwiększa energię wewnętrzną gazu oraz wzmacnia pole magnetyczne. Zjawiska te skutkują silną emisją na bardzo wysokich częstotliwościach promieniowania elektro-magnetycznego oraz mogą prowadzić do powstania wiatrów i wypływów z dysku akrecyjnego. W niektórych przypadkach dochodzi do wyrzucenia spektakularnych relatywistycznych strug (dżetów) które powstają kosztem energii rotacji czarnej dziury.

Co ciekawe, czarną dziurę można opisać za pomocą jedynie trzech parametrów. Są to masa, spin i ładunek elektryczny. W kontekście astrofizycznym ładunek powinien jednak zostać szybko zneutralizowany z uwagi na wysoką przewodność plazmy kosmicznej, czyniąc masę i spin jedynymi parametrami opisującymi czarną dziurę, niezależnie od jej pochodzenia, ewolucji i historii akrecji. Pomimo tej pozornej prostoty, czarne dziury występują w szerokim wachlarzu obiektów astrofizycznych, odgrywając aktywną rolę w kształtowaniu swojego otoczenia. Wiele pytań dotyczących szczegółów tego oddziaływania i roli aktywności pozostaje otwartych. Niniejsza praca doktorska składa się z czterech opublikowanych, powiązanych tematycznie artykułów naukowych, które zgłębiają fizykę czarnych dziur i ich relatywistycznych dżetów w różnych systemach astrofizycznych.

W pierwszym projekcie analizowaliśmy ewolucję spinu i masy czarnej dziury tuż po jej powstaniu w wyniku kolapsu supermasywnej gwiazdy. Naszym celem było określenie, w jaki sposób ostateczne wartości tych parametrów zależą od początkowego spinu czarnej dziury, momentu pędu akreującej materii i siły pola magnetycznego. Uważa się, że tego typu systemy (ang: collapsars) są źródłami długich błysków gamma. W analizie wykorzystaliśmy zmodyfikowaną wersję kodu numerycznego High Accuracy Relativistic Magnetohydrodynamics rozwiązującego równania magnetohydrodynamiczne w ramach ogólnej teorii względności.

W kolejnym projekcie skupiliśmy się na relatywistycznych dżetach blazarów (ang: blazars), klasy AGNów których emisja jest zdominowana właśnie przez relatywistycznie wzmocnioną emisję dżetu. Zainteresowały nas monochromatyczne modulacje pojawiające się w optycznej krzywej zmian blasku blazara S5 07165+714 na przestrzeni ostatnich ~ 20 lat. Wyszliśmy hipotezę, że modulacje te mogą być wynikiem soczewkowania grawitacyjnego, w którym emisja kolejnych segmentów dżetu jest soczewkowana przez ciężki układ podwójny znajdujący się w galaktyce macierzystej blazara. W zależności od odległości pomiędzy źródłem a soczewką, układ ten zawierać może czarną dziurę o masie pośredniej (ang: Intermediate Mass Black Hole). Wyniki analizy oferują nowe spojrzenie na modelowanie zmienności blazarów oraz wskazują potencjalną nową metodę identyfikacji układów podwójnych zawierających czarne dziury o masach pośrednich.

Trzeci projekt dotyczył wielozakresowej emisji zwartych obiektów symetrycznych (ang: Compact Symmetric Objects). Jest to klasa AGNów charakteryzujących się młodymi strukturami radiowymi, które są cennym źródłem informacji o początkowych etapach ewolucji radiogalaktyk. Przedmiotem naszej analizy była identyfikacja źródła emisji rentgenowskiej w tych układach. Wykazaliśmy, że w przypadku trzech obiektów wyróżniających się występowaniem fluorescencyjnej linii żelaza K_{α} w ich silnie zaabsorbowanych widmach (J1407+2827, J1511+0518 i J2022+6137), wysokoenergetyczna emisja może pochodzić ze zwartych płatów (lobów) radiowych znajdujących się wewnątrz torusa pyłowego.

Ostatni projekt zawarty w tej pracy doktorskiej dotyczy magnetyzacji relatywistycznych dżetów z radialnym ścinaniem prędkości. Poprzez rachunki matematyczne oraz obliczenia numeryczne pokazaliśmy, że na dużych odległościach od centrum galaktyki, energia kinetyczna materii przewyższa strumień Poyntinga dżetów w równowadze magneto-hydrostatycznej, opisywanych gładkim profilem ciśnienia. Ten zaskakujący wynik jest interesujący w zestawieniu z prezentowanymi w literaturze rachunkami pokazującymi, że dżety kolimowane i przyśpieszane magnetohydrodynamicznie, przy udziale ciśnienia otoczenia, pozostają zdominowane przez strumień Poyntinga, nawet na logarytmicznie dużych odległościach od centralnej czarnej dziury.

*Dla mamy.
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Contents

Abstract	i
Streszczenie	iii
Acknowledgements	vii
List of Figures	ix
List of Tables	xiii
List of Papers	xvii
List of Formulas	xix
1 Introduction	1
1.1 Dissertation outline	11
2 Accretion-induced Black Hole Spin-up Revised by Numerical General Relativistic MHD	15
3 Possible Gravitational Microlensing Events in the Optical Lightcurve of Active Galaxy S5 0716+714	37
4 On the Origin of the X-ray Emission in Heavily Obscured Compact Radio Sources	47
5 Magnetization of Relativistic Current-carrying Jets with Radial Velocity Shear	65
6 Summary and Future Prospects	75

List of Figures

1.1	Critical surfaces of a Kerr black hole.	2
1.2	GRMHD simulations of the jet launch, showing energy extraction from the rotating black hole	3
1.3	Schematic diagram of a micro-lensing event.	5
1.4	Collection of detected stellar-mass black hole masses	6
1.5	Composition of M87 observations across different scales and wavelengths.	7
1.6	Cartoon of the black holes density function.	8
1.7	$M_{\text{BH}}-\sigma_B$ and $M_{\text{BH}}-L_B$ relations.	10
2.1	Density distributions of the innermost regions of the accretion flow at the beginning of numerical simulations	20
2.2	The mass flux through the inner boundary as a function of time.	21
2.3	The black hole spin as a function of time.	21
2.4	The black hole mass as a function of time.	22
2.5	The profiles of density, specific angular momentum, and Mach number during the simulations without a magnetic field.	22
2.6	The profiles of density, specific angular momentum, and Mach number during the simulations without a magnetic field.	23
2.7	The profiles of final density, specific angular momentum, and Mach number at the end of the simulations without a magnetic field.	23
2.8	The profiles of density, specific angular momentum, and Mach number during the simulations without a magnetic field.	24
2.9	Color maps of the maximal and final black hole spin, final black hole mass, and the maximal accretion rate as a function of initial model parameters.	25
2.10	Shock front positions as a function of time.	26
2.11	Mass accretion rate, black hole mass, and spin as a function of time during the simulations with a magnetic field.	27
2.12	Mass accretion rate, black hole mass, and spin as a function of time during the simulations with a magnetic field.	27

2.13	Mass accretion rate, black hole mass, and spin as a function of time during the simulations with a magnetic field.	27
2.14	The profiles of specific angular momentum, Mach number, and density with over-plotted contours of magnetic field during the simulations with a weak magnetic field.	28
2.15	The profiles of specific angular momentum, Mach number, and density with over-plotted contours of the magnetic field during the simulations with a strong magnetic field.	29
2.16	The profiles of specific angular momentum, Mach number, and density with over-plotted contours of the magnetic field during the simulations with a strong magnetic field.	30
2.17	The profiles of specific angular momentum, Mach number, and density with over-plotted contours of the magnetic field during the simulations with a strong magnetic field.	31
2.18	Comparison of density profiles obtained in simulations with disk profiles discussed in the literature.	33
3.1	S5 0716+714 light-curves observed by TESS.	39
3.2	Close-up of three events (features) in a light curve, superimposed on each other.	40
3.3	Light curves with fitted lens models.	42
3.4	The magnification pattern along with the source trajectory and the resulting source lightcurve.	44
4.1	Left panel: The merged Chandra ACIS-S image of three J1511+0518 observations. Right panel: brightness profile of the merged image, along with the simulated averaged PSF.	51
4.2	The broadband SED and the best-fit models for J1407+2827.	55
4.3	The broadband SED and the best-fit models for J1511+0518.	56
4.4	The broadband SED and the best-fit models for J2022+6137.	57
4.5	The broadband SED and an example of the model requiring X-ray coronal emission for J2022+6137.	58
4.6	Comparison of the high-energy slope of the injection electron distribution with the radio linear size LS, and the source-intrinsic absorbing column density for the CSOs sample.	58
4.7	Left panel: schematic view of an evolved radio-loud AGN. Right panel: schematic view of a young radio source, where the X-ray continuum emission is produced within compact radio lobes.	59

5.1	Three exemplary profiles of parameters describing analysed jet model.	69
5.2	Jet magnetization as a function of boundary condition's value.	70
5.3	Jet magnetization and plasma parameter profiles for different models.	70
5.4	Cumulative jet magnetization and its share of the total energy flux for some jet models.	71
5.5	Magnetization of the jet as a function of model parameters.	72
5.6	The pressure discontinuity as a function of model parameters.	72

List of Tables

2.1	Summary of calculated HARM models.	20
2.2	Velocities of the shock fronts	26
3.1	Best-fit model parameters.	43
4.1	Chandra observations of J1511+0518.	49
4.2	Parameters of the Expanding Radio Lobe model.	52
4.3	Parameter values assumed in our modeling based on the radio and IR data.	53
4.4	Adopted range of the ambient medium density.	53
4.5	Modeling Results.	55
4.6	Fermi-LAT upper limits.	60
4.7	Details of the SMA observations.	61
4.8	List of references to the archival radio, IR, and Optical/UV data.	61

List of Papers

This thesis include following published peer-reviewed papers:

- Król, D. Ł., Sobolewska, M., Stawarz, Ł., Siemiginowska, A., Migliori, G., Principe, G., Gurwell, M.
On the Origin of the X-Ray Emission in Heavily Obscured Compact Radio Sources
2024, the Astrophysical Journal, 966, 201.
- Król, D. Ł., Stawarz, Ł., Krzesiński J. and Cheung, C. C.
Possible Gravitational Microlensing Events in the Optical Lightcurve of Active Galaxy S5 0716+714
2023, the Astrophysical Journal, 943, 171.
- Król, D. Ł., Stawarz, Ł., Begelman, M. C., Martí, J.-M., Perucho, M., Petrenko, B. A. *Magnetization of Relativistic Current-carrying Jets with Radial Velocity Shear*
2022, the Astrophysical Journal, 929, 181.
- Król, D. Ł., Janiuk, A.
Accretion-induced Black Hole Spin-up Revised by Numerical General Relativistic MHD
2021, the Astrophysical Journal, 912, 132.

Other publications not included in the thesis:

- Janiuk, A., Shahamat Dehsorkh, N., **Król, D.**
Self-gravitating collapsing star and black hole spin-up in long gamma ray bursts
2023, Astronomy & Astrophysics, 677, 23,
- Balasubramaniam, K., Stawarz, Ł., Cheung, C. C., Sobolewska, M., Marchenko, V., Thimmappa, R., **Król, D. Ł.**, Migliori, G., Siemiginowska, A.
X-Ray Emission of the γ -ray-loud Young Radio Galaxy NGC 3894
2021, the Astrophysical Journal, 922, 84.

- Król, D. Ł., Marchenko, V., Ostrowski, M., Stawarz, Ł.
An Analysis of Soft X-Ray Structures at Kiloparsec Distances from the Active Nucleus of Centaurus A Galaxy
2020, the Astrophysical Journal, 903, 107.

List of Formulas Referred to in Introduction and Summary

Gravitational radius	$r_g = \frac{GM_{\text{BH}}}{c^2}$
Dimensionless black hole spin	$a = \frac{J_{\text{BH}}}{cM_{\text{BH}}r_g}$
Schwarzschild metric	$ds_s^2 = -\left(1 - \frac{2r_g}{r}\right)c^2dt^2 + \left(1 - \frac{2r_g}{r}\right)^{-1}dr^2 + r^2d\Omega^2$
Two-sphere metric	$d\Omega^2 = d\theta^2 + \sin^2\theta d\phi^2$
Kerr metric	$ds^2 = -\left(1 - \frac{2r_g r}{\Sigma}\right)c^2dt^2 + \frac{\Sigma}{\Delta}dr^2 - \frac{4r_g a}{\Sigma}r \sin^2\theta cdt d\phi^2 + \left(r^2 + a^2 r_g^2 + \frac{2a^2 r_g^3}{\Sigma}r \sin^2\theta\right) \sin^2\theta d\phi^2$
Outer event horizon of Kerr BH	$r_h = r_g + \sqrt{r_g^2 - r_g^2 a^2}$
Static limit of Kerr BH	$r_e = r_g + \sqrt{r_g^2 - r_g^2 a^2 \cos^2\theta}$
Radius of gravitational influence of the SMBHs	$r_i = \frac{2GM_{\text{BH}}}{\sigma^2} \simeq 10^{-2} \left(\frac{M_{\text{BH}}}{10^8 M_\odot}\right) \left(\frac{\sigma_*}{200 \text{ km/s}}\right)^{-2} \text{ kpc}$
Eddington accretion rate limit	$\dot{M}_{\text{Edd}} = \frac{4\pi GM_{\text{BH}} m_p}{\epsilon c \sigma_T} \simeq 2.2 \epsilon^{-1} \left(\frac{M_{\text{BH}}}{10^9 M_\odot}\right) M_\odot \text{ yr}^{-1}$
$M_{\text{BH}}-\sigma$ relation	$\frac{M_{\text{BH}}}{10^9 M_\odot} = 0.542 \times \left(\frac{L_B}{10^{11} L_\odot}\right)^{1.21}$
$M_{\text{BH}}-L$ relation	$\frac{M_{\text{BH}}}{10^9 M_\odot} = 0.309 \times \left(\frac{\sigma_B}{200 \text{ km s}^{-1}}\right)^{4.38}$

Chapter 1

Introduction

Over a century ago, Karl Schwarzschild found an exact solution to Einstein's equations describing the space-time geometry outside a non-rotating spherical mass (Schwarzschild, 1916). Approximately 40 years later, Roy Kerr and Ezra Newman found solutions describing, respectively, a neutral (Kerr, 1963) and a charged (Newman et al., 1965) rotating spherical mass. These solutions, widely referred to as the Schwarzschild, Kerr and Kerr-Newman metrics, became the foundations of the mathematical description of black holes (BHs). Initially, BHs were considered as purely mathematical concepts. This perception quickly changed with astronomical observations made in the 1950s and 1960s that led to the correct identification of the first radio galaxies (e.g., Cygnus A; Baade and Minkowski, 1954) and quasars (e.g., 3C 273; Rees, 1966) as extragalactic sources. With accurate distance estimates, it became clear that the observed energy fluxes necessitated a new, extremely efficient mechanism of energy extraction. In the late 1970s, the Blandford-Znajek mechanism (BZ; Blandford and Znajek, 1977) was introduced as a solution. In the proposed scenario, the source of energy is the angular momentum of a spinning Kerr BH, efficiently extracted in the presence of a magnetic field supported by the accretion flow. This concept solidified BHs as real astrophysical phenomena and identified them as the central engines in various accretion-powered sources.

In a broader context, accretion of a magnetized plasma onto compact objects can be observed in many distinct types and classes of astrophysical systems, not necessarily containing BHs. These accretion processes often lead to the formation of outflows and disk winds. In systems such as close binary stars and young stellar objects, the energy from accreting matter is dissipated through viscous interactions, which heat the gas and amplify the magnetic field, leading to the formation of collimated jets characterized by sub-relativistic bulk motions. However, the launching of luminous *relativistic* jets is an almost unique and spectacular manifestation of the BH activity. These

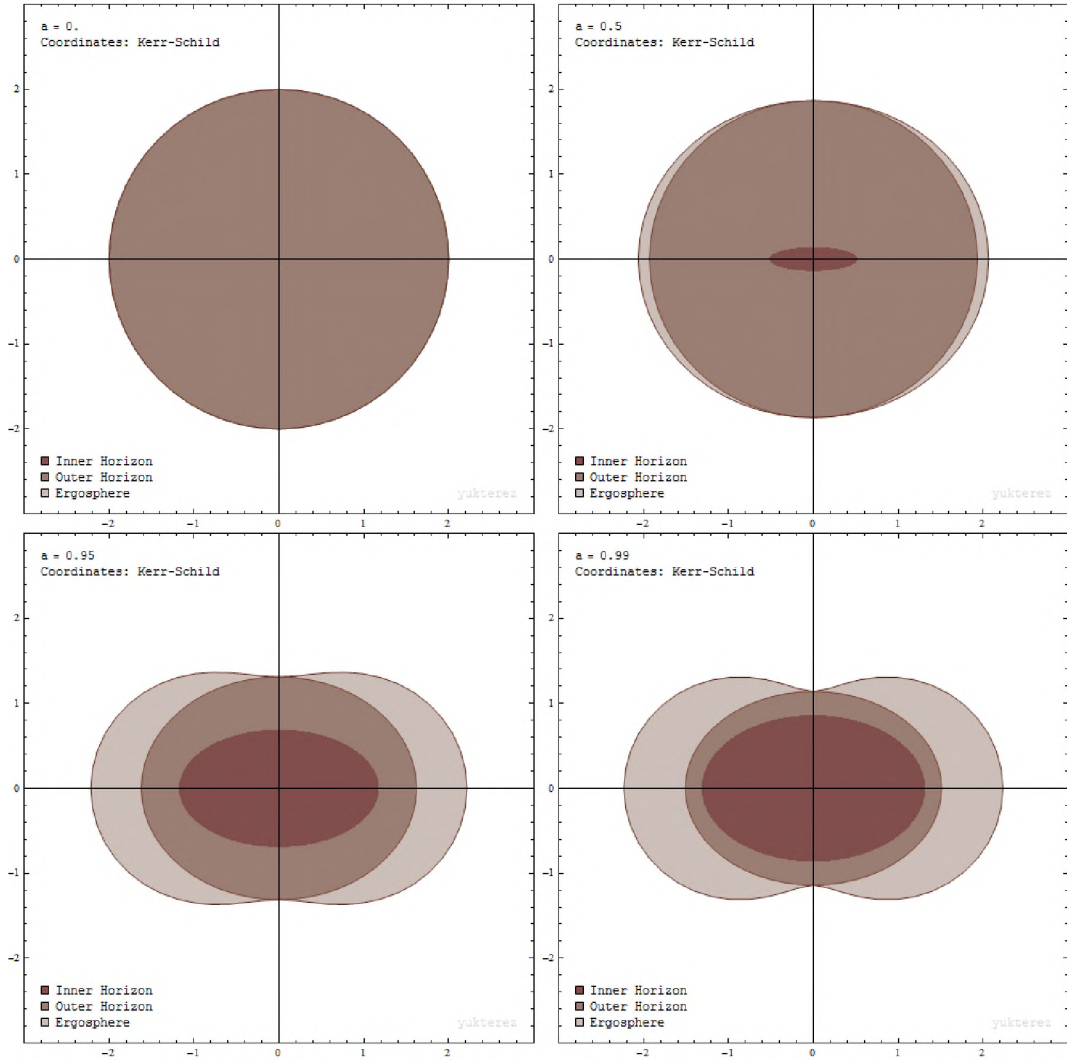


FIGURE 1.1: The geometrical structure of the Kerr's BH critical surfaces in Kerr-Schild coordinates for different dimensionless spin values, a .

Source: https://en.wikipedia.org/wiki/Rotating_black_hole

relativistic jets, when launched from supermassive black holes (SMBHs) in active galactic nuclei (AGN), can fuel giant radio lobes and emit relativistically beamed radiation, as observed in quasars and blazars. Relativistic jets can be produced by systems with a wide range of BH masses, from billions to several Solar masses in AGN and X-ray binaries (XRBs) or gamma-ray bursts (GRBs), respectively.

Stellar-mass BHs are the final stage in the evolution of the most massive stars. Interestingly, regardless of the complexity of the progenitor star's evolutionary history, or of how many subsequent mergers and accretion events the BH has undergone, according to the no-hair theorem (Gürlebeck, 2015),

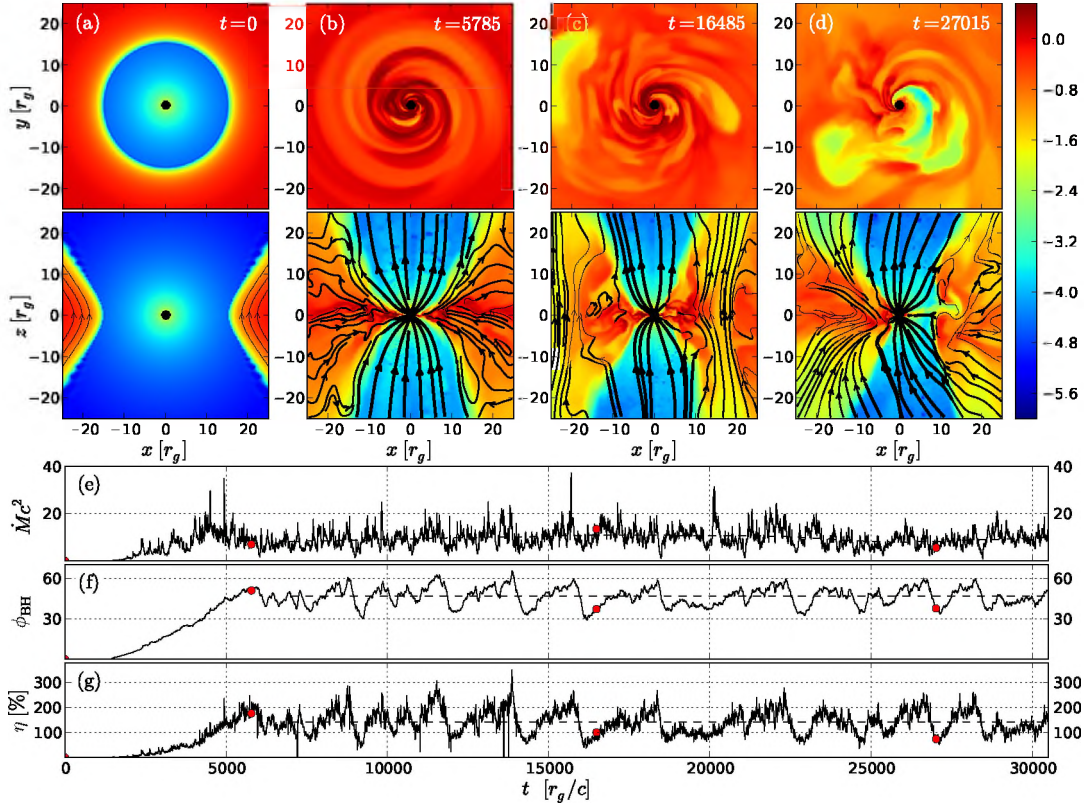


FIGURE 1.2: GRMHD simulations results from Tchekhovskoy, Narayan, and McKinney, 2011. Panels a-d) represent fluid-frame rest-mass density profile snapshots at different times of simulation, field lines in the image plane in equatorial (upper row) and meridional (lower panel) plane. The time evolution of the accretion rate and magnetic flux at the BH horizon are shown at panels e-f). Panel f) shows jet energy divided by the accretion rate, its high value—exceeding 100%—indicating energy extraction from the BH.

Source: Tchekhovskoy, Narayan, and McKinney, 2011

a BH is characterized solely by three parameters: mass, M_{BH} , angular momentum, J , and electric charge. Due to the highly conductive nature of astrophysical plasma, charged BHs are expected to quickly become neutralized by attracting opposite charges, and as such are not expected to be relevant (see in this context the discussion in Komissarov, 2022). Angular momentum is often expressed as the dimensionless spin, $a = J/cM_{\text{BH}}r_g$, where c is the speed of light, and $r_g = \frac{GM_{\text{BH}}}{c^2}$ is the gravitational radius. Spin values can range from 0 to 1. For $a = 0$, the Kerr metric simplifies to the Schwarzschild metric. The maximum value, $a = 1$, is derived from the “cosmic censorship” hypothesis (see Penrose, 1969; Krolak, 1986), which argues that singularities in the Universe have to be hidden within an horizon; for $a > 1$, the singularity would be exposed.

The line element of the Kerr metric is given by the following formula:

$$\begin{aligned} ds^2 = & -\left(1 - \frac{2r_g r}{\Sigma}\right) c^2 dt^2 + \frac{\Sigma}{\Delta} dr^2 - \frac{4r_g a}{\Sigma} r \sin^2 \theta c dt d\phi^2 + \\ & + \left(r^2 + a^2 r_g^2 + \frac{2a^2 r_g^3}{\Sigma} r \sin^2 \theta\right) \sin^2 \theta d\phi^2, \\ \Sigma = & r^2 + a^2 r_g^2 \cos^2 \theta, \\ \Delta = & r^2 - 2r_g r + a^2 r_g^2, \end{aligned}$$

where t , r , θ and ϕ are time, radial, polar and azimuthal coordinates respectively. The geometrical structure of a Kerr BH is far more complex than in the non-rotating case (see Fig. 1.1). In particular, two out of three critical surfaces are the outer event horizon, $r_h = r_g + \sqrt{r_g^2 - r_g^2 a^2}$, and the static limit, $r_e = r_g + \sqrt{r_g^2 - r_g^2 a^2 \cos^2 \theta}$. The outer event horizon, which is the event horizon in the usual sense of the term, is a “one-way” surface, once it is crossed it is impossible to go back, even for the light ray (Hartle, 2003). On the other hand, the static limit radius determines the region of space-time where no object can remain stationary relative to a distant observer (Lense-Thirring or “frame dragging” effect; Thirring, 1918). The region between these two surfaces is known as the “ergosphere.” Since the ergosphere lies outside the event horizon, it is possible for objects to escape from this region, potentially carrying away some of the BH rotational energy. This concept was first proposed by Penrose (1969). In the context of astrophysical jets production, the previously mentioned BZ mechanism is considered an electromagnetic form of the Penrose process (see e.g., Lasota et al., 2014).

Plasma accretion onto a rotating BH can be described and modeled in the framework of general relativistic magnetohydrodynamics (GRMHD), consisting of a set of notoriously difficult to solve hyperbolic partial differential equations. The BZ mechanism is a split monopole perturbation solution of the GRMHD equations in the force-free limit, which is technically valid only for slowly spinning BHs. This problem can be however addressed by the means of numerical simulations. Indeed, Tchekhovskoy, Narayan, and McKinney, 2011 showed in their GRMHD simulations an extraction of the BH angular momentum through the BZ process also in the case of very high BH spins, $a \lesssim 1$ (see Fig. 1.2).

It is important to note that BHs themselves do not emit electromagnetic (EM) radiation; therefore, any BH studies conducted through EM observations are indirect. Nevertheless, observations of a BH’s activity and its impact

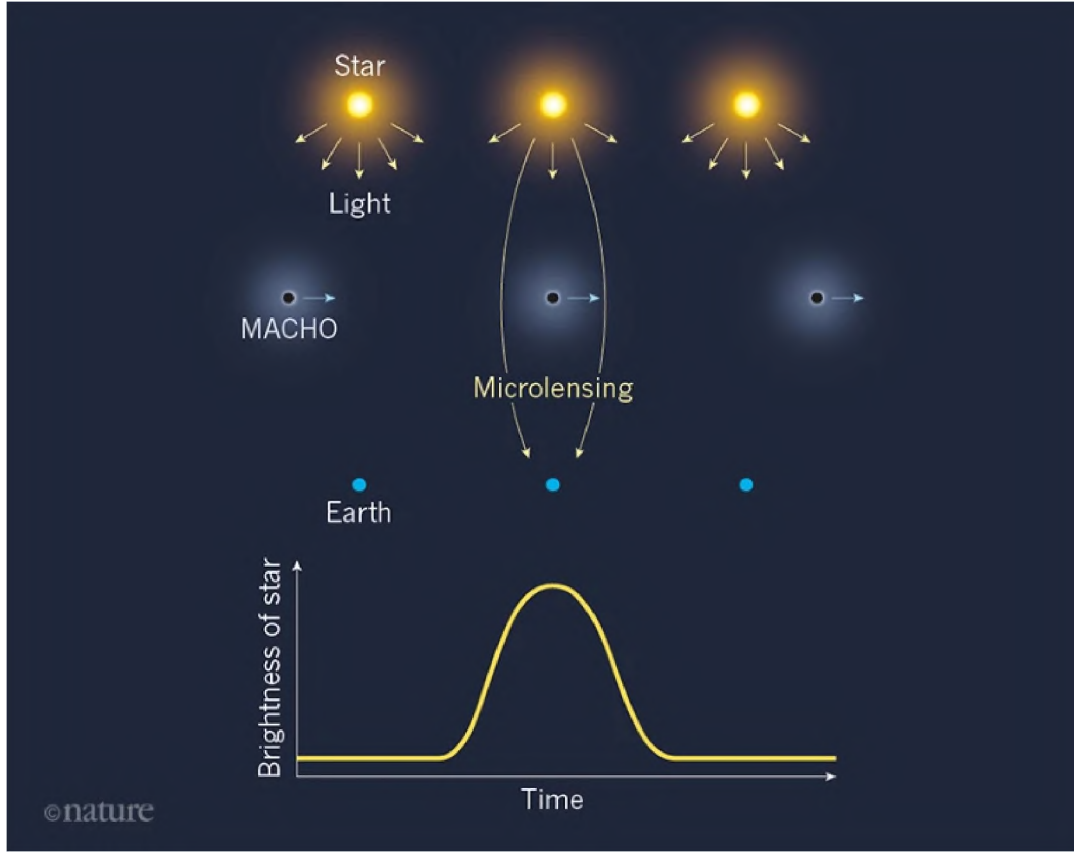


FIGURE 1.3: Schematic diagram depicting a observer (Earth), lens (MACHO in this case) source (star) geometry and the magnification pattern in the light-curve (brightness of star vs time).

Source: Pietrzyński, 2018

on surrounding astronomical systems sometimes allow for the assessment of the BH's mass and spin. In particular, stellar-mass BHs in our Galaxy can be detected due to the gravitational lensing (GL) or their activity within the X-ray binaries (XRBs). GL occurs when a BH passes through the observer's line of sight to a star. Then the path of light-rays can be bent in the space-time curved by a BH's gravitational potential (see Fig. 1.3). In some regimes of the lens mass and source-lens-observer distances it introduces a characteristic magnification pattern to the star's light-curve (Schneider, Ehlers, and Falco, 1992). XRBs are systems in which a gravitationally bound BH and a star orbit closely. EM radiation is produced when matter flows from the star to the BH, forming an accretion disk, its corona, and leading to the launch of winds and jet. In the case of both types of phenomena, detailed modeling can facilitate the estimation of BH masses, typically reported in the range of $5 - 20M_{\odot}$ (see, e.g., Rose et al., 2022; Özel et al., 2010), and, in the case of XRBs, also BH spin.

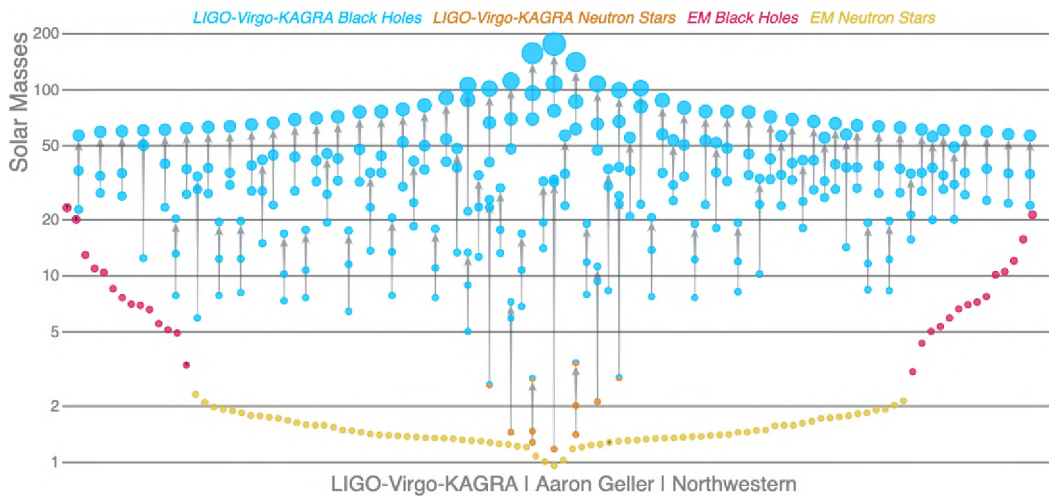


FIGURE 1.4: A collection of BH masses, measured through: gravitational wave detections (blue), XRB activity (red). Additionally plot shows known masses of neutron stars (orange and yellow).

Author: Aaron Geller, Northwestern University and LIGO-Virgo-KAGRA collaborations.

Source: Mandel and Farmer, 2022

Direct measurements of BH parameters were enabled as recently as 2015 with the first detection of gravitational waves by the Laser Interferometer Gravitational-Wave Observatory (LIGO, Abbott, Abbott, Abbott, Abernathy, et al., 2016). LIGO detects gravitational waves produced by mergers involving stellar-mass BHs and/or neutron stars. As of the beginning of 2023, LIGO has detected about 90 events, probing BHs with masses ranging from 5 to $200M_{\odot}$ (Abbott, Abbott, Abbott, Abraham, et al., 2019), see Fig. 1.4.

SMBHs lie on the other end of the BH mass distribution, with masses exceeding millions of solar masses. It is currently believed that every large galaxy in the Universe harbors a SMBH at its center (Kormendy and Ho, 2013). Depending on our distance to the host galaxy, its mass and level of SMBH activity, various methods can be employed to detect and assess SMBHs' masses. The most direct observations have been conducted by the Event Horizon Telescope (EHT), which, thanks to its ultra-high resolution, managed to observe the emission from gas accreting onto the SMBH at scales of the order of the BH event horizon. This type of study was possible for two objects: the SMBH at the center of our Galaxy, Sagittarius A*, and M87, one of the most extensively studied powerful, jetted AGN. Fig. 1.5 shows a composition of the M87 center and jet observations at different wavelengths, performed with different angular resolution (EHT MWL Science Working Group et al., 2021).

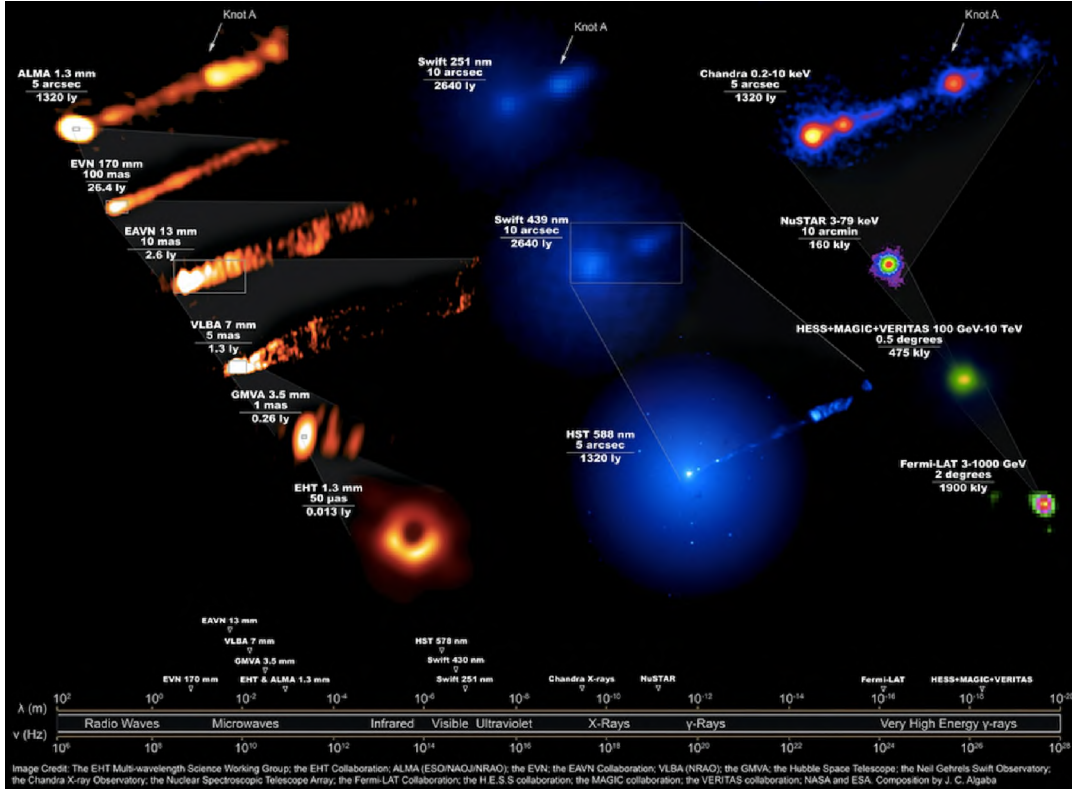


FIGURE 1.5: A composition of M87 observations at angular scales ranging from the μ arc seconds scale of the central SMBH's shadow structure, up to the arc seconds scale showing the relativistic jet's structure. Observation performed by multiple instruments in multiple EM bands, from radio (EHT, East Asia and European VLBI Networks, ALMA), through visible light/UV (The Hubble Space Telescope) to X-rays and γ -rays (Chandra X-ray Observatory, NuSTAR, MAGIC + HESS + VERITAS, The Fermi Gamma-ray Space Telescope).

Author of the composition: J.C. Algaba, Source: EHT MWL Science Working Group et al., 2021.

<https://science.nrao.edu/science/highlights/2021/multi-wavelength-properties-of-m87>

SMBH's masses can be assessed based on their gravitational impact on the gas and stars kinematics. This effect is limited to the most central part of a galaxy, i.e., to the so-called sphere of influence with a radius of

$$r_i = \frac{2GM_{\text{BH}}}{\sigma_*^2} \simeq 10^{-2} \left(\frac{M_{\text{BH}}}{10^8 M_\odot} \right) \left(\frac{\sigma_*}{200 \text{ km/s}} \right)^{-2} \text{ kpc}, \quad (1.1)$$

where σ_* is the velocity dispersion of stars or interstellar gas. Therefore, this method can be applied to only about 100 galaxies, in which angular scales corresponding to r_i can be spatially resolved by observations. Again, the situation is the most ideal in the case of Sagittarius A*, where it is possible to resolve kinematics of individual stars orbiting the Galactic Center (Gillessen et al., 2017). Reverberation mapping can further extend the sample of galaxies with somewhat direct estimations of SMBHs' masses. It is based on the

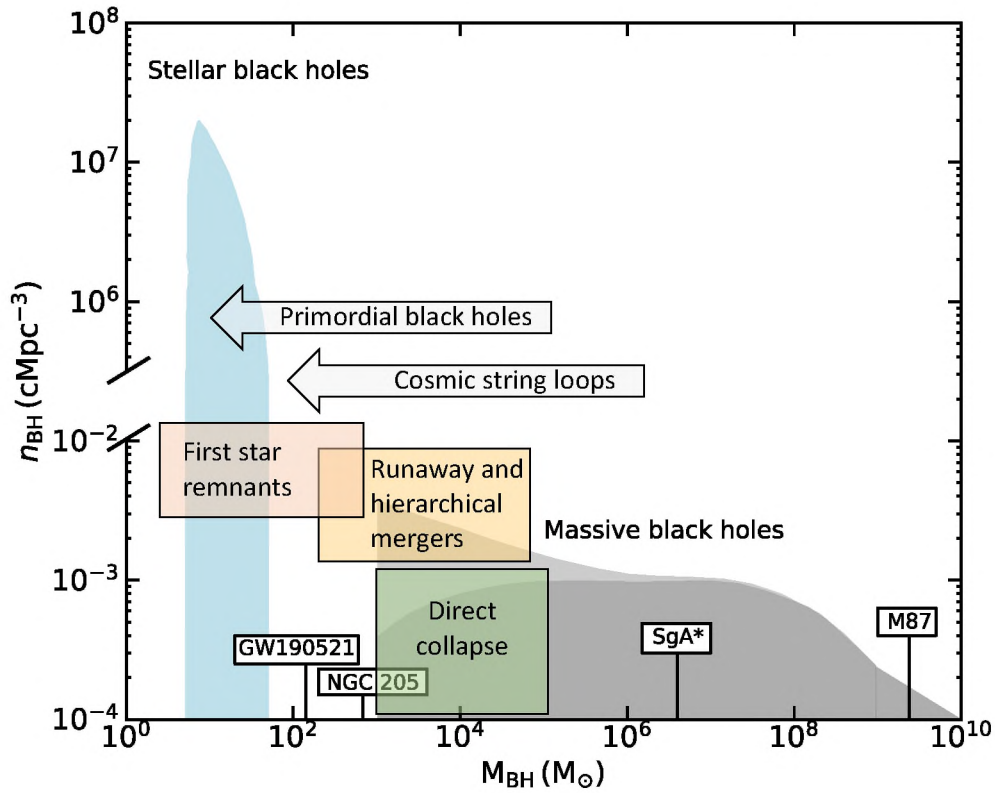


FIGURE 1.6: A cartoon showing BHs number density as a function of their mass.

Source: Volonteri, Habouzit, and Colpi, 2021.

measurements of the time lag between the variability of the disk continuum and the emission lines photoionized by the disk. Since the lines are believed to originate from the broad line region (BLR), the lag can be linked to the distance between the BLR and the nucleus (Peterson et al., 2004; Zu, Kochanek, and Peterson, 2011), while the width of the lines can be used to estimate the velocity dispersion of the line-emitting gas. Other methods are indirect and rely on established correlations between the properties of host galaxies and SMBHs' masses with available, above mentioned, direct mass estimates. These include the particularly tight and significant correlation between the SMBH mass and the velocity dispersion far beyond the sphere of influence, within the elliptical host galaxy/galactic bulge ($M_{\text{BH}} - \sigma_B$, Ferrarese and Merritt, 2000), or its luminosity of ($M_{\text{BH}} - L_B$, Kormendy and Richstone, 1995), the break frequency in the power density function of the X-ray variability (Czerny et al., 2001), and the temperature and intensity of X-ray halos (Gaspari et al., 2019).

The masses of stellar BHs and SMBHs differ by orders of magnitude (see Fig. 1.6). Many questions regarding the lack of detection of Intermediate Mass Black Holes (IMBHs), as well as the origins and evolutionary paths of

SMBHs, remain without a definitive answer (Volonteri, Habouzit, and Colpi, 2021). SMBHs are believed to originate from “seed” BHs, which were either light, with masses similar to those of stellar BHs formed in the Galaxy (i.e., $10 - 100 M_\odot$), or heavy ($10^4 - 10^5 M_\odot$). Heavy “seed” BHs could have been formed in the early Universe. The low metallicity and different physical conditions in the interstellar medium back then, may have facilitated the direct collapse of massive gas halos, and the existence of much more massive Population III stars. Recent observations of powerful quasars at very high red-shifts, performed by the James Webb Space Telescope, seem to support the scenario of heavy seeds (see e.g., Bogdán et al., 2024). These “seed” BHs are believed to increase their masses through accretion and mergers. Growth through accretion should be, in principal, restricted by the Eddington limit, $\dot{M}_{\text{Edd}}$, derived from the balance between the gravitational force on accreting matter and the radiation pressure:

$$\dot{M}_{\text{Edd}} = \frac{4\pi G M_{\text{BH}} m_p}{\epsilon c \sigma_T} \simeq 2.2 \epsilon^{-1} \left(\frac{M_{\text{BH}}}{10^9 M_\odot} \right) M_\odot \text{yr}^{-1}, \quad (1.2)$$

where ϵ is the radiative efficiency of the accretion process. High accretion rates lead to an increase in radiation pressure, which halts the accretion, setting a limit on the BH’s mass growth rate. However, in reality when disks are sufficiently dense to trap photons, objects can exhibit super-Eddington accretion rates (see e.g., Abramowicz et al., 1988; Dotan and Shaviv, 2011; Jiang, Stone, and Davis, 2019). The importance and exact course of mergers, both within stellar clusters and between galaxies, are still subjects of debate.

This picture is further complicated by the fact that SMBHs do not evolve in isolation; rather, AGN co-evolve with their host galaxies, undergoing cycles of activity influenced by various factors, including tidal disruptions of stars and galaxy mergers. The previously mentioned $M_{\text{BH}}-\sigma_B$ and $M_{\text{BH}}-L_B$ relations (see Fig. 1.7) are expressed by the following formula (Kormendy and Ho, 2013):

$$\begin{aligned} \frac{M_{\text{BH}}}{10^9 M_\odot} &= 0.542 \times \left(\frac{L_B}{10^{11} L_\odot} \right)^{1.21} \\ \frac{M_{\text{BH}}}{10^9 M_\odot} &= 0.309 \times \left(\frac{\sigma_B}{200 \text{kms}^{-1}} \right)^{4.38} \end{aligned}$$

where $L_\odot$ is the Solar luminosity in the K-filter. These tight correlations imply a very close link between the SMBH’s mass and the global features of the

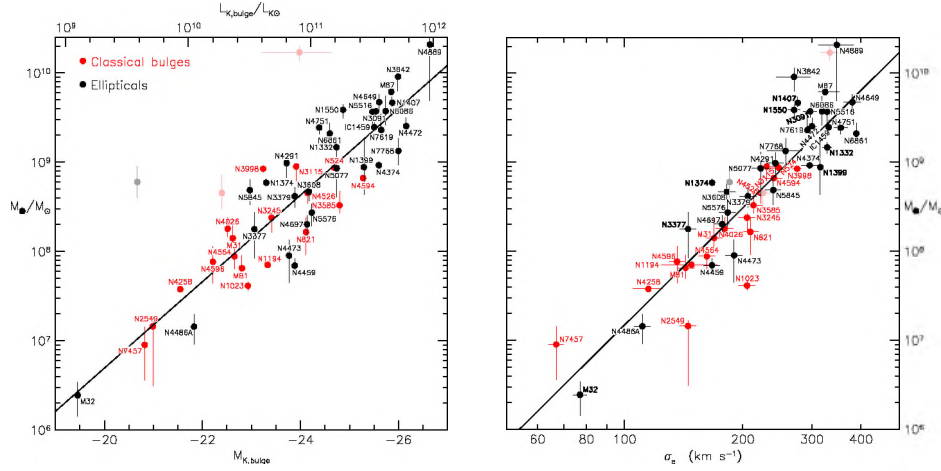


FIGURE 1.7: $M_{\text{BH}}-\sigma_B$ and $M_{\text{BH}}-L_B$ relations, which link SMBH mass to the velocity dispersion and luminosity of the bulges (red) and elliptical galaxies (black). (Kormendy and Ho, 2013)

Source: Kormendy and Ho, 2013

host galaxy, meaning that the impact of the central SMBH is noticeable at distances much larger than the gravitational influence radius. This co-evolution is also evident on cosmological scales; in particular, the co-moving number density of quasars appears to be closely related to cosmic star-formation history, as both peak around redshift $z \sim 2$ (Madau and Dickinson, 2014). It is widely believed, that this co-evolution involves a feedback between the host galaxy that provides the fuel for the SMBH growth and activity on the one hand, and the AGN impact on the host through disk outflows, emission, and jets. Highly collimated, relativistic jets may in particular propagate through host galaxies, interacting with the interstellar medium, and altering its state (Heckman and Best, 2014). Notably, some jets extend into the extragalactic medium on kiloparsec and even megaparsec scales, depositing enormous amounts of energy carried from their very launching sites, i.e., the ergospheres of supermassive BHs to the inter-galactic medium (Blandford, Meier, and Readhead, 2019).

1.1 Dissertation outline

This Ph.D. dissertation consists of four scientific articles that delve into the physics of black holes and their activity across different astrophysical systems. The different scales involved and the nature of the analysed problems, enforce a variety of applied techniques and approaches. I arranged the articles according to the increasing distance of the discussed regions from the BHs.

Chapter 2 consists of the article “Accretion-induced Black Hole Spin-up Revised by Numerical General Relativistic MHD”, which concerns black holes themselves. When a black hole emerges from a collapsing hyper-massive star, ultra-relativistic jets form, causing a gamma ray burst (a “collapsar” model). In this regime, the MHD equations resist simplification, and general relativity is pivotal. Consequently, analytical analysis becomes challenging, making GRMHD simulations the only way to study these phenomena. We present the results of the GRMHD simulations conducted with the modified version of the High Accuracy Relativistic Magnetohydrodynamics code (HARM, Gammie, McKinney, and Tóth, 2003). This adaptation tracks the evolution of the black hole’s mass and spin during accretion (Janiuk, Sukova, and Palit, 2018). Since the BH’s mass and spin immediately after the collapse remain unknown, we followed the M_{BH} and a evolution for different initial a values, the amount of angular momentum in the accreting gas, and strength of the magnetic field. We showed that a BH with high mass and spin can be produced for a wide range of initial conditions, even if the values of these two parameters were low just after the collapse. Preliminary results of the simulations without the magnetic field (section 3.2) were included in my master’s thesis.

The following dissertation chapters shift the focus towards the activity associated with SMBHs. In particular, the third chapter focuses on repetitive “U”-shaped features appearing in the optical light-curve of the blazar S5 0716+714. Blazars emission is dominated by relativistic jets, and is produced at distances $\sim 10^4 r_g \simeq 10^{-5} \times \left(\frac{M_{\text{BH}}}{10^9 M_{\odot}} \right) \text{ pc}$ from the central SMBH (Sikora, Begelman, and Rees, 1994). In general, AGN jets are highly variable over wide ranges of time scales and wavelengths. A debate exists, if some part of this variability can be attributed to the GL of jet segments (e.g., Gopal-Krishna and Subramanian, 1991). Chapter 3 contains the article titled “Possible Gravitational Microlensing Events in the Optical Light-curve of Active Galaxy S5 0716+714”. In this analysis, we examined “U”-shaped features

spotted across all filters in the Whole Earth Blazar Telescope (WEBT) data from 2003 and 2014 (Ostorero et al., 2006; Bhatta et al., 2016), and in the Transiting Exoplanet Survey Satellite (TESS) data collected in 2019/2020 (Raiteri et al., 2021). Drawing from prior studies of Peirson et al., 2022 and Vedantham et al., 2017, who attributed similar features in another blazar’s radio light-curve to gravitational lensing events (albeit on much longer timescales), we proposed that S5 0716+714’s optical features could also be due to lensing by a binary lens. The fitting shows that gravitational lensing by a single binary lens, located in the source’s host galaxy, may explain the observed phenomenon. Depending on the source-lens distance, the lens’s mass might align with an Intermediate Mass Black Hole, the missing link in between stellar mass BHs and SMBHs.

Chapter 4 consists of the article “On the Origin of the X-ray Emission in Heavily Obscured Compact Radio Sources.” We studied the broadband emission of Compact Symmetric Objects (CSOs), a subset of AGN known for their compact, young radio structures, which may provide valuable insights into the initial evolutionary stages of radio galaxies. The X-ray continuum emission in AGN is typically associated with the “disk corona”, and the narrow K_α iron line in the spectrum signifies the subsequent reflection of the X-ray continuum from the circumnuclear dusty tori. Various models have been proposed for the corona’s location and structure. In this project, we explored the origins of the X-ray emission, showing that for CSOs, the continuum might arise from radio lobes still embedded in the dusty tori. These lobes constitute a reservoir of relativistic electrons that can produce high-energy emission through Comptonization, which is then reflected from the torus. In this scenario, the observed iron lines are produced by the reflection of the lobes’ continuum from the surrounding cold dust. To test this scenario, we modeled the broadband emission of three sources with distinct K_α lines and significant absorption features in their X-ray spectra, characterized by linear sizes of radio lobes ranging from 7 to 25 pc: J1407+2827, J1511+0518, and J2022+6137. Through analytical estimates, we showed that the high-energy emission of the most compact radio lobes indeed may dominate over the disk-corona radiative output. We also successfully employed the model of emission from the expanding radio lobes (Stawarz et al., 2008) to match the observed flux levels across the broadband spectrum.

Depending on various factors, such as jet power and the surrounding medium density, jets may propagate outside of the AGN host galaxy, reaching a different stage of evolution in the extra-galactic medium (Mukherjee

et al., 2016). At this stage, jets become axisymmetric, and certain approximations concerning their velocity profiles, geometry, and magnetic field can be applied. Chapter 5 consists of the paper titled “Magnetization of Relativistic Current-carrying Jets with Radial Velocity Shear”, where we examined that stage of jets evolution through analytical calculations and semi-analytical modeling within the framework of relativistic magnetohydrodynamics (RMHD). We were particularly interested in the values of the magnetization, defined as the ratio of the magnetic energy flux to the particle energy flux along the flow. This parameter contains information on the relative importance of magnetic reconnections and strong shocks for the acceleration of jet particles to ultra-relativistic energies. We showed that relativistic jets at large distances from their launching sites, if in magneto-hydrostatic equilibrium, tend to be matter-dominated. For certain magnetic and velocity profiles, this could be proven analytically. In more complex cases, we explored numerical solutions and discovered that achieving even slightly magnetically dominated jets requires extreme and possibly nonphysical pressure gradients. Although this finding might be somewhat unexpected, it complements earlier studies on jet evolution. These studies indicate that collimating and accelerating magnetohydrodynamic jets, confined by ambient pressure, remain dominated by Poynting flux until they reach logarithmically large distances from the central black hole (Lyubarsky, 2009; Lyubarsky, 2010).

Chapter 2

Accretion-induced Black Hole Spin-up Revised by Numerical General Relativistic MHD



Accretion-induced Black Hole Spin-up Revised by Numerical General Relativistic MHD

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Abstract

We investigate the accretion-induced spin-up of the black hole via numerical simulations. Our method is based on general relativistic magnetohydrodynamics of the slowly rotating flows in the Kerr metric, where possibly transonic shock fronts may form. We account for the changing black hole mass and spin during accretion that enforces dynamical evolution of the spacetime metric. We first study nonmagnetized flows with shocks, and we also include magnetic field endowed in the gas. The aim of this study is to verify whether the high-mass black holes may be produced with large spins, even though at birth the collapsars might have contained slowly or moderately spinning cores. In this way, we put constraints on the content of angular momentum in the collapsing massive stars. Our studies are also showing that shock fronts and magnetic fields may halt accretion and limit the black hole spin-up in the exploding supernovae.

Unified Astronomy Thesaurus concepts: Bondi accretion (174); Black hole physics (159); Magnetohydrodynamics (1964)

1. Introduction

Evolution of massive stars ends when the compact remnant is formed in its core, becoming either a neutron star or a black hole. The nature of the remnant mostly depends on the progenitor mass on the main sequence (Janka et al. 2007; Ott et al. 2018) and is also dependent on its metallicity (Nomoto et al. 2003; Eldridge & Tout 2004). Massive stars heavier than $10 M_{\odot}$ are frequently born in binary systems, or in multiple stellar configurations (Kobulnicky & Fryer 2007; Sana et al. 2012). The evolution of such a system is therefore driven by the mass-exchange process, which can proceed either via stable transfer or via an unstable episode called a common-envelope phase (see review by Langer 2012). On the other hand, some fraction of supernovae might originate from single-star progenitors. In this case, the stellar evolution may involve a large amount of stellar wind being ejected at the final stages, leading to the self-stripping of the stellar envelope (Heger et al. 2003). The final mass of the remnant after the core-collapse supernova depends on whether the star evolved in a binary system or as a single object, because of an essentially different presupernova structure revealed by the different stellar evolution models (Hirai et al. 2020).

The measurements of compact remnant masses are possible now independently through the X-ray observations and gravitational wave detections. In the case of X-ray binaries, the known masses of black holes range from ~ 1.6 up to $18 M_{\odot}$ (Casares et al. 2017). In these systems, the first compact object went through a phase of mass transfer, while the progenitor must have been stripped off its hydrogen envelope (Podsiadlowski et al. 2003). It is rather unclear, though, how the envelope stripping affects the final mass of the compact remnant. As for the spin of such a black hole, it is natural to expect that the rotation of the envelope also adds the angular momentum to the stellar core (see, however, the recent study by Fuller et al. 2019 on the angular momentum evolution in stellar cores).

The masses of black holes detected by gravitational wave analysis are systematically higher than those obtained from

X-ray observations (Perna et al. 2019). The masses of merging components are between 10 and $85 M_{\odot}$, with the merger products as massive as $142 M_{\odot}$, for GW190521 (Abbott 2020). The past history of these most massive black holes is subject to many speculations. The merging black holes might have been formed in the same binary system and hence have experienced a past mass-exchange episode. The results of population synthesis modeling and comparison with the data about current merger rates seem to support such a scenario (Belczynski et al. 2016; Mapelli et al. 2017). Alternatively, the premerger history might contain dynamical encounters in dense stellar clusters (Hong et al. 2018; Samsing et al. 2018), which means that the black holes have been formed independently, possibly from single-star progenitors.

The information on the black hole spins is important to determine which of the scenarios would be favored in particular objects. If the spin vectors of merger components are aligned, then most probably the black holes have been formed in a binary system, while the dynamical encounters allow for a more random distribution of the spins (Farr et al. 2017). The preferences for isotropic spin distribution have to be verified, because the data allow only for determination of an effective spin, which is a combination weighted by the black hole masses (see, e.g., Mandel & Fragos 2020, for the case study of GW190412). In addition, the large masses of merging black holes should be correlated with their low spins, $a < 0.5$, regardless of their formation channel (Amaro-Seoane & Chen 2016).

These new insights into the black hole mass and spin distributions, resulting from the binary interactions or wind-driven stripping, are tackled with the stellar structure and evolution models and the supernova explosion physics. Moreover, the paths of black hole formation are affected by the action of magnetic turbulences (Obergaulinger & Aloy 2017). The finally resulting stellar black hole properties and the fraction of prototypical stellar cores that should result in black hole collapsars depend therefore on a combination of several factors, among which the structure of the progenitor star and

the profile of specific angular momentum are probably the most crucial.

Some information about the rotational velocity in the presupernova star may be inferred from its neutrino emission characteristics. The recent simulations show that the evolution of standing accretion shock instability (SASI), which strongly depends on both stellar rotation rate and location of the convective zones, might leave imprints on the neutrino Fourier power spectrum (Walk et al. 2020). We are still awaiting the observational verification of such events.

In this study we analyze the properties of collapsing massive stars, whose envelopes have already been mostly stripped off. We simulate numerically the central parts of a collapsar, where the massive black hole is forming via accretion of mass and angular momentum from the innermost parts of the envelope. We perform a parameter study, assuming different levels of rotation in the accreting material and different initial spins of the black hole that might have resulted from the past stellar evolution before the envelope stripping. We account for both magnetized and nonmagnetized models, and in the latter we probe the formation of transient shocks in the collapsing envelope. With the magnetohydrodynamic (MHD) simulations, we model weakly and strongly magnetized collapsars, where in the latter we adopt the condition of equipartition between thermal and magnetic energy. All of our simulations are done in full general relativity, and the coupled evolution of the black hole spin and mass during the collapse process is considered with the self-consistent updates of the Kerr metric (see Janiuk et al. 2018). The aim of our study is to provide some general constraints on the final mass and spin of the black hole, depending on the collapsar properties and progenitor’s history. We also speculate about a possible electromagnetic transient that can occur in some cases, in contrast to the events that leave no luminous counterparts.

In the paper by Perna et al. (2014) the authors study the presence of fallback disks around young neutron stars. The angular momentum distribution is taken from the presupernova models calculated with public ode MESA, for the stars of masses between 13 and 40 $M_{\odot}$. The specific angular momentum content eventually determines the post-fallback circularization radius. They identify regions in their model parameter space that lead to extended, long-lived disks around black holes, and they find that the physical conditions in these disks may be conducive to planet formation. In these simulations, the angular momentum profile is based on the reliable assumptions resulting from stellar evolution. On the other hand, these simulations were done in Newtonian gravity; hence, all the general relativistic (GR) effects, such as the circularization radius dependence on the black hole mass and spin changes, were ignored.

Recently, the problem of failed supernovae from nonrotating supergiant stars was studied by Quataert et al. (2019). They show that net angular momentum is not a necessary condition to form the accretion disk, because of convective motions in the outer parts of the supergiant and a random angular momentum in supergiant convection zones that exceeds that of the last stable circular orbit of a black hole by a factor of ~ 10 . The result is obtained using analytic estimates and Cartesian box simulations in the Bussinesq approximation.

In the current work, we follow the approach started in the work of Lee & Ramirez-Ruiz (2006). The idea described a ballistic motion of particles in the gravitational potential of the

black hole, while the authors studied a freefall of the fluid. A superimposed rotation profile was given by an analytic function, $l_{\text{spec}} \propto \sin^2(\theta)$. Neutrino cooling and contributions to the pressure from degenerate electrons were considered by those authors because the discussion was focused on the resulting gamma-ray burst power.

A corresponding initial condition setup can be found in De Villiers et al. (2005). In their work, the authors conducted a series of GR MHD 2D simulations with a 3D follow-up for chosen models. Initially, the physical system consisted of a black hole (with different spin values probed), a Keplerian disk, and the background described by the Bondi solution. The spacetime metric was fixed and did not evolve under the spin and mass evolution. A significant difference in comparison with our simulations is the presence of a Keplerian disk. Also, the authors assumed a smaller ratio between accreting mass and black hole mass, to justify the fixed metric assumption. The main focus of their work was to discuss the jet-launching process, and they obtained a jet present between 0.1 and 1.4 s, with a significant contribution from the coronal wind.

In this paper we perform the first multidimensional GR simulations of uniformly rotating, low angular momentum, magnetized flows, where we take into account the Kerr metric evolution due to changing black hole mass and spin. We build our models on the initial conditions of the slowly rotating quasi-spherical transonic flow, and we follow up on the other recent works of our group. However, in Janiuk et al. (2018) and in Murguía-Berthier et al. (2020), the magnetic fields were neglected. The latter paper, which actually also neglected the Kerr metric evolution effects, gave the main focus to the comparison of the constraints on the angular momentum resulting from the GR hydrosimulations, with the realistic models of the stellar structure and rotation profiles obtained from MESA simulations.

The structure of the article is as follows. In Section 2 we describe the theoretical framework of our simulations. In Section 2.1 we present the analytical formulation of transonic spherical accretion onto a black hole, in Section 2.2 we introduce the slow angular momentum prescription, in Section 2.4 we describe our method for the Kerr spacetime evolution, in Section 2.5 we introduce the physical scaling and parameters of our collapsar model, and in Section 2.6 we describe the initial configuration for magnetic fields used in the MHD simulations. The results of our numerical analysis are presented in Section 3, and in particular, the nonmagnetized models are presented in Section 3.2, while the models with magnetic fields are presented in Section 3.3. The discussion and conclusions are presented in Section 4.

2. Black Hole Accretion with Low Specific Angular Momentum

In the simple collapsar model that we adopt here as an initial condition, we assume that the star has had enough time to reach a quasi-steady state at the end of its stellar evolution. Hence, we adopt a quasi-spherical configuration of density, resulting from the Bondi prescription (Bondi 1952) and endowed with a small angular momentum. During the time-dependent simulation, however, we do not add mass to the envelope during the time evolution, so the accretion rate is dynamically changing. Moreover, rotation changes the matter distribution, and a disk-like structure forms at the equatorial plane. The microphysics is treated in a simplified way, and we use the adiabatic index of

$\gamma = 4/3$, which describes the stellar interior, accounting for the contribution from the photon pressure.

2.1. Initial Condition for Quasi-spherical Flow

The distribution of density as a function of radius comes from the solution of transonic accretion flow in spherical geometry and is similar to that of Suková & Janiuk (2015), Suková et al. (2017), and Palit et al. (2019). The initial density profile and the radial component of the velocity (u^r) of the material are determined by the relativistic version of the Bernoulli equation (Michel 1972; Hawley et al. 1984). In this formalism, the critical point (r_s), where the flow becomes supersonic, is set as a parameter. Here we take the value of $r_s = 80r_g$.

The fluid is considered a polytrope with a pressure $P = K\rho^\gamma$, where ρ is the density, $\gamma = 4/3$ is the adiabatic index, and K is the constant specific entropy, in this case taken to be that of a relativistic gas.

Once the critical point, r_s , is set, the velocity at that point is

$$(u_s^r)^2 = \frac{GM}{2r_s}, \quad (1)$$

where r is the radial coordinate, M is the mass of the black hole, and the sound speed is

$$c_s^2 = \frac{\gamma \frac{P_s}{\rho_s}}{1 + \frac{\gamma}{\gamma-1} \frac{P_s}{\rho_s}}. \quad (2)$$

The constant specific entropy can be obtained using the sound speed:

$$K = \frac{c_s^2}{\rho_s^{\gamma-1}\gamma}, \quad (3)$$

where the density is set by the mass accretion rate $\dot{M}$:

$$\rho = \frac{\dot{M}}{r^2 u^r}. \quad (4)$$

The radial velocity profile is obtained by numerically solving the equation (Shapiro & Teukolsky 1986)

$$\left(1 + \frac{\gamma}{\gamma-1} \frac{P}{\rho}\right)^2 \left(1 - \frac{2GM}{r} + (u^r)^2\right) = \text{constant}, \quad (5)$$

and the proper velocity of the fluid at infinity is zero, while it is equal to the speed of light at the horizon.

In construction of the numerical solution, we integrate the Equation (5) upward and downward of the critical point, and the normalized accretion rate $\dot{M} = 1$ is used (note that our code works in dimensionless units of $G = c = 1$). To convert the quantities into physical units, we adopt the parameterization of density in the accreting gas as in Janiuk et al. (2018).

2.2. Initial Conditions for Small Angular Momentum Configuration

The accreting material is endowed with small angular momentum scaled to the one at the circularization radius of r_{circ} , which is the innermost stable circular orbit (ISCO) radius (equal to $6r_g$ for a nonrotating black hole).

The angular momentum is further scaled with the fraction of the angular momentum at the ISCO, S , and scaled with polar angle to have its maximum value on the equatorial plane, at

$$\theta = \pi/2:$$

$$l = Sl_{\text{isco}} \sin^2 \theta, \quad (6)$$

where l is the specific angular momentum, defined as $l = u_\phi = g_{\phi\mu} u^\mu$. Here l_{isco} is the specific angular momentum at the ISCO of the black hole, and θ is the polar coordinate. In order to get the angular velocity at each point, we use the auxiliary energy ϵ_{isco} and auxiliary angular momentum at the ISCO (l_{isco}):

$$\epsilon_{\text{isco}} = \frac{1 - 2/r_{\text{isco}} + a/r_{\text{isco}}^{3/2}}{\sqrt{1 - 3/r_{\text{isco}} + 2a/r_{\text{isco}}^{3/2}}} \quad (7)$$

$$l_{\text{isco}} = \frac{r_{\text{isco}}^{1/2} - 2a/r_{\text{isco}} + a^2/r_{\text{isco}}^{3/2}}{\sqrt{1 - 3/r_{\text{isco}} + 2a/r_{\text{isco}}^{3/2}}}, \quad (8)$$

where a is the spin of the black hole, and the radius of the ISCO (r_{isco}) changes self-consistently with changing spin. In these units, black hole mass is $M = 1$.

We can then construct the angular velocity in the Boyer–Lindquist coordinates as

$$u^\phi = -g^{t\phi} \epsilon_{\text{isco}} + g^{\phi\phi} l_{\text{isco}}, \quad (9)$$

where the components of the metric are for a Kerr black hole $g^{t\phi} = -2ar/(\Sigma\Delta)$ and $g^{\phi\phi} = (\Delta - a^2 \sin^2 \theta)/(\Sigma\Delta \sin^2 \theta)$, and $\Sigma = r^2 + a^2 \cos^2 \theta$, $\Delta = r^2 - 2r + a^2$.

Normalization of angular momentum content in the collapsar's envelope is done with S , which is a free parameter of our model. Note that $S = 0$ reduces to the Bondi spherical accretion.

2.3. Time-dependent General Relativistic MHD Code

We use the GR MHD code *HARM* (Gammie et al. 2003; Noble et al. 2006). In this version of the code, we use the method developed in Janiuk et al. (2018) to account for the change of Kerr metric due to the dynamically increasing black hole mass and its spin.

The *HARM* code is a finite-volume, shock-capturing scheme that solves the hyperbolic system of the partial differential equations of GR MHD. The numerical scheme is using the plasma energy momentum tensor, $T^{\mu\nu}$, with contributions from gas and electromagnetic field

$$T_{(m)}^{\mu\nu} = \rho h u^\mu u^\nu + p g^{\mu\nu} \quad (10)$$

$$T_{(em)}^{\mu\nu} = b^\kappa b_\kappa u^\mu u^\nu + \frac{1}{2} b^\kappa b_\kappa g^{\mu\nu} - b^\mu b^\nu \quad (11)$$

$$T^{\mu\nu} = T_{(m)}^{\mu\nu} + T_{(em)}^{\mu\nu}, \quad (12)$$

where u^μ is the four-velocity of gas, u denotes internal energy density, b^μ is magnetic four-vector, and h is the fluid specific enthalpy. The continuity and momentum conservation equations read

$$(\rho u^\mu)_{;\mu} = 0; \quad T_{\nu;\mu}^\mu = 0. \quad (13)$$

They are brought in conservative form, by implementing a Harten, Lax, van Leer (HLL) solver to calculate numerically the corresponding fluxes.

In terms of the Boyer–Lindquist coordinates, (r, θ, ϕ) , the black hole is located at $0 < r \leq r_h$, where $r_h = (1 + \sqrt{1 - a^2})r_g$ is the horizon radius of a rotating

black hole with mass M and angular momentum J in geometrized units, $r_g = GM/c^2$, and a is the dimensionless Kerr parameter, $a = J/(Mc)$, $0 \leq a \leq 1$.

The *HARM* code does not perform the integration in the Boyer–Lindquist coordinates, but instead in the so-called modified Kerr–Schild ones: $t, x^{(1)}, x^{(2)}, \phi$ (Noble et al. 2006). The transformation between the coordinate systems is given by $r = R_0 + \exp[x^{(1)}]$ and $\theta = \frac{\pi}{2}(1 + x^{(2)}) + \frac{1-h}{2}\sin[\pi(1 + x^{(2)})]$, where R_0 is the innermost radial distance of the grid, $0 \leq x^{(2)} \leq 1$, and h is a parameter that determines the concentration of points at the midplane. In our models we use $h = 0.3$ (notice that for $h = 1$ and a uniform grid on $x^{(2)}$ we obtain an equally spaced grid on θ , while for $h = 1$ the points concentrate on the midplane). The exponential transformation in the r -direction leads to higher resolution in the innermost regions. The inner and outer radii of our grid are located at $0.98r_h$ and 10^3r_g , respectively. Our grid resolution is 256×256 .

2.4. Kerr Spacetime Evolution

The changing black hole spin and mass affect the spacetime metric. In our simulation, the spacetime is described by a sequence of quasi-stationary Kerr solutions, where the metric is changed discretely between the consecutive time steps, according to mass and spin increments, $\Delta M = (M^i/M^0 - 1)$, and $\Delta a = (J/M^i - a^{i-1}/M^i\dot{E})\Delta t$, where M^i denotes the current black hole mass at time $t > 0$, M^0 denotes initial black hole mass at $t = 0$, and $\dot{J}$ and $\dot{E}$ are the flux of angular momentum and energy flux transmitted through the black hole horizon at a given time (see Janiuk et al. 2018). The six nontrivial components of the Kerr metric, namely, $g_{tt}, g_{tr}, g_{t\phi}, g_{rr}, g_{r\phi}, g_{\phi\phi}$, are updated at every time step, and they get new values.

This procedure is developed to take into account the change of black hole spin and mass, as well as the metric update, but does not consider the effects of self-gravity of the diluted envelope and the effect of its changing mass on the metric. We argue that the self-force of the envelope would have negligible impact on our simulation results; however, we plan to address this issue in more detail and make quantitative comparisons in future work.

We notice that the external material does have an impact on the overall Kerr metric evolution; however, in the present simulations we neglect this effect. The adopted density of matter in the accreting cloud is large (see Section 2.5 below). Nevertheless, the large size of the computational domain ensures that the cloud is still not as compact as the central black hole, so the black hole still dominates the gravitational potential in this system.

The effect of the collapsing cloud on the metric evolution is driven by the self-force, which is induced on the matter element located at every grid point by the mass contained inside the sphere of that radius. It effectively acts as an additional perturbative term, apart from the increasing mass of the black hole. This perturbative term is changing with time during the collapse (see Palit et al. 2020). We verified that this term is large only at the beginning of collapse, when the cloud mass is large, and very quickly decreases with time.

2.5. The Physical Simulation Setup

The code *HARM* operates in dimensionless units $G = c = M = 1$, so the units do not appear in the metric coefficients. Because we want to model a particular case of collapsing star, we convert our simulation results to physical scale, which allows us to obtain the proper intuition about astrophysical objects of interest. Following Janiuk et al. (2013), we use the value of the total (volume-integrated) initial mass of the accreting matter to scale the density in the spatial domain, and we use the unit conversion system as

$$L_{\text{unit}} = r_g = \frac{GM}{c^2} = 1.48 \cdot 10^5 \frac{M}{M_\odot} \text{cm}$$

$$T_{\text{unit}} = \frac{r_g}{c} = 4.9 \cdot 10^{-6} \frac{M}{M_\odot} \text{s}$$

for the spatial and time units, respectively. The density scale is related to the spatial unit as $D_{\text{unit}} = M_{\text{scale}}/L_{\text{unit}}^3$.

In this work, the collapsing stars are modeled with the following global parameters: initial mass of the core black hole M_{BH}^0 , the cloud mass enclosed within the simulation volume M_{cloud} , and the initial black hole dimensionless spin A_0 . The fiducial parameters are $M_{\text{BH}}^0 = 3 M_\odot$ and $A_0 = 0.3 - 0.85$, while the cloud mass is $25 M_\odot$, as resulting from its size in geometrical units and adopted density scaling. The accretion rate onto the black hole, measured as the mass flux transported through the horizon, is varying with time. Its value is expressed in physical units of $M_\odot \text{s}^{-1}$, as converted from the density scaling and time unit.

2.6. Magnetic Field Configuration for MHD Models

We assume that the initial accreting cloud is embedded in a poloidal magnetic field, prescribed with the vector potential of

$$A_\phi \propto (1 - \cos(\theta)), \quad (14)$$

where the normalization comes from the ratio between the gas pressure, given by $p_{\text{gas}} = (\gamma - 1)u$, and magnetic pressure, given by $p_{\text{mag}} = \frac{1}{2}b_\mu b^\mu$. Our code parameter is the ratio of initial gas to magnetic pressure, $\beta_0 \equiv p_{\text{gas}}/p_{\text{mag}}$. As a consequence, the magnetized flows considered here are not in equilibrium, and the angular momentum is transported also through the rotationally supported disk in the equatorial plane, thus affecting the accretion.

We examine here models with initial $\beta_0 = 1.0$ and $\beta_0 = 100$, and we also probe different values of A_0 , equal to 0.3, 0.5, and 0.85 for the initial spin of the black hole. In addition, every set of magnetized collapsar models includes different rotation parameters, S . The β parameter is normalized to the gas pressure at the inner boundary of the grid, whose position depends on A_0 (due to the scaling of the ISCO).

3. Results

3.1. General Structure of the Flow

The overall evolution follows a similar pattern for all hydrodynamical models that have been investigated. We always start from a spherically symmetric distribution of density, which is quickly broken by the imposed azimuthal velocity component that is maximal at the equator. In consequence, the distribution of density becomes flatter toward

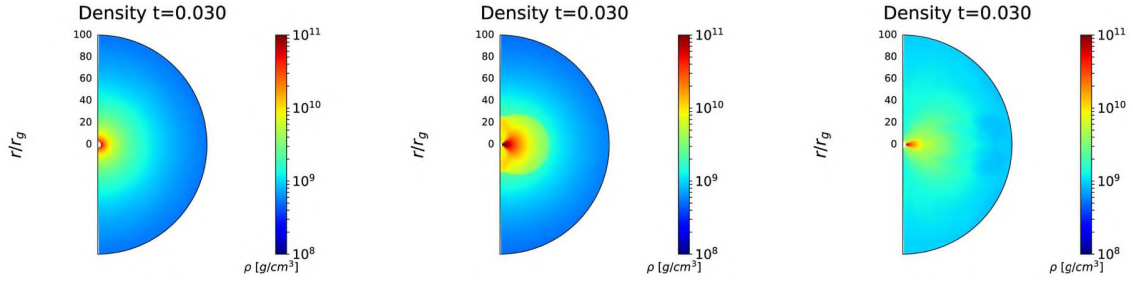


Figure 1. Density distribution in the innermost regions of the accretion flow, within $100 r_g$. Snapshots show models with $A_0 = 0.3$ and rotation parameter of $S = 0.4$ (NS-T-04), $S = 1.0$ (NS-T-10), and $S = 1.4$ (NS-T-14), from left to right, as taken at the beginning of these simulations. Time snapshots were taken at time $t = 0.03$ s.

the equator. For the radial velocity of the flow, we always start from transonic Bondi solution, and the sonic radius is initially located at $80 r_g$. In most cases, the angular momentum imposed on the flow is changing drastically the topology of solutions, and the sonic surface shape becomes aspherical (hence, we use the term “sonic surface” rather than the “point”). Also, in some cases multiple transonic solutions can be found, depending on the model, and an inner sonic point is found closer to the black hole, where a transient shock develops. Inside this region, the “bubble,” or “mini-disk,” is formed. The mini-disk is rotationally supported, so in general its size increases for a higher angular momentum content. It is present for a relatively long simulation time, and its innermost, densest part is enclosed with $\sim 20 r_g$. Its temporal behavior and quantitative properties are, however, sensitive to the global parameters of the model and the presence or absence of magnetic fields.

The angular momentum distribution at the initial state is given by the parameter S and specific function of the polar angle. The black hole rotation is initially parameterized by the dimensionless Kerr parameter A_0 , and then the black hole spin changes during the simulation. It may either decrease or increase, and the final spin is measured when the simulation ends (the end time of most simulations is fixed at $20,000 M^3$).

In Figure 1 we show the initial distributions of density in the cloud depending on the rotation parameter. The subcritical models do not contain enough angular momentum in the envelope to form a disk bubble at the equatorial plane. Hence, the distribution of density in the gas is spherically symmetric and resembles the Bondi flow. Material from both polar and equatorial regions contributes to the growth of black hole mass, as it cannot be kept by rotation on the orbit and will flow under the horizon. We note, however, that even if there is a small amount of angular momentum in the gas, there will be dissipation of energy at the equator, because of shocks and compression (see Murguía-Berthier et al. 2020). For critical rotation in the envelope, a rotationally supported torus may form at the equator, and material can accrete only through the polar regions.

Below we present results for two sets of models: nonmagnetized and magnetized. The summary of adopted model parameters and the results are given in Table 1.

3.2. Simulations with Nonmagnetized Accretion Flows

In Figure 2 we present the time evolution of the main important quantity, the mass accretion rate onto the black hole, for all nine models without magnetic field. The models shown

³ Gravitational times, $t_g = GM/c^3$, in units of black hole mass. For the black hole of $3 M_\odot$ the physical time of our simulation would be ~ 1.5 s.

Table 1
Summary of the Models

Model	Black Hole Spin A_0	S l/l_{crit}	β_0	t_i (s)	M_{BH}^i	A_f	$\dot{M}_{\text{max}}^i$ ($M_\odot \text{ s}^{-1}$)
HS-T-04	0.85	0.4	∞	2.1	15.45	0.15	45.1
HS-LM-04	0.85	0.4	100	1.4	17.35	0.13	49.6
HS-HM-04	0.85	0.4	1.0	1.3	17.30	0.13	49.6
LS-T-04	0.5	0.4	∞	1.9	16.09	0.16	46.4
LS-LM-04	0.5	0.4	100	1.4	18.07	0.13	51.0
LS-HM-04	0.5	0.4	1.0	1.4	18.07	0.13	51.1
NS-T-04	0.3	0.4	∞	2.0	16.22	0.17	46.7
NS-LM-04	0.3	0.4	100	1.4	18.20	0.14	51.4
NS-HM-04	0.3	0.4	1.0	1.4	18.20	0.14	51.4
HS-T-10	0.85	1.0	∞	4.5	14.30	0.40	41.1
HS-LM-10	0.85	1.0	100	1.4	15.90	0.33	47.7
HS-HM-10	0.85	1.0	100	1.4	15.90	0.34	43.6
LS-T-10	0.5	1.0	∞	2.4	14.95	0.43	42.6
LS-LM-10	0.5	1.0	100	1.3	17.65	0.34	49.7
LS-HM-10	0.5	1.0	1.0	1.5	16.90	0.35	46.2
NS-T-10	0.3	1.0	∞	2.8	12.22	0.57	34.1
NS-LM-10	0.3	1.0	100	1.2	17.60	0.36	49.7
NS-HM-10	0.3	1.0	1.0	1.2	17.14	0.37	45.8
HS-T-14	0.85	1.4	∞	1.5	11.90	0.65	42.0
HS-LM-14	0.85	1.4	1.0	1.4	15.70	0.50	44.1
HS-HM-14	0.85	1.4	1.0	1.5	15.20	0.51	43.2
LS-T-14	0.5	1.4	∞	1.4	12.43	0.71	48.3
LS-LM-14	0.5	1.4	1.0	1.3	16.55	0.53	45.7
LS-HM-14	0.5	1.4	1.0	1.3	16.05	0.52	42.7
NS-T-14	0.3	1.4	∞	1.4	12.60	0.74	79.3
NS-LM-14	0.3	1.4	1.0	1.2	16.50	0.55	45.9
NS-HM-14	0.3	1.4	1.0	1.2	16.20	0.57	43.9

Note. The HS models refer to a highly spinning black hole at birth, the LS models represent a moderately spinning black hole, and NS models are for a slowly spinning black hole at birth. The angular momentum magnitude, S , and the spin of the initial black hole are given as the initial state parameters. In magnetized models, denoted with *HM*, the plasma β was adopted to 1.0, *LM* models correspond to $\beta = 100$, while in thermal models, denoted with *T*, the magnetic field was neglected, $\beta = \infty$. The initial mass of the accreting cloud was set to $25 M_\odot$, while the final mass reached at time t_f was close to zero. The last three columns give the resulting final black hole mass, its final spin, and time-averaged mass accretion rate.

differ with respect to rotation magnitude in the accretion flow and the initial black hole spin. Our fiducial models were run until the end time of $t = 1.5$ – 4.5 s.

We observe that the accretion rate onto a black hole is small initially but then rises, and the steepest rise is always found in

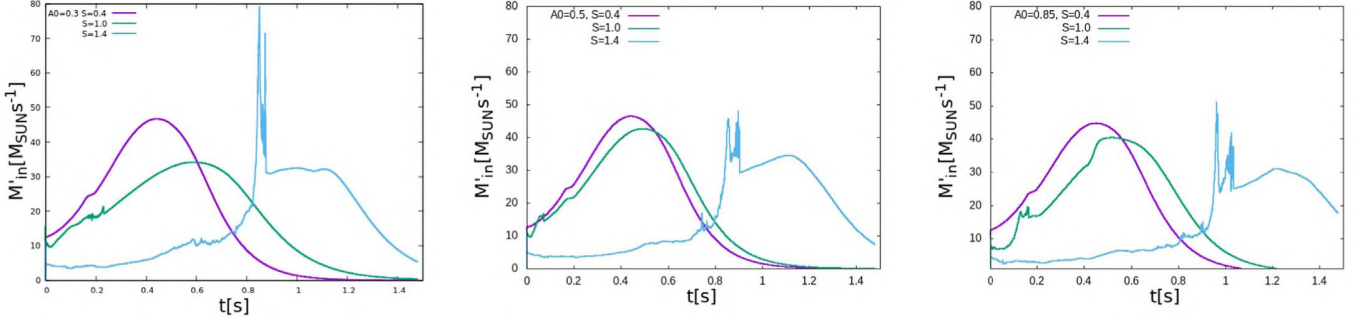


Figure 2. Time dependence of the mass flux through the inner boundary for three different initial black hole spins. Left panel: $A_0 = 0.3$; middle panel: $A_0 = 0.5$; right panel: $A_0 = 0.85$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines $S = 1.0$.

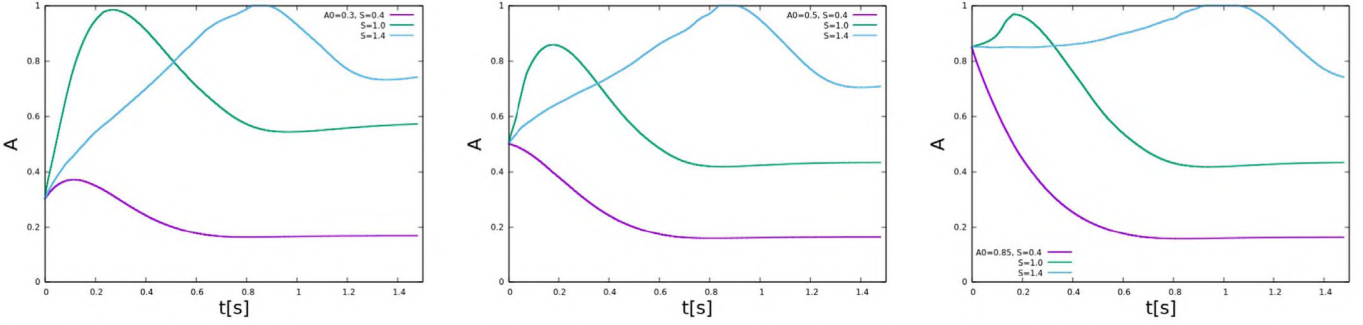


Figure 3. Time dependence of the black hole spin for three different initial black hole spins. Left panel: $A_0 = 0.3$; middle panel: $A_0 = 0.5$; right panel: $A_0 = 0.85$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines $S = 1.0$.

the subcritical rotation models ($S = 0.4$). The critical and supercritical rotation of the flow results in a somewhat slower rise of the accretion rate and a smaller value at the peak. This is because a large part of the material is kept by rotation in the centrifugally supported mini-disk, while only the flow accreting from the polar regions contributes to the mass accretion rate through the inner radius, $\dot{M}_{in}$. Moreover, for critical rotation some fluctuations of the accretion rate appear with low amplitudes, at the initial times of the simulation. They appear for all simulations with $S = 1.0$, but at different times, depending on the initial value of the spin parameter. For supercritical rotation, $\dot{M}$ at the beginning is steady and low in comparison to other cases. Small-amplitude fluctuations develop later in time (from ~ 0.6 to ~ 0.8 s). Also, for $S = 1.4$ models, the accretion rate oscillations with very large amplitude appear transiently, before the $\dot{M}$ becomes roughly constant for a period of time and then starts decreasing. In this case the moment of appearance of this feature is weakly sensitive to the value of the initial black hole spin. The largest amplitude of oscillations is obtained for the model with small initial spin, $A_0 = 0.3$.

Evolution of the black hole spin is presented in Figure 3. For subcritical rotation only in the case of $A_0 = 0.3$ do we observe an increase of the black hole spin at any stage of the simulation, when it grows to the value of ~ 0.37 . Then, it starts to decrease and finally saturates at the value of $A \sim 0.17$. In the case of other initial spin parameters, the value of black hole spin starts to decrease from the very beginning of the simulation and saturates at a similar value. In the case of critical rotation of the envelope there is no simple correspondence between the initial and maximal values of spin parameters. The highest maximal spin is achieved in the case of $A_0 = 0.3$ ($A_{max} \sim 0.98$); intermediate initial spin, $A_0 \sim 0.5$, gives the lowest value of

maximal spin ($A_{max} \sim 0.87$); and the highest initial spin, $A_0 \sim 0.85$, gives a maximal spin value during evolution at the level of $A_{max} \sim 0.92$. The final spin value is similar for every critical case with values from ~ 0.4 to ~ 0.57 . In the case of supercritical rotation models it is worth comparing spin evolution to the accretion rate evolution described above. For all A_0 values, the spin is growing since the beginning of the simulation and achieves the maximal value ~ 1 when the high-amplitude oscillation of accretion rate appears. Spin maintains its maximal value during those oscillations. Then, it starts to drop to the final value, which is around ~ 0.7 for all supercritical models. Exact values can be found in Table 1.

We conclude that the main factor that determines evolution of the spin and accretion rate during the simulation is the angular momentum content in the envelope, i.e., the value of the S parameter. Models are further dependent on the A_0 value, and it seems to shift the moment in time at which a characteristic feature arises and also changes its duration and the amplitude. Nevertheless, the occurrence of those characteristic features is preserved.

In Figure 4 we present the evolution of the black hole mass during simulations. For every initial spin value the highest black hole mass is achieved for $S = 0.4$. Initial spin $A_0 = 0.3$ is the only case when the supercritical rotation is more efficient for mass gain than critical. For the initial part of the simulation in every case the black hole mass increases linearly. In the case of supercritical rotation the slope of $M_{BH}(t)$ changes, which corresponds to the beginning of high-amplitude oscillations of accretion rate.

The 2D structure of accretion flow in our simulations is as follows.

In Figure 5 we show the 2D distributions of density, angular momentum, and Mach number at the end of simulation of the

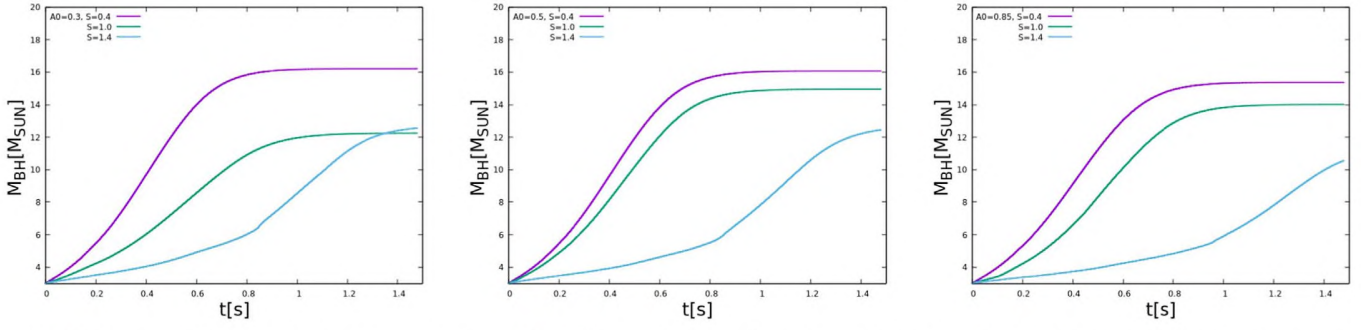


Figure 4. Time dependence of the black hole mass for three different initial black hole spins. Left panel: $A_0 = 0.3$; middle panel: $A_0 = 0.5$; right panel: $A_0 = 0.85$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines $S = 1.0$.

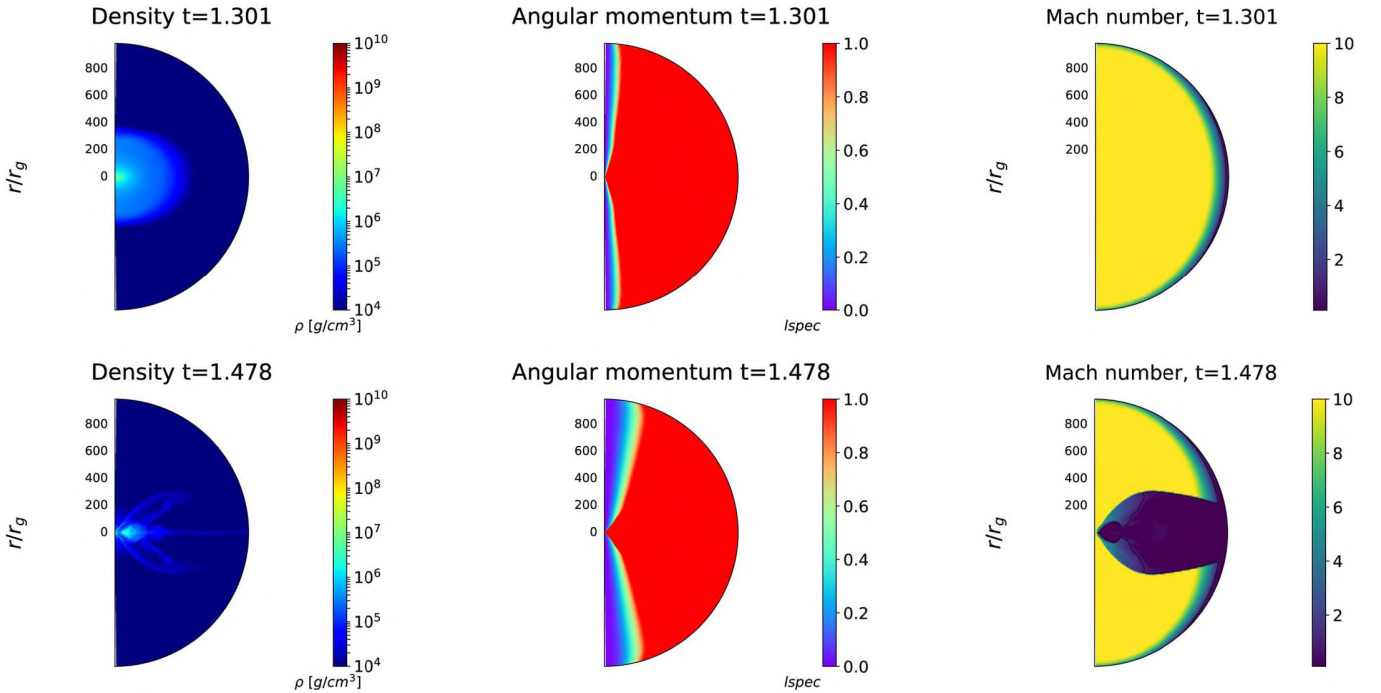


Figure 5. Results, from left to right, for the density, specific angular momentum, and Mach number distributions, for one of the models with neglected magnetic fields. The initial black hole spin was $A_0 = 0.3$, and the rotation parameter was $S = 0.4$ (NS-T-04 model, nonmagnetized). The color maps are taken in the end of the simulation. Note that in the middle panels the spatial scale of the plot is zoomed in to $100 r_g$. In addition, the contour of $M = 1$ is shown with a black line.

most slowly rotating flow with $S = 0.4$. The model shows the case of initial black hole spin $A_0 = 0.3$ (model NS-T-04). Rotation in the envelope was below critical, so the structure of the flow was close to spherical, for most of the time. A somewhat flattened, Bondi-like structure accreted onto the black hole uniformly, and the envelope became almost empty by the time $t \sim 1.5$ s. The mass of the black hole reached about $16 M_\odot$, while its spin has dropped down to about $A = 0.17$, because of accretion of very low angular momentum material (see above). An interesting feature (a “butterfly shape”) developed in the equatorial plane of the envelope at the end of the collapse, after time $t \sim 1.35$ s, which we show in Figure 5.

The fiducial model (NS-T-04) was run until time $t = 1.4$. In addition, we ran a longer simulation to verify that the “butterfly shape” feature develops into a disk-like structure at the equatorial plane, which is sustained for up to 2 s. The duration of this feature additionally depends on the A_0 value.

In the equatorial plane, a remnant material with relatively high angular momentum reached the innermost regions of the collapsar and is kept by rotation. This mini-bubble with subsonic accretion speed survived at the equator for a long time, preventing the cloud from getting completely evacuated. The supersonic matter overpasses the bubble at higher latitudes and sinks under the black hole horizon.

Evolution of the density structure in the collapsing star corresponds to features visible in $\dot{M}$ behavior. As mentioned above, they are determined mostly by the S -parameter value. Simulations for the same rotation magnitude S and different initial spin A_0 give similar results, and characteristic moments are just shifted in time.

As an example of a simulation with a critical initial angular momentum, $S = 1.0$, in the envelope, let us discuss the case with the highest initial spin value, $A_0 = 0.85$. At the beginning of the calculations we observe accumulation of matter in the innermost part of the cloud, forming also multiple sonic surfaces in the flow. It corresponds to slower growth of the $\dot{M}$

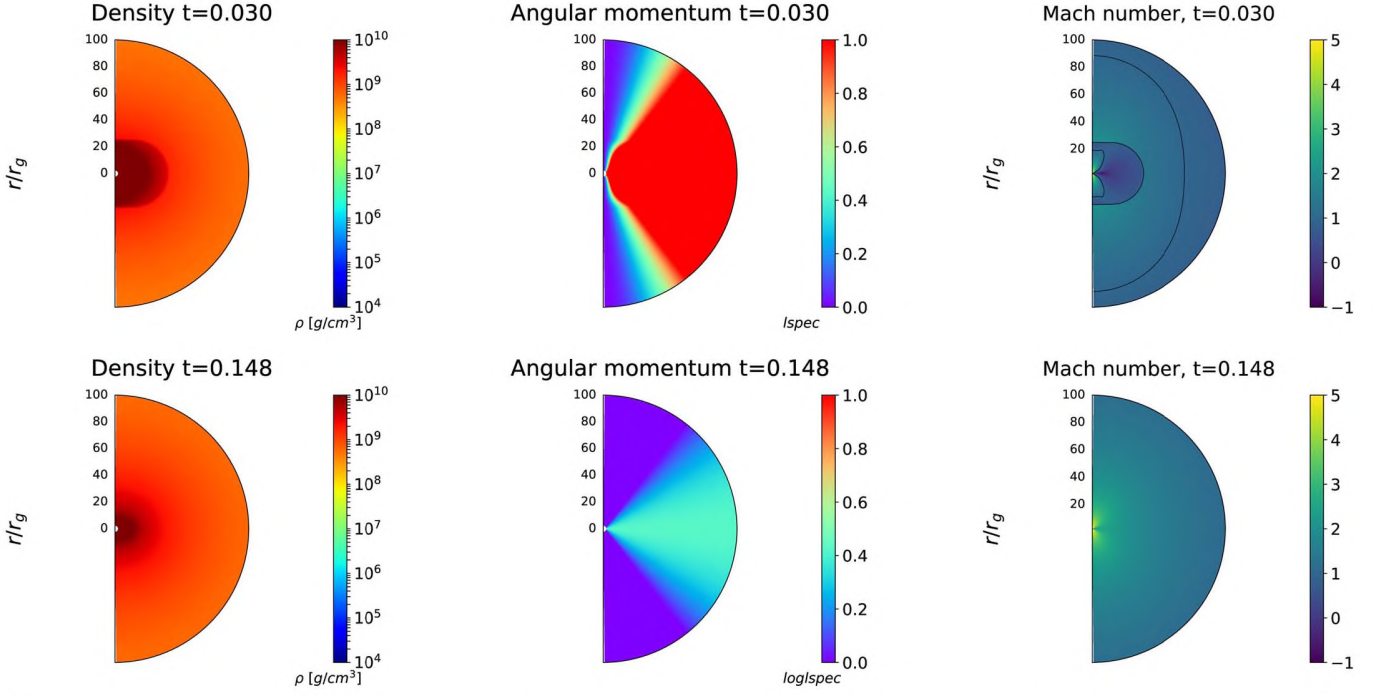


Figure 6. Results, from left to right, for the density, specific angular momentum, and Mach number distributions, for the model with neglected magnetic fields and initial black hole spin of $A_0 = 0.85$ and the rotation parameter $S = 1.0$ (model HS-T-10). The color maps show profiles at the beginning of the simulation.

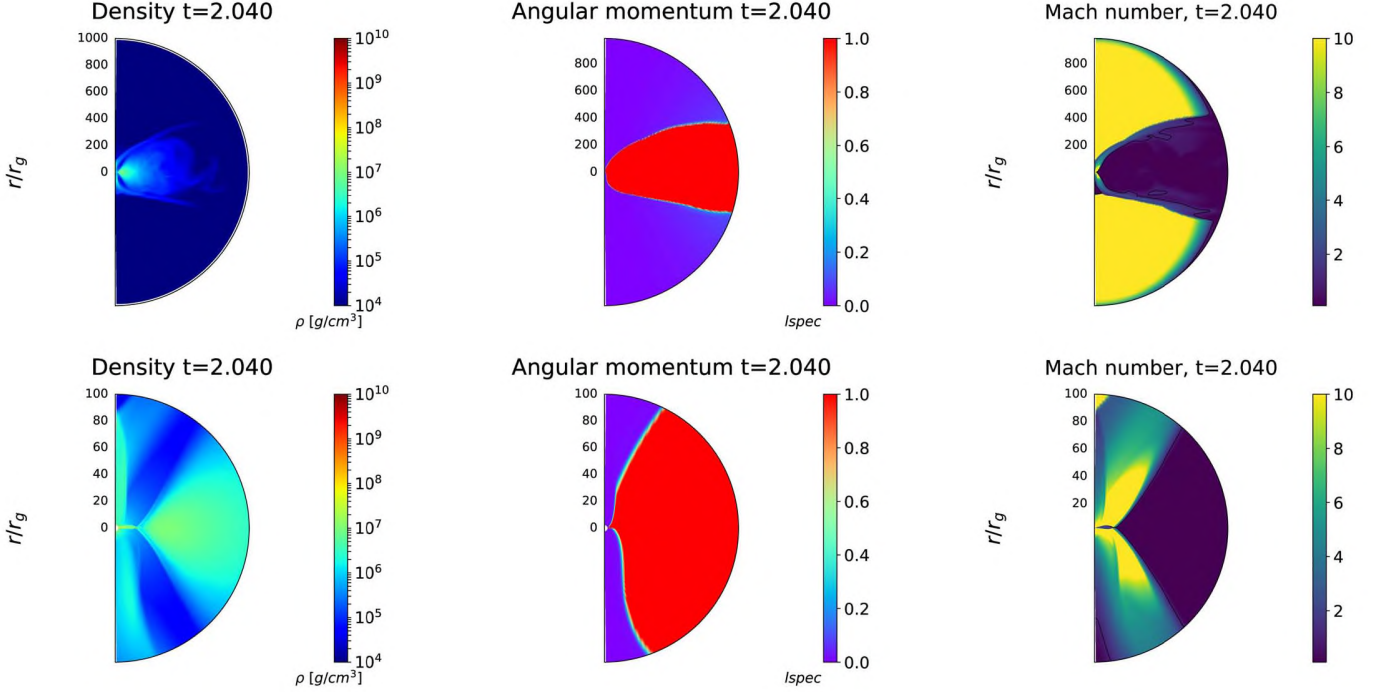


Figure 7. Results, from left to right, for the density, specific angular momentum, and Mach number distributions, for the model with neglected magnetic fields and initial black hole spin of $A_0 = 0.85$ and the rotation parameter $S = 1.0$ (model HS-T-10). The color maps show the disk structure that appears in this simulation and is sustained for ~ 4 s. The contour of $M = 1$ is marked with a black line. (Note the different spatial scale of density and angular momentum profiles in the first and second rows.)

than in the case of subcritical rotation. Subsequently, matter from denser inner regions gets accreted onto the black hole (reflected in the growth of the accretion rate), and multiple sonic surfaces no longer exist. There is only one such surface. Density and angular momentum profiles corresponding to this initial stage of the simulation are shown in Figure 6. Later, the

shape of the cloud is mostly spherical. At the final times, around $t = 1.5$ s the density distribution flattened toward the equator, creating a mini-disk structure. The additional run of this simulation, carried for a longer time, revealed the formation of the elongated disk structure. In this model, the disk was present for the longest time of ~ 3 s. In the case of

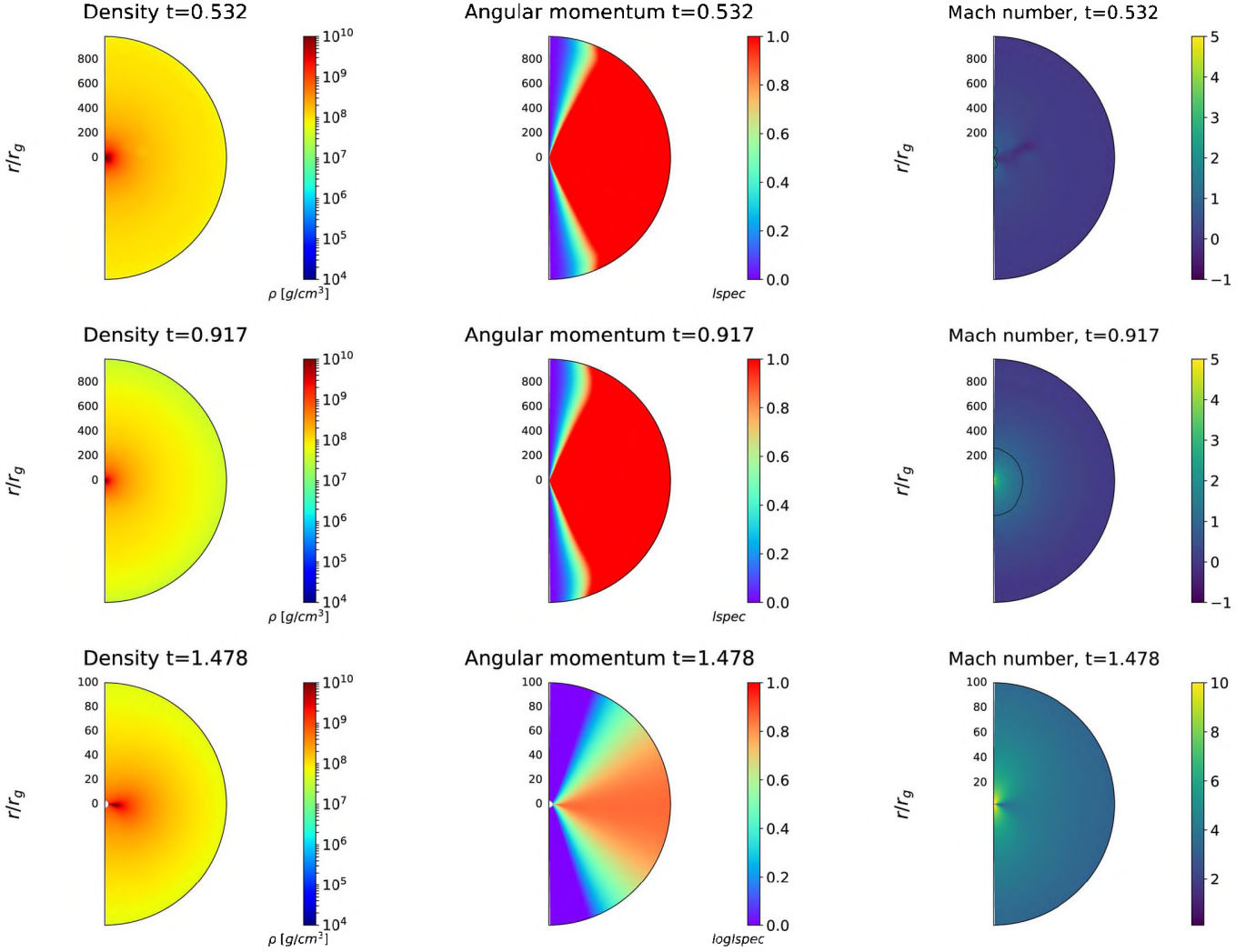


Figure 8. Results, from left to right, for the density, specific angular momentum, and Mach number distributions, for the model with neglected magnetic fields and initial black hole spin of $A_0 = 0.3$. The rotation parameter was $S = 1.4$ (model NS-T-14), and there was no magnetic field. The color maps are taken at two intermediate times, $t = 0.53$ and $t = 0.92$, and at $t = 1.48$, which is at the end of the simulation. In addition, the contour of $M = 1$ is shown with a black line. (Note the different spatial scales of density and angular momentum profiles in the third row.)

different values of initial spins, such structure did not sustain longer than for ~ 1 s. In Figure 7 we show the elongated disk structure that was formed and sustained.

In Figure 8 we present selected profiles of density, specific angular momentum, and radial Mach number (M_r) distributions, for the model with $A_0 = 0.3$ and supercritical rotation, $S = 1.4$. Selected snapshots correspond to different stages of evolution of the accretion rate that is reflected in the flow hydrodynamical structure. The first row of Figure 8 shows the snapshots taken at $t = 0.532$ s. It corresponds to a steady $\dot{M}$ value. The sonic surface is very close to the black hole at the equator and forms an “eight-shape” structure, elongated in the polar direction. Close to the equatorial plane a drop in the M_r value is visible. The middle row shows a snapshot at $t = 0.917$ s. It corresponds to the high-amplitude oscillations in the accretion rate. Density profiles around that time flicker. The sonic surface becomes spherical, and with time it starts to move to the outer boundary of the mesh. The last row shows the final structure of the flow at the end of the simulation. In supercritical simulations with $S = 1.4$, we do see only a mini-disk structure at the innermost region of the flow, up to

$r \sim 10r_g$ at time $t = 1.5$ s. At larger scales, the density profile is spherically symmetric. Now the angular momentum is largest close to the black hole, while the density of the entire cloud has dropped significantly in the whole envelope, leaving only a small centrifugally supported dense remnant disk at the equator below $\sim 20r_g$ (note different spatial scales of density and angular momentum maps in the last row in Figure 8, on which a small-scale disk structure is visible).

To summarize the parameter space of our models, in Figure 9 we present color maps showing the results of final and maximal global quantities obtained for our grid of models, plotted with the parameter-space contours. The resulting black hole mass and final spin, obtained at the end of the simulation, are depicted with a color scale (bottom left and right panels, respectively). We also show the maximal spin and maximal accretion rate values reached during the simulation (top left and right panels, respectively).

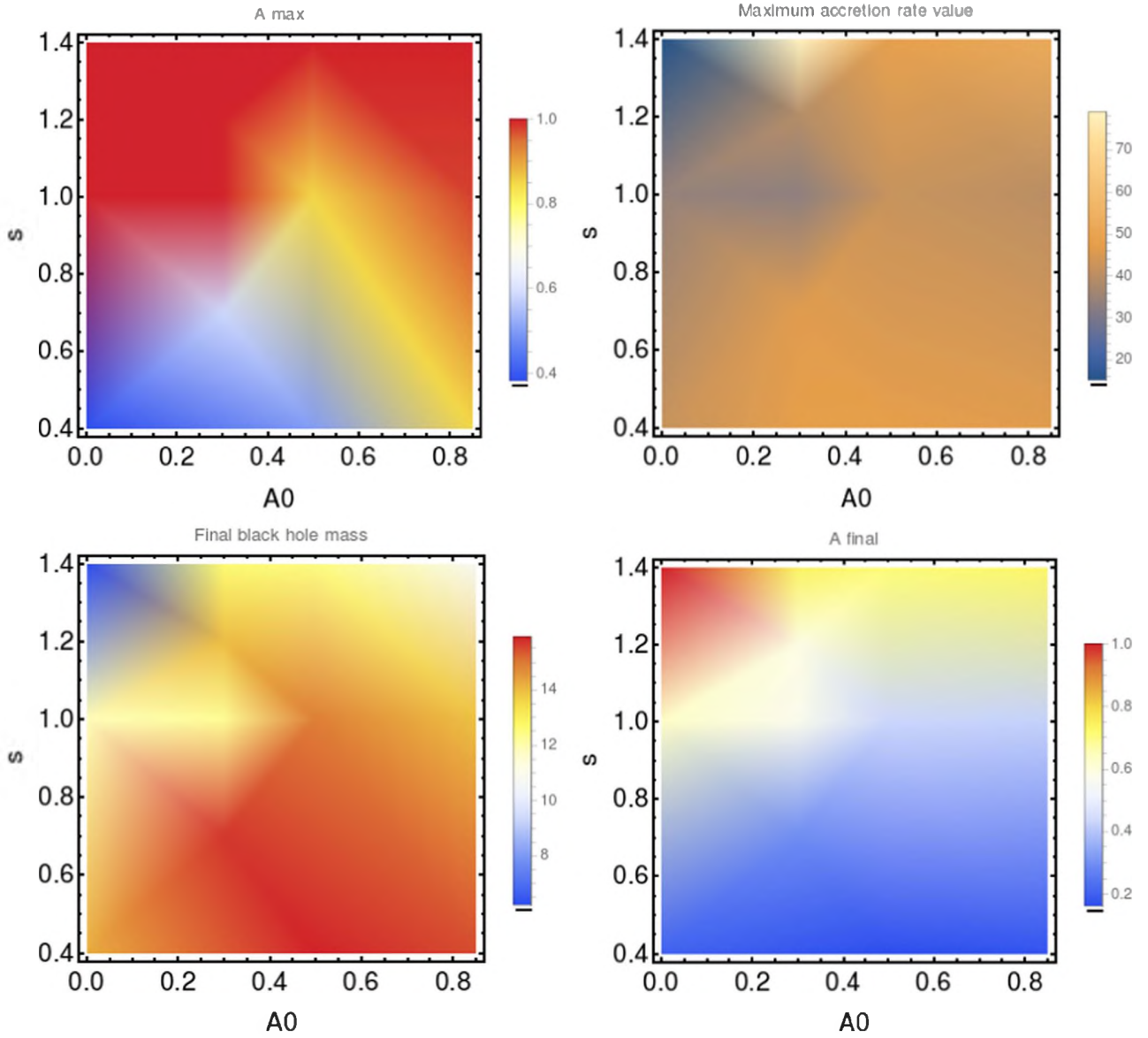


Figure 9. Parameter space studied in our simulations, and results for maximal spin during the simulation (top left panel), maximal accretion rate (top right panel), and final values of mass (bottom left panel) and black hole spin (bottom right panel). The maps show results with nonmagnetized flows.

3.2.1. Shocks

At critical rotation, the shock discontinuity forms and pressure builds up to slow down the incoming material near the equator. For an even higher rotation parameter, $S = 1.4$, the accretion flow forms a geometrically thin structure at the equatorial region. Here, a transient shock structure may form to redistribute the angular momentum, but the shock is still subsonic for $S < 2$. The shock for the case $A_0 = 0.3$ and $s = 1.4$ appears only for $t = 0.857$ s at $r = 16.37r_g$.

We investigated the Mach number profiles in order to check for the shock's existence. Our algorithm for shock finding was simple, i.e., we were looking for the radii where the flow changes from supersonic to subsonic. Additionally, we checked whether this change in the flow is connected with a steep change in the thermodynamic properties of the flow. We rejected the apparent shocks that appear with a supersonic velocity value only in one (the outermost) grid cell and are

numerical artifacts. We investigated the shock presence in the equatorial plane and found that for almost every combination of A_0 and S parameters a shock appeared at some point. The only exception was the model with $A_0 = 0.3$ and $S = 1.0$, in which a shock did not appear at any stage of the simulation. For $A_0 = 0.3$ and $S = 1.4$ a shock appears only in one moment, at time 0.857, so that it lasts less than 0.03 s. For $A_0 = 0.3$ and $S = 0.4$ shocks start to appear at the final stages of the simulation, and they are placed at radii around $100 - 500r_g$.

In the case of $A_0 = 0.5$, for $S = 0.4$ the shock appears at $r \sim 100r_g$ in the end of the simulation; for $S = 1.0$ it can be found at a similar time but closer to the black hole, at $r \sim 20r_g$. In addition, a transient shock is present at the beginning of the simulations below $50r_g$. In the case of supercritical rotation, the shock was found below $50r_g$ at intermediate times.

In the case of high initial black hole spin, the situation is similar. For $S = 0.4$ the shock appears at the end of the simulation at high r , in the case of critical rotation the shock

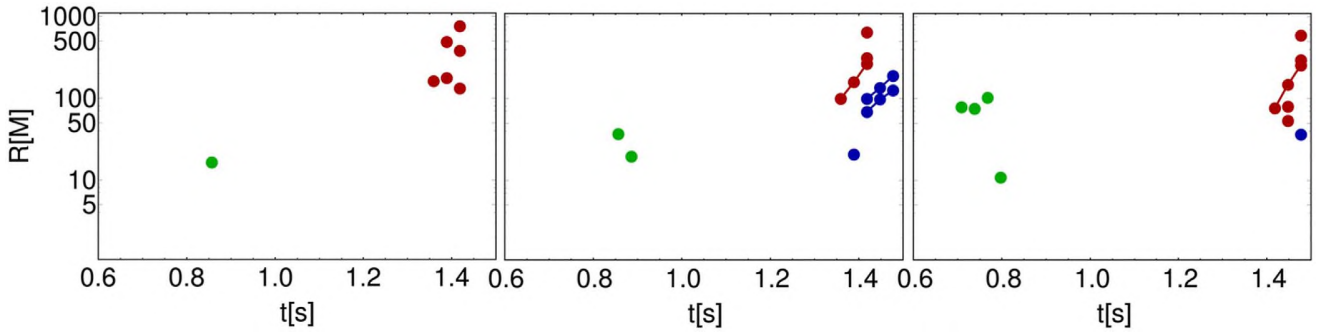


Figure 10. Shock positions as a function of time. Red circles: $S = 0.4$; blue circles: $S = 1.0$; green circles: $S = 1.4$. Panels from left to right present models with various initial spin of the black hole: $A_0 = 0.3$, $A_0 = 0.5$, and $A_0 = 0.85$. Note that plots present time from 0.8 s to the end of the simulation, so shocks that appear as initial condition relaxation are not visible. They do appear for critical initial rotation for $A_0 = 0.85$ and $A_0 = 0.5$ before 0.1 s at $\sim 20r_g$ and $\sim 50r_g$.

appears at the very beginning and at the end of the simulations, and for supercritical rotation it appears for a moment at an intermediate time.

Shock positions evolving in time are shown in Figure 10. In simulations with $A_0 = 0.5$ and $S = 0.4$, $A_0 = 0.5$ and $S = 1.0$, and $A_0 = 0.85$ and $S = 0.4$ (models LS-T-04, LS-T-10, and HS-T-04, respectively) the shocks that appear at later stages are standing over some time and then expand to the outer boundary of the grid. We calculated their velocities, which are given in Table 2. We compare them to the escape velocity, corresponding with the position of the shock. Because this ratio is always below 20%, we conclude that such shocks are too slow to be able to dissipate enough energy and disrupt the outer envelope.

3.3. Simulations with Magnetized Accretion Flows

In this section, we present the simulations of magnetized stellar envelopes, where the low angular momentum flow is embedded in a poloidal magnetic field. The specific pattern of the evolution depends on the β value, but also on the combination of A_0 and the S parameters. We performed two sets of simulations for the same combinations of A_0 and S as for the case without magnetic field discussed above. Tested values of β were 1 and 100. We present the evolution of the global parameters for both β values in Figures 11–13. The evolution of black hole mass and spin depends strongly on the rotation of the envelope, parameterized by S .

For $S = 0.4$, evolution is similar to nonmagnetized models, and subcritical models with $\beta = 1$ and $\beta = 100$ are nearly identical; moreover, they behave in a similar way for every value of the initial black hole spin. The accretion rate smoothly grows to the value of $\sim 50 M_\odot \text{ s}^{-1}$ (Figures 11–13; see purple solid and dashed lines), and then it starts to decrease. Spin increases at any stage of the simulations only for $A_0 = 0.3$, and in this model the final black hole mass is the highest. It is also higher than in the case of a nonmagnetized envelope.

In the case of critical and supercritical rotation, the evolution pattern for various initial parameters is more complicated. First, the β value affects the appearance of the oscillations in the accretion rate. For a less magnetized envelope these oscillations do not appear for critical rotation, and both accretion rate and black hole mass evolution look similar to the subcritical case. The evolution is smooth; accretion rate grows to the values below $\sim 50 M_\odot \text{ s}^{-1}$, which are slightly smaller than values reached in the case of subcritical rotation; and the shape of the evolution is very similar. The achieved final black hole masses

Table 2
Summary of the Models

A_0	S l/l_{crit}	v c	v_e c
0.5	0.4	0.041	0.26
0.5	1.0	0.014	0.34
0.5	1.0	0.022	0.28
0.85	0.4	0.045	0.28

Note. Velocities of shock fronts and escape velocities for different models.

are somehow smaller than in the subcritical case. However, the difference is notable in the case of the spin. Critical initial rotation of the envelope for every value of A_0 results in maximal spins achieved during simulation, which are significantly larger than the initial ones.

For supercritical rotation of the envelope, the oscillations of the accretion rate appear around 0.4 s for every A_0 parameter. The accretion rate growth is steepest for $A = 0.3$. For every A_0 value, the oscillation is not related to the $\dot{M}$ maximum value; it occurs later during the simulation, and the peak is at about $\sim 45 M_\odot \text{ s}^{-1}$. Final black hole masses are lower than in critical and subcritical cases, but higher for $\beta = 100$ than in the corresponding model with $\beta = 1$. Also, the maximal value of the spin is higher than for the stronger magnetization, but still in none of the cases is maximal spin reached, and spins are generally smaller than in the case of a nonmagnetized envelope.

For $\beta = 1$, in both critically and supercritically rotating envelopes during the growth of $\dot{M}$, we can observe oscillations of its value. After the peak of $\dot{M}$, the accretion rate smoothly decreases. In the case of the model parameterized with $A_0 = 0.85$ and $S = 1.4$, the accretion rate evolution is different than for all the other models. After initial oscillations of its value, we observe a large spike at the time ~ 0.5 s. Then, it drops to the previous level, saturates on it for some time, grows smoothly, and then starts to decrease. The final black hole masses for $\beta = 1$ are always higher than in the case of nonmagnetized accretion (see Table 1), but somewhat smaller than in the case of $\beta = 100$. In particular, for the initial parameters $A_0 = 0.85$ and $S = 1.4$, we get $M_{\text{BH}} = 15.2 M_\odot$, compared to $11.9 M_\odot$ for the nonmagnetized model and $15.70 M_\odot$ for $\beta = 100$. Final black hole spins, on the other hand, are on average lower, e.g., for $A_0 = 0.85$ and $S = 1.4$ we have $A_{\text{final}} = 0.51$ and 0.65 for HS-M-14 (magnetized) and HS-T-14 (nonmagnetized) models, respectively. Plots illustrating the evolution of $\dot{M}_{\text{in}}$, A , and M_{BH} are shown in Figures 11–13.

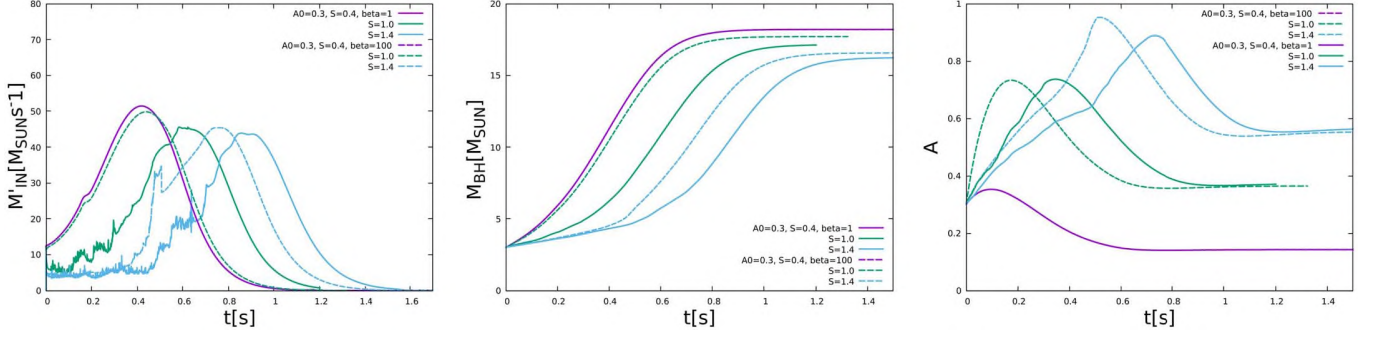


Figure 11. Time dependence of the mass accretion rate onto the black hole (left panel), black hole mass (middle panel), and its spin evolution (right panel) for the magnetized collapse case with $\beta = 100$ (dotted line) and $\beta = 1$ (solid lines). The initial black hole spin was $A_0 = 0.3$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines are for $S = 1.0$.

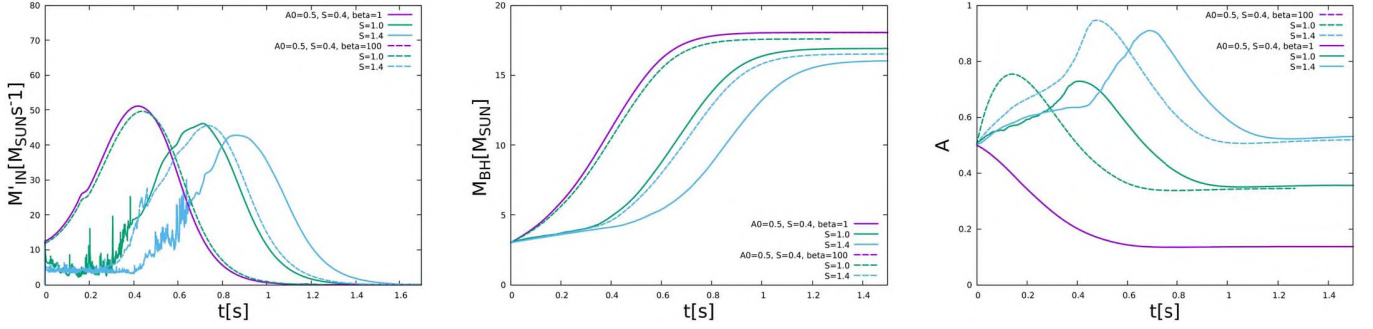


Figure 12. Time dependence of the mass accretion rate onto the black hole (left panel), black hole mass (middle panel), and its spin evolution (right panel), for the magnetized collapse case with $\beta = 100$ (dotted lines) and $\beta = 1$ (solid lines). The initial black hole spin was $A_0 = 0.5$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines are for $S = 1.0$.

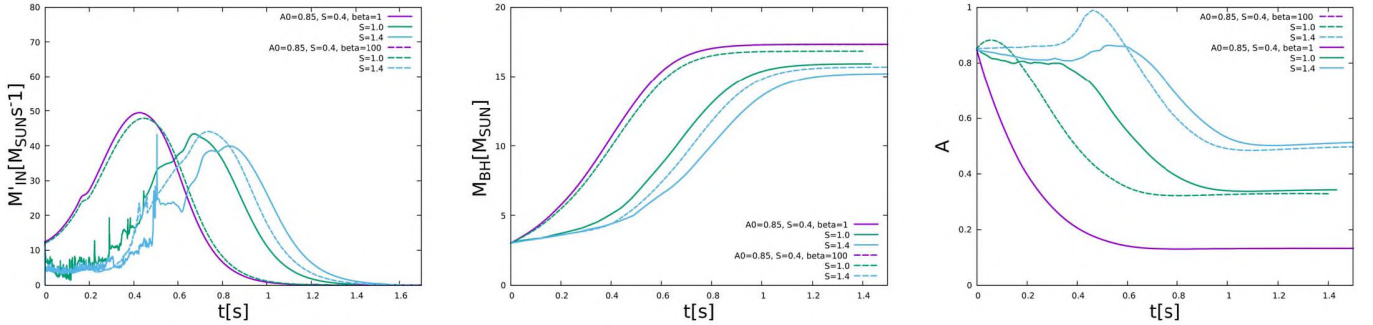


Figure 13. Time dependence of the mass accretion rate onto the black hole (left panel), black hole mass (middle panel), and its spin evolution (right panel), for the magnetized collapse case with $\beta = 100$ (dotted lines) and $\beta = 1$ (solid lines). The initial black hole spin was $A_0 = 0.85$. Cyan lines denote the models with $S = 1.4$, while the violet lines denote $S = 0.4$ and green lines are for $S = 1.0$.

Final and maximal values of those parameters are listed in Table 1.

3.3.1. Weakly Magnetized Envelope: $\beta = 100$

For model $A_0 = 0.3$ and $S = 0.4$, the density, radial Mach number, and angular momentum during the collapse evolve similarly to the case without magnetic field. They are spherically symmetric throughout most of the time, and at the end of the simulation they tend to flatten toward the equator. This behavior is the same for every value of the initial spin. Interestingly, this course of evolution is maintained in simulations of a nonmagnetized envelope and from the $\beta = 1$ case. The situation changes in the case of the models with $S = 1.4$. The evolution of the profiles is similar for every A_0 value. In the

first ~ 0.4 s there are some deviations from spherical symmetry. Contours of A_ϕ are disturbed at the central part of the grid. In the radial Mach number profile, the sonic surface has an eight-like shape. Those features start to change around $t = 0.4$ s, which corresponds to the oscillation in accretion rate. Sonic surface shape starts to vary, and it gets separated and splits into two parts: one of a small circular shape at the innermost part of the grid, and one at the scales of $\sim 200r_g$. At a later time, the inner contour gets accreted and the outer part gets spherically symmetric and moves toward the outer part of the grid. Contours of A_ϕ go back to the original configuration, and all the profiles are smoothed and spherically symmetric. This is an effect of evacuating most of the mass from the envelope.

We present selected profiles in Figure 14.

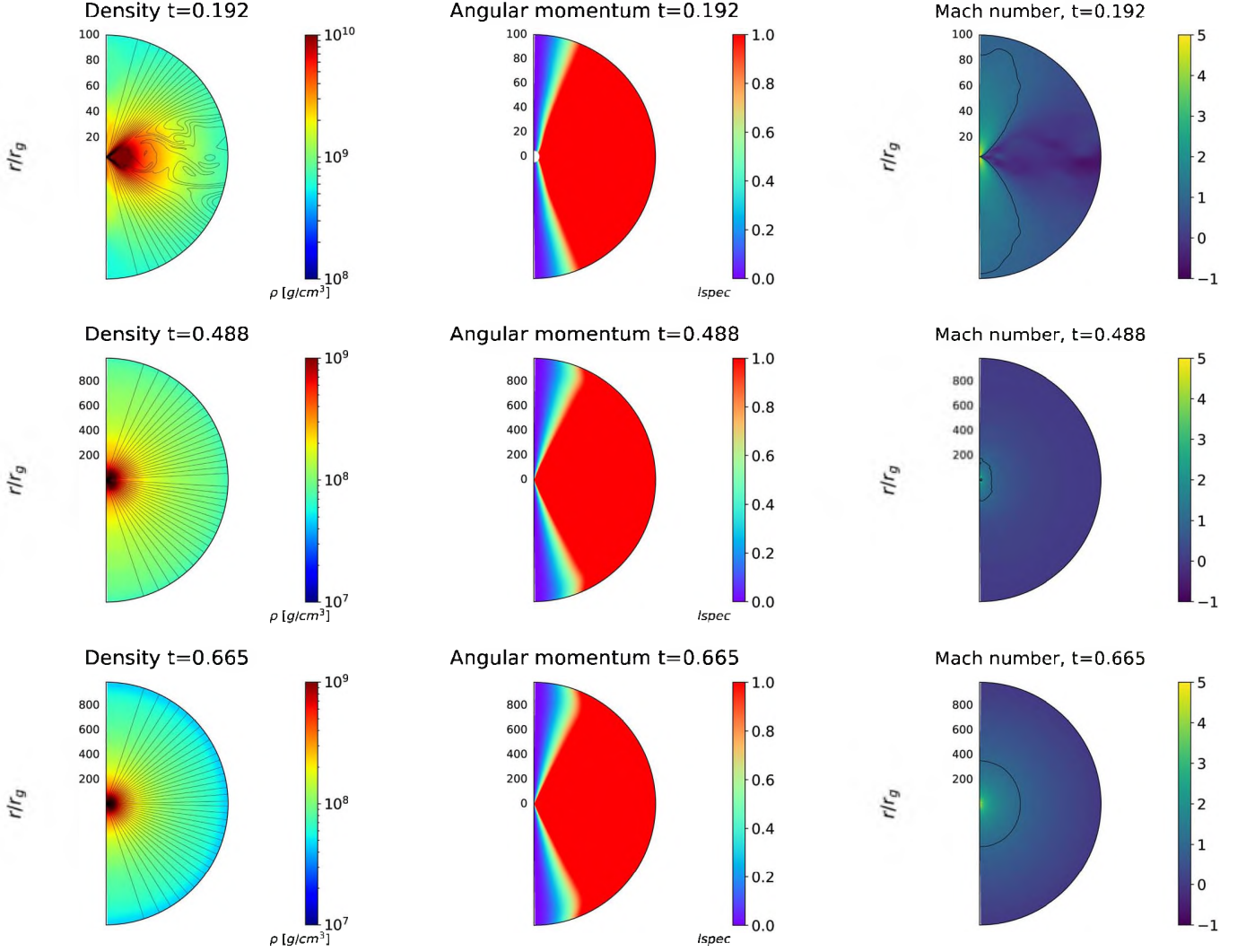


Figure 14. Results, from left to right, for the density with overplotted contours of magnetic field A_ϕ , specific angular momentum, and Mach number distributions, for the model with magnetic fields and initial black hole spin of $A_0 = 0.3$. The rotation parameter was $S = 1.4$ and $\beta = 100$ (NS-LM-14). The color maps are taken at times $t = 0.192$ s, $t = 0.488$ s, and $t = 0.665$ s, which correspond to the time before, right after, and after the spike in accretion rate. (Note that the first row shows a different spatial scale of $100 r_g$, for which also density scale is adjusted.) In addition, the contour of $M = 1$ is shown with a black line.

The presence of magnetic field and its strength influence the results in many different ways. In general, it affects the global parameters' evolution and their dependence on the initial conditions. In particular, we get the highest peak value of accretion rate for a nonmagnetized envelope with $A_0 = 0.3$ and $S = 1.4$ ($\sim 80 M_\odot s^{-1}$). The peak value appears during oscillations and does not persist, so it does not affect significantly the value of the final black hole mass. The presence of magnetic field modifies the shape of $\dot{M}$ evolution; the least affected case is subcritical rotation of the envelope, and for supercritical rotation we see substantial changes. Moreover, for magnetized envelopes we get significantly smaller maximal spins of the black hole. There is no magnetized model in which $A = 1$ is achieved. Spins in simulations with weaker magnetic field are higher than for $\beta = 1$. Interestingly, the highest black hole masses are reached for weaker magnetic field, and they are slightly higher than those in strongly magnetized envelopes, while the models without magnetic field result in the significantly smaller M_{BH}^f (see details in Table 1). Some particular sets of initial

parameters do not follow these general tendencies. They were described in previous sections. We attribute this result to the action of black hole rotation, which winds up the magnetic field lines and launches Poynting-dominated jets more efficiently, while the uncollimated wind outflows are decaying. The black hole spin does not affect the results directly, but its action is moderated by the magnetic fields.

3.3.2. Strongly Magnetized Envelope: $\beta=1$

In this subsection we discuss in more detail the distributions of density and magnetic field angular momentum and radial Mach number. We focus on the representative models in selected, characteristic snapshots representative for simulations, which correspond to changes in the $\dot{M}$ plot.

In Figure 15 we present selected profiles of density, Mach number, and specific angular momentum for simulations with $S = 0.4$ and $A_0 = 0.3$. The flow is magnetized, and magnetic field lines are overplotted on density maps. Here we show the models with the lowest rotation parameter. For most of the simulation time, the density distribution is spherically

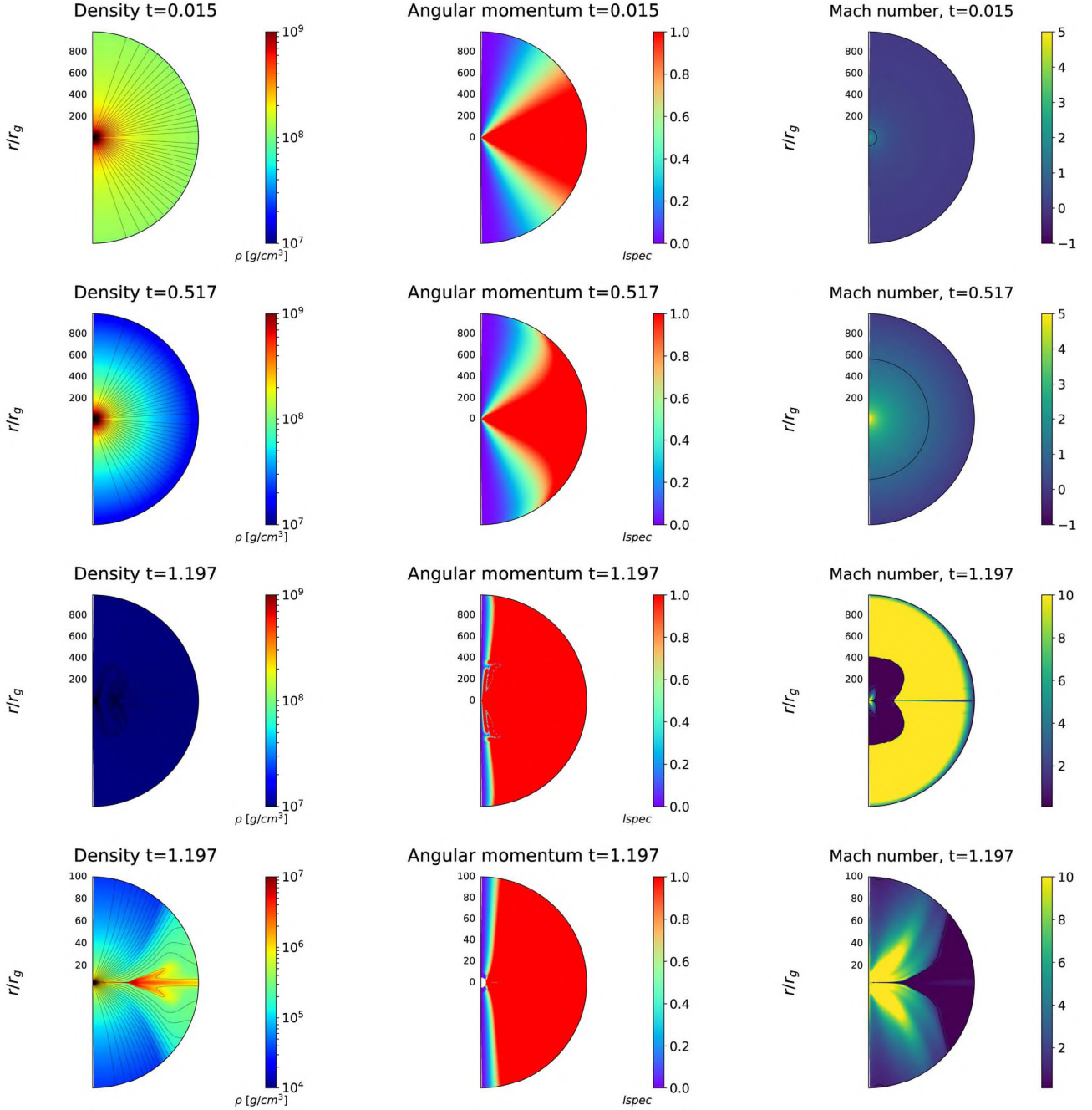


Figure 15. Results, from left to right, for the density, specific angular momentum, and Mach number distributions, for the model with magnetic field characterized by $\beta = 1$ and initial black hole spin of $A_0 = 0.3$. The rotation parameter was $S = 0.4$ (model NS-HM-04). The first row presents profiles at the beginning of the simulation. In the second row we show the profile from the intermediate stage of the simulations: $t = 0.517$. In the third and fourth rows we present the color maps taken at time $t = 1.197$, which corresponds to final stages of the simulation. In addition, the contour of $M = 1$ is shown with a black line.

symmetric, similarly to the $S = 0.4$ case, where the simulations were without magnetic field. The shape of the contours of A_ϕ stays radial and is unchanged during that time. The matter slowly accretes radially toward the black hole, the density in the whole domain uniformly decreases, and the flow becomes supersonic, as can be seen by comparison between first and second snapshots. The flow structure changes around the time 1.1 s. Now the density profile flattens toward the equator, and the shape of the A_ϕ contours gets disturbed. At the end of the simulation there is a structure created that is symmetrical about the equator and manifests a pileup of matter in the density

profile, with turbulent magnetic field depicted by the contours of A_ϕ . The radial Mach number profile shows that the flow here is highly subsonic, and the angular momentum increase means that it is rotating down to the black hole horizon. The last row in this figure shows the state of the simulation at time 1.19 s, zoomed in to the innermost part of the flow. Here we notice that the very thin region located at the equator is in fact rotating, while the material with small angular momentum and supersonic speeds can reach the black hole and disappear under the black hole horizon. At the same time, open magnetic field lines start forming, while they bend toward the polar axis. This

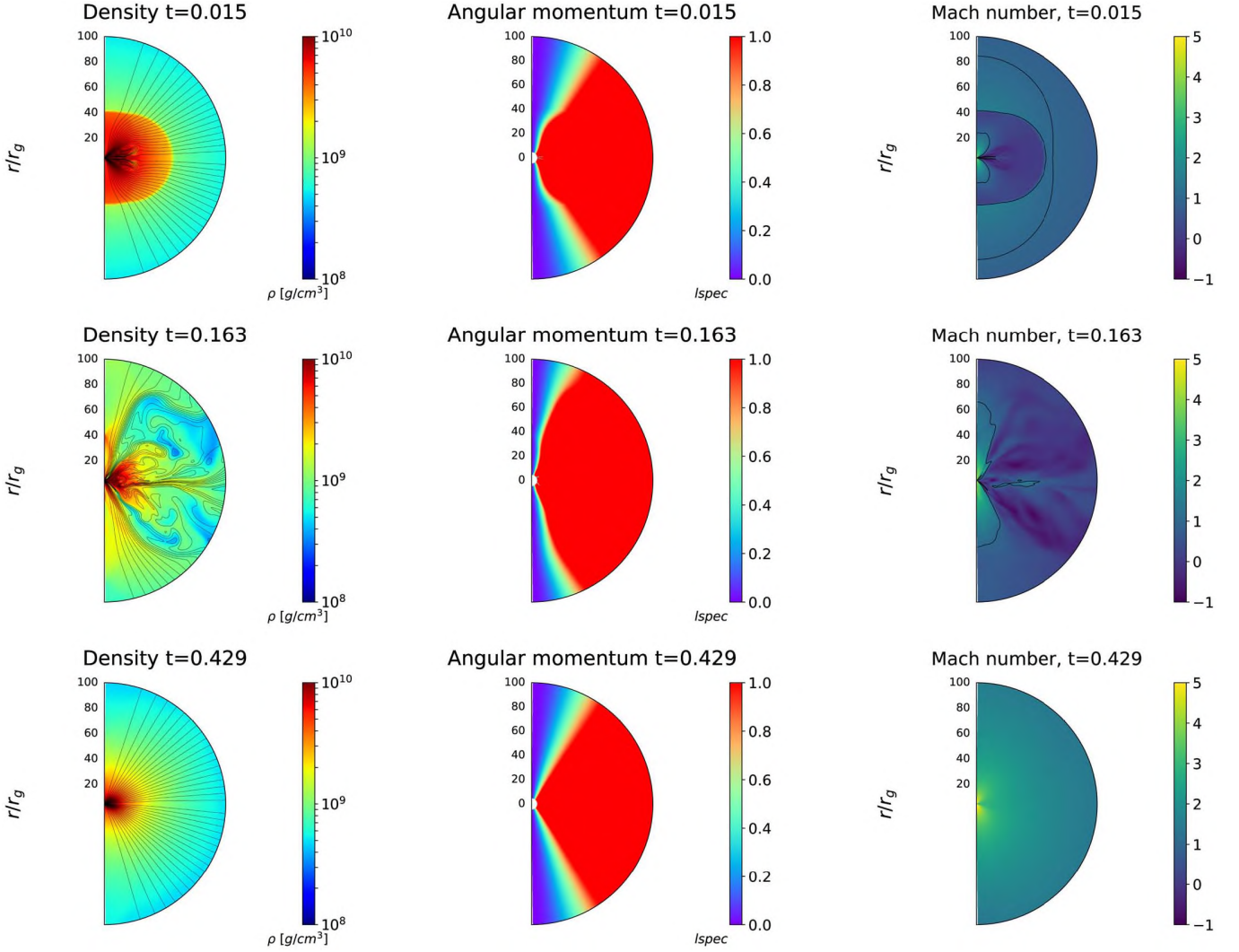


Figure 16. Results, from left to right, for the density with contours of A_ϕ , specific angular momentum, and Mach number distributions, for the model with magnetic field normalized by $\beta = 1$ and initial black hole spin of $A_0 = 0.5$. The rotation parameter was $S = 1.0$ (model LS-HM-10). The first row shows profiles at the beginning of the simulation; middle row profiles are taken at time $t = 0.163$ s, which corresponds to time of oscillations of $\dot{M}$; and the bottom panel shows profiles after the end of the oscillations. In addition, the contour of $M = 1$ is shown with a black line.

evacuated region can be envisaged as the place where the jet could be born in the collapsar’s envelope. The amount of angular momentum in this model is, however, not sufficient to sustain the jet for a very long time.

In Figures 16 and 17 we show the magnetized models with $S = 1.0$ and $S = 1.4$. First, we present profiles for the representative case: $A_0 = 0.5$ and $S = 1.0$. The magnetic field structure, depicted by contours of constant A_ϕ , is distracted at the time when oscillations of accretion rate value occur, while when they stop, the structure of magnetic field returns to its original shape (see Figure 16). As the Mach number profile shows, at the beginning of the simulations, the multiple sonic surfaces appear. Density and angular momentum profiles also follow those structures. Then, during the low-amplitude oscillations, the multiple sonic surfaces merge and form a single structure as is shown in the second row of Figure 16. The sonic point is very close to the black hole at the equator, while the shape of the sonic surface elongates with the latitude and gets disturbed, becoming an “eight-shape.” The magnetic field lines close to the polar region are bent and tend to an open field

line configuration. As mentioned above, this may suggest that the flow in the collapsing star will try to launch a magnetically supported jet along the evacuated low-density funnel. The angular momentum in the flow is, however, rather low everywhere, and only in localized places does the turbulence allow the matter to rotate faster than critical speed. Finally, after oscillations it becomes again spherically symmetrical and moves toward the outer boundary. Density and angular momentum profiles are smoothed, and as mentioned above, the shape of A_ϕ contours turns back to the initial state. Those profiles are shown in the bottom row of Figure 16.

The situation starts to change for critical initial rotation— $S = 1.0$. In Figure 15 we present selected profiles of density, Mach number, and specific angular momentum for simulations with $S = 0.4$ and $A_0 = 0.3$. The flow is magnetized, and magnetic field lines are overplotted on density maps. Here we show the models with lowest rotation parameter. For most of the simulation time, the density distribution is spherically symmetric, similarly to the $S = 0.4$ case, where the simulations were without magnetic field. The shape of the contours of A_ϕ

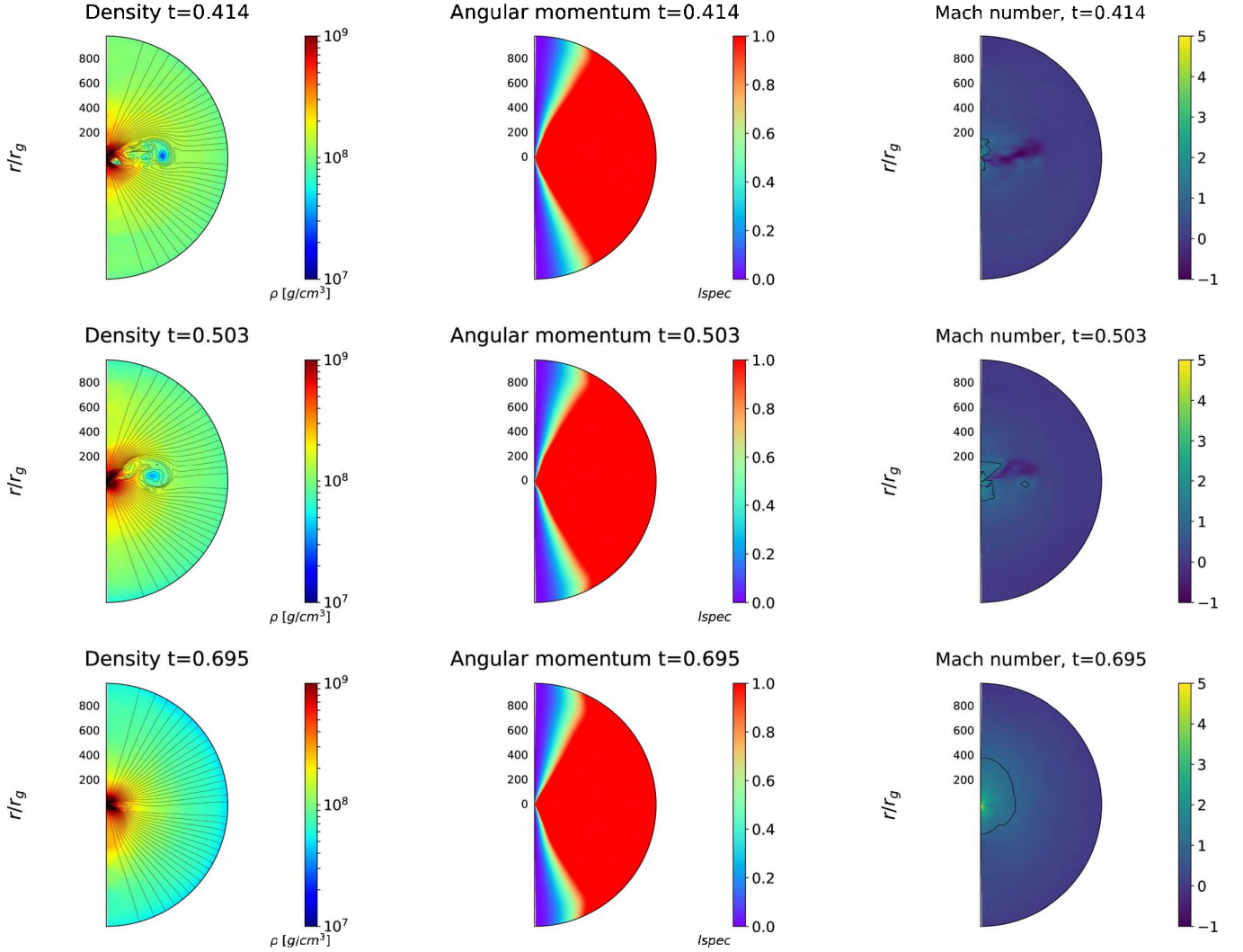


Figure 17. Results, from left to right, for the density with overplotted contours of magnetic field A_ϕ , specific angular momentum, and Mach number distributions, for the model with magnetic field characterized by $\beta = 1$ and initial black hole spin of $A_0 = 0.85$. The rotation parameter was $S = 1.4$ (model HS-HM-14). The color maps are taken at times $t = 0.414$, $t = 0.503$, and $t = 0.695$, which correspond to the time before, during, and after the spike, respectively. In addition, the contour of $M = 1$ is shown with a black line.

stays radial and is unchanged during that time. The matter slowly accretes radially toward the black hole, the density in the whole domain uniformly decreases, and the flow becomes supersonic, as can be seen by comparison between the first and second snapshots. The flow structure changes around the time 1.1 s. Now the density profile flattens toward the equator, and the shape of the A_ϕ contours gets disturbed. At the end of the simulation there is a structure created that is symmetrical about the equator and manifests a pileup of matter in the density profile, with turbulent magnetic field depicted by the contours of A_ϕ . The radial Mach number profile shows that the flow here is highly subsonic, and the angular momentum increase means that it is rotating down to the black hole horizon. The last row in this figure shows the state of the simulation at time 1.19 s, zoomed in to the innermost part of the flow. Here we notice that the very thin region located at the equator is in fact rotating, while the material with small angular momentum and supersonic speeds can reach the black hole and disappear under the black hole horizon. At the same time, open magnetic field lines start forming, while they bend toward the polar axis. This

evacuated region can be envisaged as the place where the jet could be born in the collapsar's envelope. The amount of angular momentum in this model is, however, not sufficient to sustain the jet for a very long time.

As mentioned before, for initial parameters $A_0 = 0.85$ and $S = 1.4$, the evolution of accretion rate is different from all the other models with $S = 1.4$ and $S = 1.0$. In Figure 17 we show the distribution of density with A_ϕ contours, angular momentum profiles, and radial Mach number. The snapshots were taken before, around the time of, and after the highest spike in the accretion rate. It is visible that in this case density profile disturbances appear on much larger scales than in the case of the models with subcritical and critical initial rotation. Before the spike there are disturbances in density profile with higher values in central regions and A_ϕ contour structure close to the equator plane with smaller density inside the sonic surface before the spike is perturbed. Near the equatorial plane there is an area with lower M_r values. What makes these profiles different from those previously presented is a scale of those disturbances. For the $A_0 = 0.85$ and $S = 1.4$ model, they are

reaching $r \sim 500r_g$, and for models with critical rotation, they are reaching $\sim 200r_g$. A spike in $\dot{M}$ corresponds to accretion of the matter from the poles. After this spike, accretion rate stays constant for around 0.1 s. During that time, remaining disturbances in the density profile are accreted, A_ϕ contours are getting back to the initial shape and the sonic surface is transforming to spherically symmetrical. At further stages of the simulation all profiles are getting smoothed and the sonic surface moves toward the outer boundary of the grid. In the case of the other simulations with $S = 1.4$, a similar spike in $\dot{M}$ does not occur. Evolution of the profiles is slightly different, in particular, the A_ϕ structure is less stable and loses its shape, but it also extends at its maximum to $r \sim 500r_g$.

3.4. Toroidal Component of the Magnetic Field

We start simulations with A_ϕ being the only nonzero vector potential component, which means that the B_ϕ component is absent. However, the toroidal magnetic field component appears shortly after simulation starts. We examined the B_ϕ distribution for characteristic models, presented in Sections 3.3.2 and 3.3.1.

1. Model NS-HM-04 ($\beta = 1$, $A_0 = 0.3$, and $S = 0.4$). For most of the evolution the flow is spherically symmetric, with the strongest toroidal magnetic field developing at the central part of the flow close to the black hole. Around time $t = 0.3$ s, a sphere of weaker magnetic field starts to move inward. Around time $t = 0.7$ s it stops around $r = 700r_g$. This feature maintains its position until $t = 1$ s. Then, it continues to move inward and at time $t = 1.197$ s the flow loses its spherical symmetry, for the toroidal field as well as other quantities.
2. Model LS-HM-10 ($\beta = 1$, $A_0 = 0.5$ and $S = 1.0$). Shortly after the beginning of the simulation, the toroidal component of the magnetic field appears. It follows similar small-scale structures to that in the density profile presented in Figure 16. The structure relaxes around $t = 0.4$ s. From this moment the flow evolution is spherically symmetric. Around $t = 0.5$ s a sphere with a lower value of B_ϕ starts to accrete inward. It reaches $r = 400r_g$ at $t = 1.449$ s, and then the flow loses its symmetry.
3. Model HS-HM-14 ($\beta = 1$, $A_0 = 0.85$ and $S = 1.4$). The toroidal component of the magnetic field that appears shortly after the beginning of the simulation has a similar shape to the profile of density seen in Figure 17. The low-density regions near the equator correspond to high toroidal magnetic field. This kind of a structure is maintained until $t \sim 0.7$ s. Then, the flow recovers its spherical symmetry, which is retained until $t \sim 1.4$ s, and a sphere with lower magnetic field strength is moving inward. Then, the symmetry is broken.
4. Model NS-LM-14 ($\beta = 100$, $A_0 = 0.3$ and $S = 1.4$). In this simulation with weaker magnetic field we observe disturbances in the innermost part of the flow (below $r \sim 100r_g$) at the beginning of the simulation, when the high-density regions are more highly magnetized. In this case, however, the matter distribution does not follow the B_ϕ distribution. Around $t = 0.488$ s the flow retains spherical symmetry. At $t \sim 0.7$ s a sphere with weaker magnetic field starts moving inward, and then at

$t \sim 0.95$ s this sphere stops at $r = 750r_g$ and remains there until the end of the simulation.

As noted above, the qualitative evolution of the toroidal magnetic field component depends strongly on initial parameters of the model.

Our results do not depend significantly on the adopted grid resolution. We have run an additional simulation with twice bigger resolution in both the r and θ directions, namely, 512×512 points, and compare the results quantitatively. We repeated simulations with $\beta = 100$, $A_0 = 0.3$, and $S = 1.0$ with this higher grid resolution. The comparison of the evolution of accretion rate, spin, and black hole mass shows that during the first ~ 0.6 s of the collapsar evolution, when the accretion rate increases, the black hole mass increase and spin increase are exactly the same. Both runs resulted in the same maximum spin value. The mean relative difference of black hole mass, accretion rate, and spin throughout the simulation is about 0.002%. The final black hole mass difference is 0.8% and is higher for higher resolution. There were no big differences in the profiles of density, radial Mach number, and angular momentum either, as checked at the characteristic moments of the collapsar evolution.

4. Discussion

In this section, we discuss the main assumptions, results, implications, and limitations of our models. Our simulations are divided into two sets with the *Magn* class of models, referring to a moderately magnetized accreting cloud, and the *Therm* class, where the cloud magnetization is neglected. Table 1 presents a summary of the models and their initial parameters, namely, the black hole spin and rotation of the envelope, and the final black hole mass and spin at the end of simulation. Also, the maximum accretion rate reached during the simulation is mentioned in Table 1.

We note that the normalization of rotation in the envelope is not done fully independently from the black hole spin value at the initial state. In reality, these two quantities should depend on each other. The angular momentum adopted in the simulation is assumed to be a fraction (S) of the momentum at the last stable circular orbit. The radius of this orbit depends on the spin value; therefore, the models with the same S parameter and different A_0 value have different magnitudes of angular momentum in the envelope. However, the picture is more complex, and our models show a nonlinear behavior of black hole spin and accretion rate evolution, depending on the initial setup.

The Alfvén velocity and the sound speed in the flow are in general related through $\sqrt{\beta} = \frac{c_s}{v_A}$. Initial conditions in our simulations prescribe the magnetic field potential as a function of θ only, and the field strength is normalized with the β parameter (Equation (14)). The sound speed is given by Equation (2), with thermodynamical parameters resulting from the Bondi solution, and depends on r . We checked that at the beginning of the simulation the ratio of $\frac{c_s}{v_A}$ calculated for both magnetic field normalizations gives $\frac{\beta_{100}}{\beta_1} = 10$ in every point of the grid, to confirm the viability of embedding the transonic Bondi accretion flow in the poloidal magnetic field normalized to the gas pressure.

There are no specific outflow/jet signatures in the magnetized models. A structure with dominating open magnetic field

lines can be seen in the zoom-in view up to $r = 100 r_g$ of the model LS-HM-10. This transient feature is accompanied by the existence of magnetized, turbulent torus of high density, located in the equatorial plane. Nevertheless, the polar regions are also having high density and do not resemble evacuated, low-density funnels present in other types of simulations (e.g., Sapountzis & Janiuk 2019). These authors started their models with a magnetized accreting torus embedded in a quasi-vacuum. In contrast, here we start our models with an accreting dense spherical cloud, so the jet break through such a dense envelope would be very difficult. We can speculate that if the torus structure existed for a longer time, the evacuation of polar regions on freefall timescale would eventually provide a clear structure for a long gamma-ray burst engine with a jet. However, it was not found in the currently presented simulations.

Our model neglects self-gravity effects of the accreting cloud and assumes that the dominant gravitational potential is enforced by the central black hole, whose point mass and spin are changing. The secondary effects and metric perturbations due to the envelope self-gravity might, however, be important, especially at the beginning of the collapse. To assess this effect quantitatively, we made an exemplary test calculation, limited to the nonmagnetized accretion (we adopted an initially nonspinning black hole and critical rotation of the envelope). We found that initially, for the first $\sim 3000 t_g$, the simulation results do not differ. Then, the main change in the evolution of black hole mass and spin is starting at about the time of $5000 M$, when both mass and spin increase much faster than in the simulation neglecting self-gravity. Finally, at time $\sim 30,000 t_g$, the models merge again, and ultimately the final mass and spin of the black hole are the same as found before. We note that this test simulation still neglected higher-order terms in metric evolution, i.e., connection coefficients and derivatives. The precise investigation of self-gravity effects is postponed to a separate future work.

In our simulations we use the initial condition of the relativistic Bondi solution, for the radial components, and on top of that we add a net angular momentum, centered on the equatorial plane. This initial condition describes a quasi-spherical accretion onto a black hole, and initially the assumption is made that the stellar envelope was in a steady state. However, because we do not add mass to the envelope during the time evolution, the accretion rate is dynamically changing. Moreover, by adding rotation to the flow, we relax the steady-state assumption, and as a result the mini-disk forms. We use the adiabatic index of $\gamma = 4/3$, which describes the microphysics in the stellar interior, including a large contribution from the photon pressure, rather than the isothermal index of $\gamma = 1$, which would be describing the protostellar disks. Also, as argued above, the dynamically changing Kerr background metric sufficiently describes the evolution of gravitational potential due to the increase of black hole mass and its spin changes. This is because the gravitational potential of the black hole dominates the spacetime, while the self-gravity of the mini-disk is a rather small effect. It could be treated as a small perturbation, but this effect is currently beyond the scope of the present work.

The magnetic field assumed in our models is normalized either to $\beta = 100$, which means thermally dominated accretion flows, or to $\beta = 1$, which means actually quite strong magnetization and equilibrium ratio of the magnetic and gas

pressures. The gas pressure is set by the solution of the Bondi transonic accretion flow, which is our initial condition for the simulation. Only if the magnetic pressure does not add a high contribution to this flow structure can we safely assume that the solution holds, and the gas falls supersonically to the black hole downstream of the sonic point, ultimately reaching the speed of light at the black hole horizon. We perturb this flow in time, by adding a slow rotation to the gas; hence, our sonic point is moving, and the shape of the sonic surface is disturbed. However, when the magnetic pressure is substantially large, the initial solution for the transonic accretion may not be a good approximation. Already at the beginning, multiple critical points might exist in the flow. The possible existence of shocks in low angular momentum flows connected with the presence of multiple critical points in the phase space has been studied from different points of view during the past 30 yr. However, the theoretical works that describe the fundamental properties of the low angular momentum accretion and that usually treat the steady solution of the equations have so far been carried out only for nonmagnetized and usually also nonviscous flows (Abramowicz & Zurek 1981; Abramowicz & Chakrabarti 1990; Das 2002; Das & Czerny 2012). Hydrodynamical models of the low angular momentum accretion flows have been studied already in two and three dimensions, e.g., by Proga & Begelman (2003) and Janiuk et al. (2008, 2009). In those simulations, a single, constant value of the specific angular momentum was assumed, while the variability of the flows occurred owing to, e.g., nonspherical or nonaxisymmetric distribution of the matter. The level of this variability was also dependent on the adiabatic index (see also Palit et al. 2019). However, these studies considered nonmagnetized flow only, which limits the applicability of these models to observational data and to the stellar evolution and presupernova studies (Heger et al. 2005). The range of strength of the magnetic field varies from very low values in the interstellar medium to extremely high values of $\sim 10^{15}$ G in the magnetars, which are extremely magnetized neutron stars. During the accretion process, the magnetic field in the star is likely to be amplified in the innermost region of the black hole, so a negligible amount of magnetic field in those regions, where the shock and the sonic point could be located, is thus not anticipated. The only method, however, to describe the effect of strong magnetic fields in the context of transonic accretion with low angular momentum is via numerical simulations, because no analytic solution of this problem exists, to the best of our knowledge. Hence, setting up the initial condition with strong magnetic fields overimposed on the slowly rotating transonic Bondi flow is a simple but working approach that we make here.

The physical conditions in the flow, namely, its density distribution and temperature, are governed by the global parameters, such as the mass of the accreting cloud and accretion rate, black hole mass and spin, the sonic radius, and the magnetic field normalization. The density normalized to the mass scale of the cloud enclosed in our computational domain is chosen to match the densest parts of the accretion torus, as can be found in the gamma-ray burst central engines ($\rho \lesssim 5 \times 10^{10} \text{ g cm}^{-3}$). Also, the chosen value of the sonic radius sets the temperature of the flow. At the beginning of the simulations, we have the temperature range above $3.76 \times 10^{10} \text{ K}$, consistently with the physical conditions in the neutrino-cooled torus formed inside the collapsar, as modeled, e.g., by Lopez-Camara et al. (2009). We note that

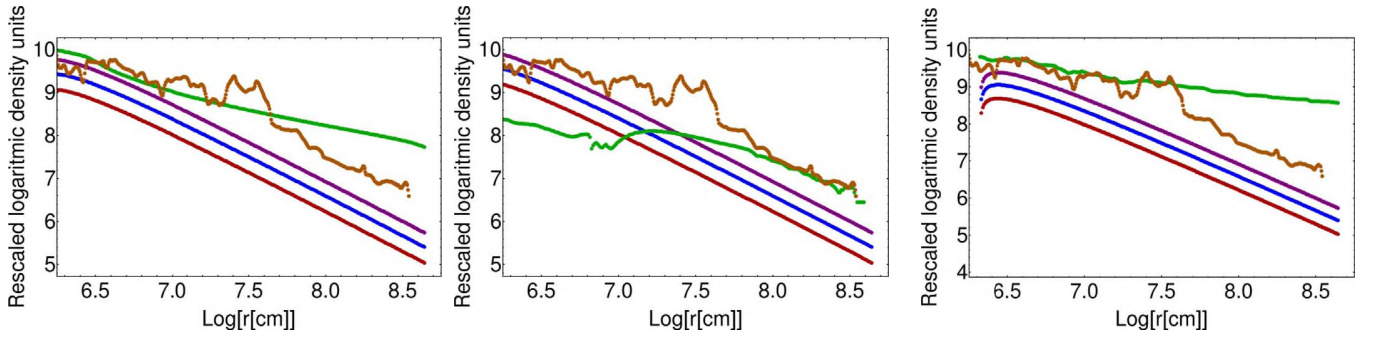


Figure 18. Comparison of the density profiles at the equator obtained during our simulations (green) with density profiles of the Shakura–Sunyaev disk with $\alpha = 0.1$ (red), $\alpha = 0.3$ (blue), and $\alpha = 1.0$ (purple). Additionally, we plotted the density of the evolved magnetized torus model, started from the Chakrabarti solution, taken at time $t = 1500r_g$ (brown lines). Density units are rescaled. The left panel shows our simulation without magnetic field with $A_0 = 0.3$ and $S = 1.0$, and in the right panel we show the simulation with $\beta = 1$, $A_0 = 0.5$, and $S = 1.0$.

presupernova models from Heger et al. (2005) and Heger et al. (2000) show values of the temperatures in the inner parts of the star around 8×10^9 K, while the densities of these presupernova stars are $\rho \sim 8 \times 10^9$ g cm $^{-3}$.

The position of the shock front and its expansion velocity may serve to estimate the energy dissipation rate in the shocks. In the case of simulations without magnetic field we detect four shock fronts that move outward, and we can determine their velocities. Two shock fronts were found in the model with $A_0 = 0.5$ and $S = 1.0$, and their velocities are $0.014c$ and $0.022c$. For the models with subcritical envelope rotation we obtained higher velocities of the shock fronts: $0.04c$ and $0.044c$ for $A_0 = 0.5$ and $A_0 = 0.85$, respectively. This trend seems opposite in comparison to Murguia-Berthier et al. (2020), where the shock front velocities were growing for the increasing angular momentum in the envelope. However, their simulations had always higher angular momentum assumed (corresponding to $S = 1.5$, $S = 1.75$, and $S = 2.0$). Moreover, the fundamental difference in the physics of the code was that the mass and spin of the black hole, and hence Kerr metric coefficients, did not change. Such an approach is justified only if the black hole mass is much larger than the adopted envelope’s mass and if the simulation timescale is short. In our model, such an assumption is no longer valid. We note that the shock velocities obtained here are on average one order of magnitude smaller than the escape velocities corresponding to positions of the shocks, which means that they cannot significantly halt the collapse.

Interestingly, we find that the evolution pattern in most of the models is similar, although one particular initial parameter is crucial and affects the final outcome. For the nonmagnetized envelopes, the evolution mostly depends on stellar rotation rate (S value). In particular, the large-amplitude oscillations of the accretion rate appear only for supercritical rotation $S = 1.4$. These models also result in the smallest final black hole masses and largest spins, regardless of the initial values of the spin.

In the case of a magnetized envelope, things change, and evolution is not only determined by the rotation rate. For subcritical rotation, the magnetic fields seem to not affect much the collapsars’ evolution. The distributions of density and Mach number are similar to the respectively modeled nonmagnetized cases. The accretion rate oscillations appear, however, in both critical and supercritical rotation models, in the $\beta = 1$ case, while the amplitude of the peaks is lower than in thermally dominated models without magnetic fields. In contrast, in the case of $\beta = 100$ the accretion rate in simulations

with $S = 1.0$ behaves more smoothly and is very similar to the subcritical case. This is due to mostly almost spherically symmetric distribution of density and only slight flattening when a mini-disk forms at the end of collapse.

The grid resolution of our simulations is not enough to probe the magnetorotational instability in great detail. Nevertheless, several results reveal the important coupling between the magnetic fields and rotation of the collapsing matter. For the supercritical rotation, the magnetic fields affect the collapse process in strongest way. In the $S = 1.4$ models, the spikes and oscillations at accretion rate are present around 0.4 s for every initial spin value, A_0 . Among all the magnetized models, one specific combination of $A_0 = 0.85$ and $S = 1.4$ presents a somewhat different evolution pattern. For $\beta = 100$, the model results in the maximum spin of the black hole, which reaches $A = 1$ temporarily (however, it drops to $A \sim 0.65$ at the end of simulation). In other words, the rotation of the envelope influences the collapse depending on the specific β value.

In some models we obtained disk-like structures at the equatorial plane, which sustained for a period of time during MHD simulations. We can compare their structure, i.e., equatorial density profiles, with the standard solutions known from accretion disk theory (Shakura & Sunyaev 1973, hereafter SS73). Here we present the standard, stationary accretion disk model with $\alpha = 0.3, 0.5$, and 1 . It is shown with three different lines in Figure 18. Our simulated disk-like profiles are shown for three chosen simulations: (i) small mini-disk that is formed around $t = 1.5$ s in the simulation with $A_0 = 0.3$ and $S = 1.4$, shown in the left panel; (ii) $A_0 = 0.85$ and $S = 1.0$, shown in the middle panel, where our simulated disk appears around $t = 1.5$ s and remains for a few seconds; and (iii) disk structure that is formed in the simulation with $\beta = 1$, $A_0 = 0.5$, and $S = 1.0$, presented in the right panel. The simulation results are marked with the green solid line. Disk (i) presents the snapshot taken at time $t = 1.48$ (see Figure 8), disk (ii) shows the snapshot taken at $t = 2.04$ (see Figure 7), and disk (iii) shows the snapshot taken at $t = 0.163$ (see Figure 16).

We notice that the disk forming from the nonmagnetized collapsar model with supercritical rotation speed agrees with the SS73 disk only in the inner parts of the flow. Then, the density profile flattens in the outer parts. In this simulation, the angular momentum is not transported out by viscous forces, while the rotation of the envelope is sub-Keplerian. Also, the same holds for the nonmagnetized simulation of the mini-disk with critical rotation. Also, the model with magnetic field and

critical rotation diverges from the SS73 disk structure at the outer parts of the flow. This model, however, is best matched by another “standard” accretion model. We compared it with the density profile of the evolved Chakrabarti (1985) torus structure, taken from the GR MHD simulation, and evolved at late time, $t = 1500M$ (see Janiuk et al. 2021). Profiles shown in Figure 18 have arbitrarily normalized density units.

Finally, let us speculate about the relevance of our findings to the observed black hole systems and their parameter estimates. In Figure 9 we presented the parameter space of the nonmagnetized models, depicting the results depending on the initial black hole spin and rotation strength. As the results show, the final black hole mass is anticorrelated with the final black hole spin, which is in a good agreement with the results from gravitational wave observations (Safarzadeh et al. 2020). They show that most heavy (stellar-mass) black holes tend to be spinless, and their spins do not exceed $a \sim 0.4$, as given by the probability distribution analysis (Roulet & Zaldarriaga 2019). It should be noted, however, that our results consider only black holes formed in a collapsing star, without the interaction with any companion in a binary system. The parameter space used in our analysis (see Figure 9) covers different values of an initial black hole spin at the beginning of collapse and the angular momentum content in the envelope. In general, these values might represent the state of the collapsar as reached at some evolved state in a binary system, when the envelope and/or stellar core has been spun up by the companion. This state, however, refers to only a very close contact, or a common-envelope phase, rather than a detached system with long time accretion.

In the case of the black holes accreting in the low- and high-mass X-ray binaries (LMXBs and HMXBs), even in the wind-fed HMXBs the accreted material has enough angular momentum from the orbit so that it forms a disk around the black hole before it is accreted (only at larger scales will the accretion flow be Bondi-like). Typical black holes in HMXBs are expected to accrete only up to a few tenths or at most $1 M_{\odot}$ over their lifetimes, and the common belief is that no spin-up is possible via accretion on these systems.

On the other hand, the measured spins of black holes in LMXBs cover the whole range of spin parameters (Fragos & McClintock 2015). Our simulations cover the parameter space in which the specific angular momentum accreted by the black hole is not much larger than that of the ISCO. Nevertheless, the maximum black hole spin achieved temporarily in all the nonmagnetized runs is very high, if only rotation is larger than critical. Therefore, the constraint that we find on the angular momentum content (or the magnetic fields in the case of the magnetized flow) might be relevant for the evolutionary history of some particular black hole binary systems, such as GRS 1915+105 (McClintock et al. 2006).

In systems like GRS 1915+105, the spin is “directly” measured from the spectral fitting (in this microquasar we have an $a = 0.95$ estimate). In addition, the black hole mass for this microquasar is largest among known sources of this type in our Galaxy (over $12 M_{\odot}$). In others, like GX 339-4, the indirect method for the black hole spin estimate, based on the estimated jet power, is used (for this system it results in $a = 0.25$). For several further LMXBs there is no observational constraint for the black hole spin, while the estimates of the black hole mass give values from 7.5 (for XTE J1118+480) up to $9.0 M_{\odot}$ (for GS 2023+338).

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Chapter 3

Possible Gravitational Microlensing Events in the Optical Lightcurve of Active Galaxy S5 0716+714



Possible Gravitational Microlensing Events in the Optical Lightcurve of Active Galaxy S50716+714

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Abstract

A well-known active galaxy of the blazar type, S5 0716+714, is characterized by a particularly high-variability duty cycle on short timescales at optical frequencies. As such, the source was subjected to numerous monitoring programs, including both ground-based as well as spaceborne telescopes. On closer inspection of the most recent accumulation of the data provided by the Transiting Exoplanet Survey Satellite, we have noticed several conspicuous events with “volcano-like” symmetric shape, all lasting for several hours, which closely resemble the achromatic events detected with the previous Whole Earth Blazar Telescope campaigns targeting the source. We propose that those peculiar features could be due to the gravitational microlensing of the innermost segments of the precessing jet in the system by a binary lens. We study the magnification pattern of the lens with the inverse-ray shooting method and the source trajectory parameters with the Python package MULENSMODEL. In this way, we were able to fit successfully all the selected events with a single lens, adjusting slightly only the source trajectory parameters for each lensing event. Based on the fit results, we postulate the presence of a massive binary lens containing an intermediate-mass black hole, possibly even a supermassive one, and a much less massive companion (by a factor of $\lesssim 0.01$) located within the host galaxy of the blazar, most likely the central kiloparsec region. We discuss the major physical implications of the proposed scenario regarding the quest for the intermediate-mass and dual supermassive black holes in active galaxies.

Unified Astronomy Thesaurus concepts: Black holes (162); Gravitational lensing (670); Active galaxies (17); Blazars (164); Active galactic nuclei (16)

1. Introduction

Gravitational lensing has wide application in cosmology and astrophysics (for reviews, see, e.g., Blandford & Narayan 1992; Narayan & Bartelmann 1996). Among many others, it was proposed that microlensing can possibly explain at least some part of the variability of active galactic nuclei (AGNs), in particular the achromatic, very short-timescale flux changes in blazar sources (e.g., Subramanian et al. 1985; Schneider & Weiss 1987; Gopal-Krishna 1991; for a review, see Wambsganss 2006). Lately, this idea has been somewhat resurrected by Vedantham et al. (2017) and Peirson et al. (2022), who proposed that the peculiar yearlong features in the radio lightcurve of blazar J1415+1320, consisting of a symmetric double-horn minimum (designated by the authors as a “volcano-type” shape event), could be explained by “milli-lensing” events involving a binary lens.

Two-point-mass lensing is a well-understood phenomenon (Schneider & Weiss 1986). For example, gravitational lensing by a binary star leading to the double-horn shape in the lightcurve of a lensed source has been discussed by Mao & Paczynski (1991) in the specific context of Galactic stars and planets. The authors pointed out that when the source is projected close to the caustics, an extra pair of unresolved microimages is formed, creating complicated patterns in the source lightcurves, including the characteristic double-horn features.

In this paper, we present another take on the story. Namely, we identify peculiar symmetric features in the optical lightcurve of the blazar S5 0716+714, lasting for several hours each, with similar double-horn flux profiles. We propose that all of these could in fact be due to the lensing of the innermost segments of the precessing jet in the source by a binary lens consisting of an intermediate-mass black hole (IMBH) or even a supermassive-mass black hole (SMBH) with a less massive companion (mass ratio $\lesssim 0.01$) located within the host galaxy of the blazar.

S5 0716+714 is a very well-known blazar located at redshift $0.227 < z < 0.254$ (Dorigo Jones et al. 2022). For modern cosmology with $H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1}$, $\Omega_m = 0.3$, and $\Omega_\Lambda = 0.7$, this corresponds to angular-size distance of $d_A \simeq (0.75 - 0.82) \text{ Gpc}$ and conversion scale of $\sim (3.64 - 3.96) \text{ kpc arcsec}^{-1}$. The blazar is characterized by a particularly high duty cycle on short timescales at optical frequencies and as such was subjected to numerous intranight monitoring programs, including several campaigns on behalf of the Whole Earth Blazar Telescope (WEBT) Collaboration³ (e.g., Ostorero et al. 2006; Bhatta 2016), and most recently with the Transiting Exoplanet Survey Satellite (TESS) by Raiteri et al. (2021).

On closer inspection of the most recent TESS lightcurves of the target, we have noticed at least six several-hours-long conspicuous events, whose “volcano-like” symmetric shape resembled the features discussed by Vedantham et al. (2017) and Peirson et al. (2022) in the radio lightcurve of J1415+1320. We note, however, that the J1415+1320 radio events are much longer, lasting up to even a few years. Interestingly,

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³ <http://www.oato.inaf.it/blazars/webt>

the identified TESS events resemble closely achromatic features detected previously in the source optical lightcurve with WEBT (Bhatta 2016). As we argue below, all of those could be explained in the framework of the microlensing scenario.

2. TESS Data

For our analysis, we use short-cadence (120 s) TESS mission (Ricker et al. 2015) data for S5 0716+714 (TESS Input Catalog source, TIC 147796121) downloaded from the Barbara A. Mikulski Archive for Space Telescopes.⁴ The target was observed in four seasons:

1. 2019-12-25–2020-01-20, sector 20, camera 2, ccd 2
2. 2020-06-09–2020-07-4, sector 26, camera 4, ccd 4
3. 2021-06-25–2021-07-23, sector 40, camera 4, ccd 3
4. 2022-01-01–2022-01-27, sector 47, camera 2, ccd 2

corresponding to the four panels of Figure 1. We adopted BTJD = BJD – 2457000. Only the first observational season was the subject of the detailed analysis by Raiteri et al. (2021).

The Target Pixel Files data (containing the original CCD time series data in the shape of postage stamp image cutouts) from all the seasons were checked for possible field stars that could fall within the apertures/masks used by the TESS Science Processing Operations Center pipeline (TESS-pipeline) to extract the lightcurves.

S5 0716+714 appears to be a well-isolated object, and only faint (>19 mag) field stars can be found in the TESS-pipeline masks. This is confirmed by the CROWDSAP factor (describing the amount of light in the aperture attributed to the target object), which is between 0.962 and 0.982 in all observing seasons. A small but negligible light contamination comes from neighboring bright stars outside of the TESS-pipeline masks. Therefore, for further analysis, we used Pre-search Data Conditioning Simple Aperture Photometry fluxes (systematics-corrected photometry using cotrending basis vectors), which are lightcurves prepared by the TESS-pipeline. The TESS fluxes F were converted to R -band magnitudes following Raiteri et al. (2021), namely $\text{mag}_{\text{TESS}} = -2.5 \log F + 20.42$. For all subsequent analyses, we binned the data in 10 minute intervals. The resulting lightcurves of S5 0716+714 are presented in the four panels of Figure 1.

After a close inspection of the acquired TESS lightcurves, we found six notable events displaying symmetric double-horn minima. In fact, these lightcurve shapes resemble events previously found in ground-based WEBT campaigns in 2003 (Ostorero et al. 2006) and 2014 (Bhatta 2016), where “for about 6 hr (...) the source suddenly exhibited a strongly reduced level of flux variability, resulting in a plateau in all four bands’ light curves” (Bhatta 2016, Section 3.1.4 therein). One of the selected TESS events, with the superimposed WEBT 2003 and 2014 events (shifted vertically by -1.45 mag and -1.3 mag, respectively), is presented in Figure 2. Note that the offsets of ~ 1.3 mag, translating to a factor of ~ 3 in flux, imply that the double-horn events can be found at a relatively wide range of flux levels of the source. Moreover, as we argue below, such a repeating pattern does not conform with the stochastic, red-noise type variability of the blazar (e.g., Rani et al. 2013; Bhatta 2016; Raiteri et al. 2021).

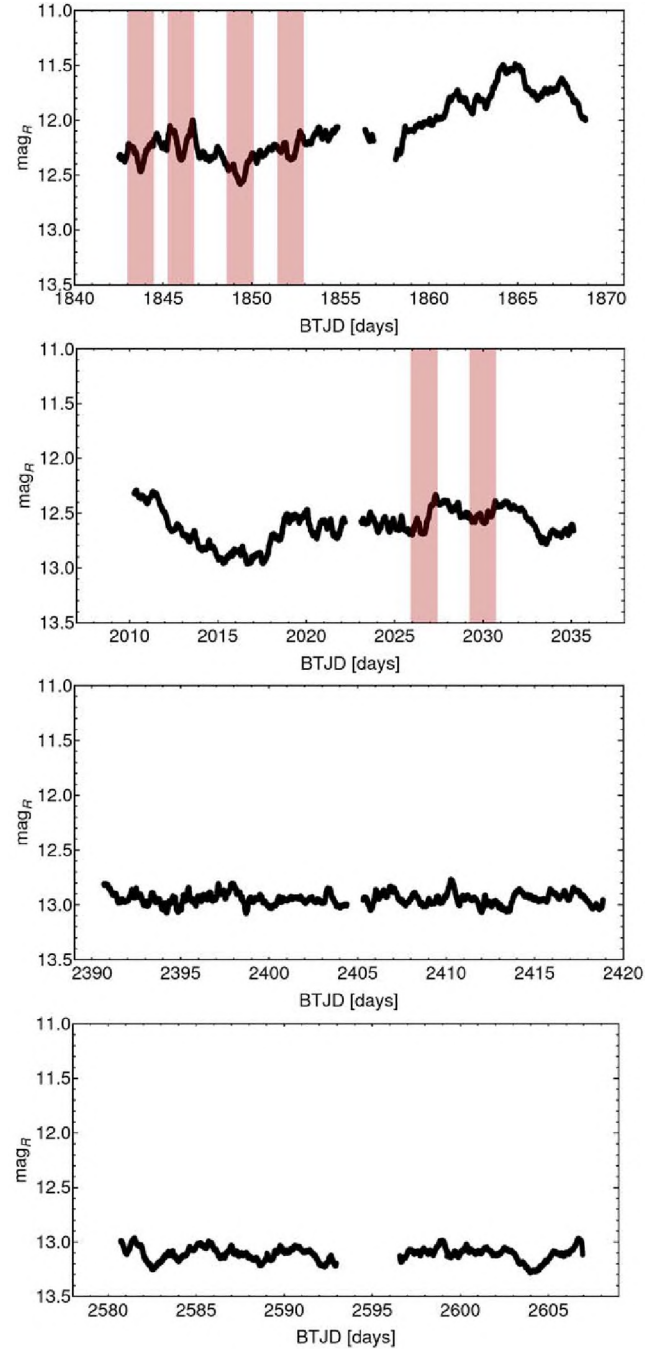


Figure 1. Four monthlong TESS observations of S5 0716+714 investigated in this paper. The TESS data were converted to R -band magnitudes and binned in 10 minute intervals (see Section 2).

We emphasize the achromatic character of the 2014 WEBT plateau event that was clearly demonstrated by Bhatta (2016). As for the TESS data, while here we show only the TESS lightcurve converted to R -band magnitudes, the accompanying WEBT monitoring—even though with a more sparse coverage and not overlapping with the entirety of the TESS pointings—confirms the achromatic character of the flux changes also in the 2019/2020 data set (see Figure 3 in Raiteri et al. 2021). In particular, four out of the six events selected here for analysis have the accompanying WEBT data, which clearly show that

⁴ <https://archive.stsci.edu/index.html>

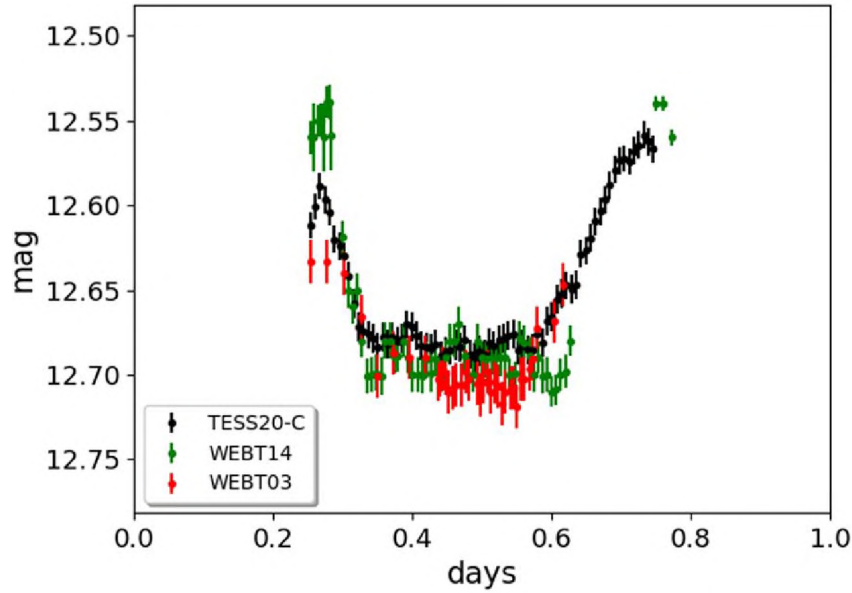


Figure 2. Zooms of the TESS event on MJD 9026 (black data points), with the superimposed data for the plateau events present in the 2003 WEBT data (red data points; from Ostorero et al. 2006, shifted vertically by -1.45 mag), and 2014 WEBT data (green data points; from Bhatta 2016, shifted vertically by -1.3 mag). The WEBT data points were taken directly from Figure 9 in Bhatta (2016).

the flux changes traced by the dense TESS sampling, are followed closely in all the B , V , R , and I filters separately.

All in all, during the total of four months of the TESS monitoring time, we identified six potential microlensing events in the optical lightcurve of S5 0716+714. The question therefore arises on the frequency and the duty cycle of such events. In this respect, first we note that there may be more similar features in the TESS lightcurve, which are, however, more difficult to spot because of the superimposed flaring activity of the jet contaminating the magnification pattern. And indeed, due this reason, several other possible features in the TESS data were not selected for the analysis presented below: when viewed with a sufficient zoom, they either appeared asymmetric, or the plateau during the flux minimum phase displayed some small-amplitude variability. Second, we note that whenever an intense monitoring campaign with sufficiently long time coverage (of at least a handful of days) and sufficiently dense sampling (like the WEBT 2003/2014 or the TESS campaigns) had targeted the source, at least one microlensing candidate feature could be seen. On the other hand, the majority of the intranight observations of the blazar with a single ground-based telescope, lasting for several hours per night at most, are typically insufficient to cover the entirety of a single magnification event, precluding any robust identification of such.

To sum up, we conclude that the symmetric double-horn features in the optical lightcurve of S5 0716+714, with peak-to-peak duration of about 12 hr and the flux minimum phase lasting for a few/several hours, are promising microlensing-event candidates. Those events appear relatively frequent although one needs more good-quality monitoring data to estimate robustly the duty cycle and clustering properties of such.

3. Binary Lens Modeling

In the first stage of our analysis, we used the inverse-ray shooting method to obtain magnification patterns for different

binary parameters. If we consider objects with masses M and qM , where $q < 1$, at the respective positions $\mathbf{x}_M = [0, 0]$ and $\mathbf{x}_{qM} = [0, d]$ on the lens plane, the lens equation takes the form

$$\mathbf{y} = \mathbf{x} - \frac{\mathbf{x}}{|\mathbf{x}|^2} - q \frac{\mathbf{x} - \mathbf{d}}{|\mathbf{x} - \mathbf{d}|^2}, \quad (1)$$

where $\mathbf{y}$ and $\mathbf{x}$ are two-dimensional angular coordinates of the rays on the source and lens planes (Schneider et al. 1992). In the above, all the quantities are in units of the Einstein angle

$$\Theta_E = \left(\frac{4GM}{c^2} \frac{D_{LS}}{D_L D_S} \right)^{1/2}, \quad (2)$$

where D_S , D_L , and D_{LS} are the angular distances to the source, to the lens, and between the source and the lens, respectively. Note that the Einstein radius $R_E = \Theta_E D_L$.

To obtain magnification patterns of the lens characterized by given q and d in the inverse-ray shooting method, we evenly distributed a set of rays on the lens plane and calculated the position of each ray on the source plane. For each bin of rays, the magnification is the ratio between the number of rays that ended up in it in the presence of and in the absence of the lens. In our investigation of the parameter space, we have calculated magnification patterns for various binary mass ratios from the range $q \in [0.001, 0.5]$ and separation $d \in [0.5, 1.5]$. We note that for small separation d , magnification occurs predominantly in a compact area surrounding the center of the mass of a binary lens. If, in addition, the mass ratio is small (in practice, below 0.001), the magnification pattern resembles a point-mass lens. As the mass ratio increases, the area of magnification is also enlarged, creating various shapes whose exact appearance depends on the separation parameter. The volcano-type events with durations and magnifications resembling the ones presented in Figure 2 could be obtained for the combination $q \lesssim 0.01$ and $d \simeq 0.7$.

The binary parameters, q and d , determine the magnification pattern. However, in order to obtain a specific magnification

curve matching a given feature present in the lightcurve, we need four parameters defining the source trajectory (the impact parameter between the source and the center of the mass of the lens in units of the Einstein angle, u_0 ; the Einstein crossing time, t_E ; the time of the closest approach between the source and the lens, t_0 ; and the angle between the source trajectory and binary lens axis, α), as well as the source parameter ρ , which is the radius of the source expressed as a fraction of the Einstein radius. In the next step of the analysis, we have therefore performed a formal fitting of those parameters to the six selected events of the source lightcurve, utilizing the Python package MULENSMODEL for modeling microlensing (Poleski & Yee 2019). For magnification calculation we have chosen the VBBL method (Bozza 2010; Bozza et al. 2018) implemented in MULENSMODEL; fitting was performed by minimizing χ^2 , using the algorithm described in Gao & Han (2012).

Due to a large number of parameters, it was impossible to perform fitting with all of the model parameters allowed to vary at the same time. For this reason, keeping in mind the general features of the explored magnification patterns from the inverse-ray shooting studies, in the first step of the analysis with the MULENSMODEL, we investigated the lens parameters by fixing q and d at different values for various source trajectories and compared the quality of the obtained fits. It was relatively straightforward to find in this way lenses that reproduced either the first two (TESS19-A and TESS19-B) or the last four (TESS20-A to TESS20-D) events; however, finding a lens that can fit all six events for similar source trajectories was challenging, and the best results could be obtained for the $q \sim 0.005$ and $d \sim 0.75$ combination. To derive the final, more precise values of the lens parameters, in the next step of the analysis, we fit the TESS20-D event for a given (fixed) representative source trajectory, with q and d allowed to vary (initial values of 0.01 and 0.7, respectively), resulting in the best-fit values of $q = 0.00545$ and $d = 0.0727$. In the final, third step of the fitting procedure, the fit was performed for all the selected events with the fixed lens parameters $q = 0.00545$ and $d = 0.0727$ by thawing the source trajectory parameters t_E , t_0 , u_0 , and α . It should be noted at this point that the particular choice of the ρ parameter does not affect the MULENSMODEL fitting results, as long as the source appears point-like, meaning $\rho \ll 1$. The final model was chosen based on the χ^2 statistic.

As shown in Figure 3 and summarized in Table 1, the selected events could all be fitted reasonably well in this way with only slight changes in the source trajectory parameters u_0 , t_E , α , and t_0 . This is further visualized in Figure 4, where we present the magnification map of the lens along with the source trajectory, along with the resulting source lightcurves, corresponding to the MULENSMODEL. The emerging set of the source trajectory (on the lens plane) parameters is obviously not unique, and it is possible that comparably good fits could be obtained for the other values of u_0 , t_E , α , and t_0 . The point of the fitting exercise presented here is, however, to check whether *the same single binary lens* is sufficient to explain all the selected events. Moreover, we claim that this lens has to be characterized by the parameters close to the selected values of $q \lesssim 0.01$ and $d \sim 0.7$.

4. Discussion and Conclusions

The distance to the active galaxy S5 0716+714, $z = 0.31 \pm 0.08$, has been determined by Nilsson et al. (2008), based on the “standard candle” assumption for the detected host galaxy. Danforth et al. (2013), based on the detection of the narrow

Ly α forest, confirmed the 2σ confidence redshift range for the blazar $0.2315 < z < 0.322$. More recently, Dorigo Jones et al. (2022) improved the Ly α forest analysis and obtained a 95% confidence redshift range of $0.227 < z < 0.254$. With such, the duration and the magnification of the TESS microlensing events in the S5 0716+714 optical lightcurve dictate that the lens is located in close proximity to the source, $D_{LS} \ll D_L$, and in particular within the host galaxy of the blazar as otherwise, the implied lens parameters (mass and separation) would become unrealistic. Hence, below we take $D_S \simeq D_L \simeq d_A$.

Moreover, the lens should be located within the central parts of the host as otherwise the probability for the jet component to cross the folded caustic would be rather low, in contrary to the frequency of the microlensing events in the TESS lightcurve.

In general, S5 0716+714 exhibits a stochastic variability at optical frequencies, with power spectral density (PSD) extending as a colored-noise continuum over several decades of the variability timescales, from months and years down to hours and minutes (see, e.g., Rani et al. 2013; Bhatta 2016; Raiteri et al. 2021). In particular, the PSD analysis for the first season of the TESS campaign returns the best-fit power-law slope of $\alpha = 2.01 \pm 0.04$ (Raiteri et al. 2021), consistent with the overall red-noise behavior. In addition, however, by means of analyzing the source periodograms, the autocorrelation function, and the structure function, Raiteri et al. (2021) found the characteristic variability timescales in the system of about 1.7, 0.5, and 0.2 days. It is not clear if these timescales could be linked to the microlensing events as discussed here, especially as there are indications that some of the rapid optical flares in the source covered with WEBT are chromatic; yet 0.2 and 0.5 days do correspond, at least roughly, with the duration of the flux drop/rise phases and the total duration of the analyzed microlensing-event candidates, respectively. Hence we cautiously conclude that the timing analysis of the TESS data do not rule out repetitive—even in a quasiperiodical fashion—microlensing-type structures in the red-noise lightcurve of S5 0716+714.

The overall colored-noise nature of the S5 0716+714 optical variability indicates that the observed emission of the blazar is produced by various emission components distributed within an extended segment of the jet for which the scales range from tenths and hundredths of parsecs down to the smallest scales set by the gravitational radius of the SMBH in the galactic nucleus:

$$r_g = \frac{GM}{c^2} \simeq \left(\frac{M}{10^8 M_\odot} \right) \text{ au} \quad (3)$$

(see, in this context, Begelman et al. 2008). We posit that the microlensing events (as the ones analyzed here) become conspicuous whenever the innermost segments of the jet, with linear sizes $R \gtrsim r_g$, dominate the radiative output of the source. A very complex kinematics of the S5 0716+714 jet—with frequent changes in the jet position angle, as well as a plethora of radio emission features on milliarcsec scales displaying various and often superluminal radial and nonradial motions (e.g., Britzen et al. 2010; Larionov et al. 2013; Rani et al. 2015)—means that one should not hope to see any strict regularity in such microlensing events as the source trajectory parameters u_0 , t_E , and α , may change/evolve with time. The encouraging finding, in this context, is that our fitting returns a rather narrow range for those parameters, in particular the Einstein crossing time each time of a similar order of a few/several days.

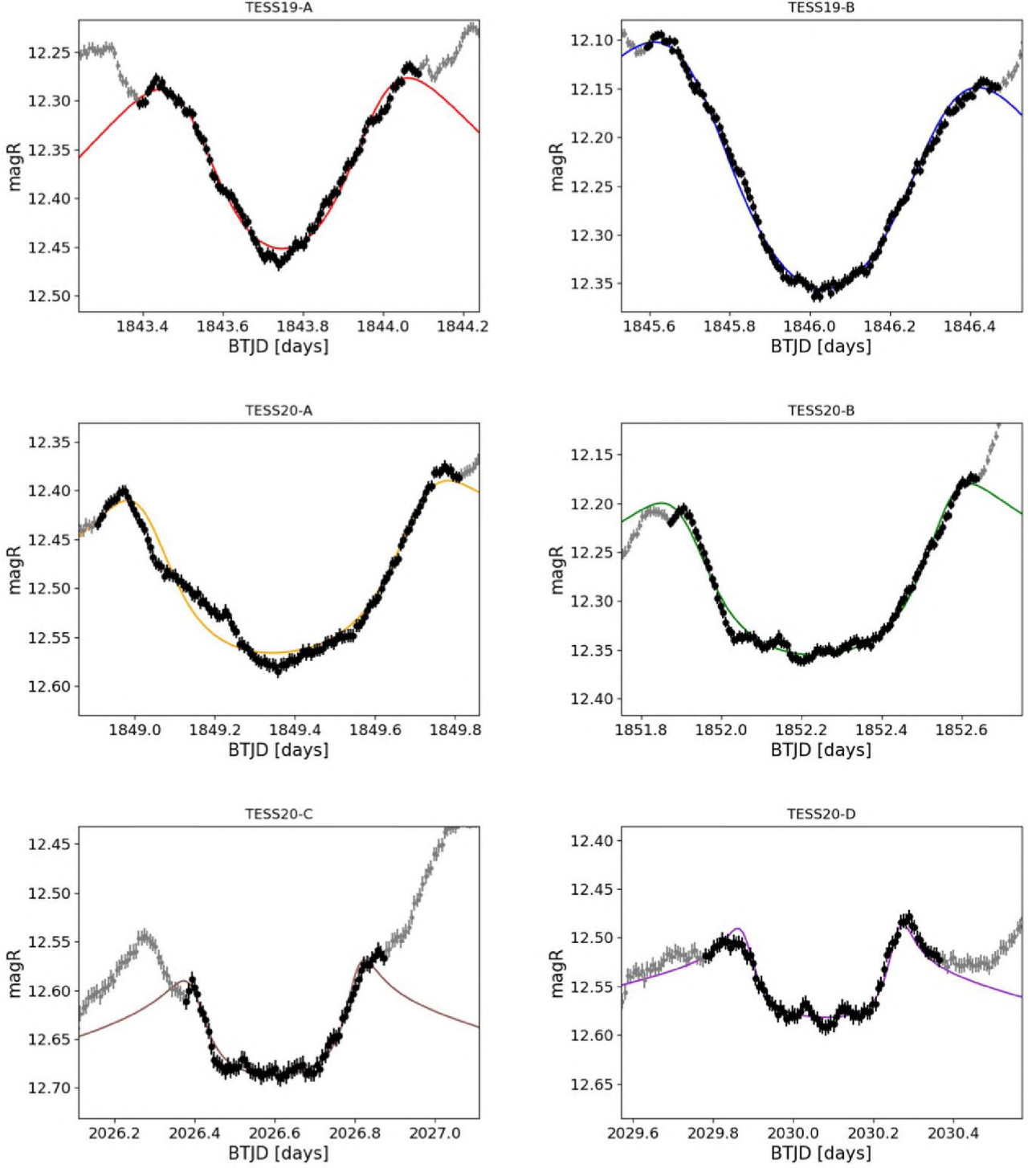


Figure 3. The MULENSMODEL fits (as described in Section 3 and summarized in Table 1) to the six selected events in the TESS lightcurve of S5 0716+714 (see Figure 1).

Keeping in mind the above considerations, for the analyzed microlensing events, we therefore set the limit $R = \rho R_E \geq r_g$, which, combined with Equation (2) gives

$$\left(\frac{M}{10^4 M_\odot} \right) \geq \left(\frac{M_\bullet}{10^8 M_\odot} \right)^2 \left(\frac{\rho}{0.01} \right)^{-2} \left(\frac{D_{LS}}{100 \text{ pc}} \right)^{-1}. \quad (4)$$

On the other hand, since by definition $R_E = c\beta_{\text{app}} t_E$, where β_{app} is the apparent velocity of the source emitting region through the lens, we also have

$$\left(\frac{M}{10^4 M_\odot} \right) \simeq 10^4 \left(\frac{\beta_{\text{app}}}{30} \right)^2 \left(\frac{t_E}{2 \text{ d}} \right)^2 \left(\frac{D_{LS}}{100 \text{ pc}} \right)^{-1}. \quad (5)$$

Table 1
MULENSMODEL Best-fit Parameters

Parameter	TESS19-A	TESS19-B	TESS20-A	TESS20-B	TESS20-C	TESS20-D
ρ	0.013 ^f	0.013 ^f	0.013 ^f	0.013 ^f	0.013 ^f	0.013 ^f
q	0.00545 ^f	0.00545 ^f	0.00545 ^f	0.00545 ^f	0.00545 ^f	0.00545 ^f
d	0.727 ^f	0.727 ^f	0.727 ^f	0.727 ^f	0.727 ^f	0.727 ^f
u_0	0.36	0.35	0.48	0.78	0.75	0.54
t_E	3.0	4.0	3.4	2.4	1.5	1.7
α	89	94	82	88	87	94
t_0 [BTJD]	1843.78	1845.92	1849.62	1852.30	2026.66	2030.00
t_0 [UTC]	2019-12-26 06:36	2019-12-28 09:58	2020-01-01 02:46	2020-01-03 19:05	2020-06-26 03:55	2020-06-29 12:04
χ^2/dof	1.91	1.5	3.75	2.10	0.93	0.89

Note.

^f parameter frozen in the fitting procedure.

Unfortunately, the SMBH mass in S5 0716+714 is unknown, and optical measurements of β_{app} are not available. However, for a blazar of the BL Lac type, we expect in general $M_{\bullet} \gtrsim 10^8 M_{\odot}$ (Plotkin et al. 2011), and for S5 0716+714 in particular $\beta_{\text{app}} \lesssim 30$ has been measured at radio wavelengths (Rani et al. 2015). With such, setting in addition $\rho = 0.01$ and also $t_E \simeq 2$ days as the average crossing time (see Table 1), the lower and upper limits are set by Equations (4) and (5), respectively and correspond to the range of the binary lens mass $10^4 < M/M_{\odot} < 10^8$ for the distance of the lens from the active nucleus $D_{\text{LS}} \simeq 100$ pc and $10^2 < M/M_{\odot} < 10^6$ for $D_{\text{LS}} \simeq 10$ kpc. We emphasize that the apparent velocity of the jet features inferred from radio observations with milliarcsecond resolution may not necessarily provide a proper characterization of the innermost segments of the jet dominating the optical emission subjected to the microlensing events (as proposed here). Indeed, in the framework of the well-established magnetohydrodynamics (MHD) model for relativistic jets launched by accreting black holes, the jet bulk velocity increases from the jet base due to a gradual collimation and acceleration of the plasma outflow by magnetic forces (e.g., Lyubarsky 2009, and references therein). Hence, apparent velocities measured in the radio domain at larger (parsec) distances from the center, $\beta_{\text{app}} \lesssim 30$ in this case, provide us a very robust *upper limit* for the proper motion velocity of the optical emission regions.

Another constraint on the lens could, in principle, follow from the requirement that the binary parameters should be stable during a single lensing event. However, the orbital velocity of the companion,

$$v_{\text{orb}} = \sqrt{\frac{GM}{d R_E}} \sim 370 \left(\frac{M}{10^4 M_{\odot}} \right)^{1/4} \left(\frac{D_{\text{LS}}}{100 \text{ pc}} \right)^{-1/4} \text{ km s}^{-1}, \quad (6)$$

is sufficiently low for the range of M and D_{LS} discussed above, that the distance traveled by the companion during the Einstein crossing time is always less than the separation of the lens. In other words, the condition $v_{\text{orb}} t_E \ll d R_E$ is easily satisfied for a wide range of lens masses located at larger distances, given the

Einstein radius of the lens

$$R_E \simeq 100 \left(\frac{M}{10^4 M_{\odot}} \right)^{1/2} \left(\frac{D_{\text{LS}}}{100 \text{ pc}} \right)^{1/2} \text{ au}. \quad (7)$$

On the other hand, the configuration of a binary lens could change over the 20 yr timescale (since the WEBT 2003 event), especially for the lower ranges of the mass M and the source–lens distance $D_{\text{LS}} < 100$ pc. Changes in the position of the $0.00545 M$ companion on the orbit around M will lead to the rotation of the magnification pattern. But a double-horn feature in the lightcurve can be obtained even for $\alpha \simeq 90^\circ \pm 20^\circ$ at the expense of a slight asymmetry of the “horns” (i.e., of the magnified flux maxima). Moreover, the characteristic microlensing pattern can still be produced even with larger rotation angles by means of more substantial changes in the source trajectory parameter introduced by the precession of the jet on a similar timescale of years/decades; indeed, observational evidence for such a precession in S5 0716+714 has been discussed in, e.g., Bach et al. (2005) and Nesci et al. (2005). These two effects—namely, the orbital modulation of the lens and the jet precession, both taking place on comparable timescales, leading altogether to a systematic evolution in the magnification patterns for the lensing events—could, in fact, offer an exciting opportunity for an observational verification of the model proposed, if only the target could be monitored at optical and radio frequencies over a longer period of time with sufficiently dense coverage.

Yet another consequence of the event detected by WEBT in 2003 is that one may require the Einstein radius to be large enough to cover the blazar over a timescale of, at least, $\tau \gtrsim 20$ yr despite the host galaxy rotation. For a typical blazar with a $M_{\bullet} \sim 10^8 M_{\odot}$ black hole, the velocity dispersion of the host is expected to be around $\sigma_{\text{vel}} \simeq 200 \text{ km s}^{-1}$. With such, assuming the lens is beyond the central SMBH sphere of influence (central ~ 10 pc region), the condition $\sigma_{\text{vel}} \tau < R_E$ gives

$$\left(\frac{M}{10^4 M_{\odot}} \right) > 10^2 \left(\frac{\sigma_{\text{vel}}}{200 \text{ km s}^{-1}} \right)^2 \times \left(\frac{\tau}{20 \text{ yr}} \right)^2 \left(\frac{D_{\text{LS}}}{100 \text{ pc}} \right)^{-1}, \quad (8)$$

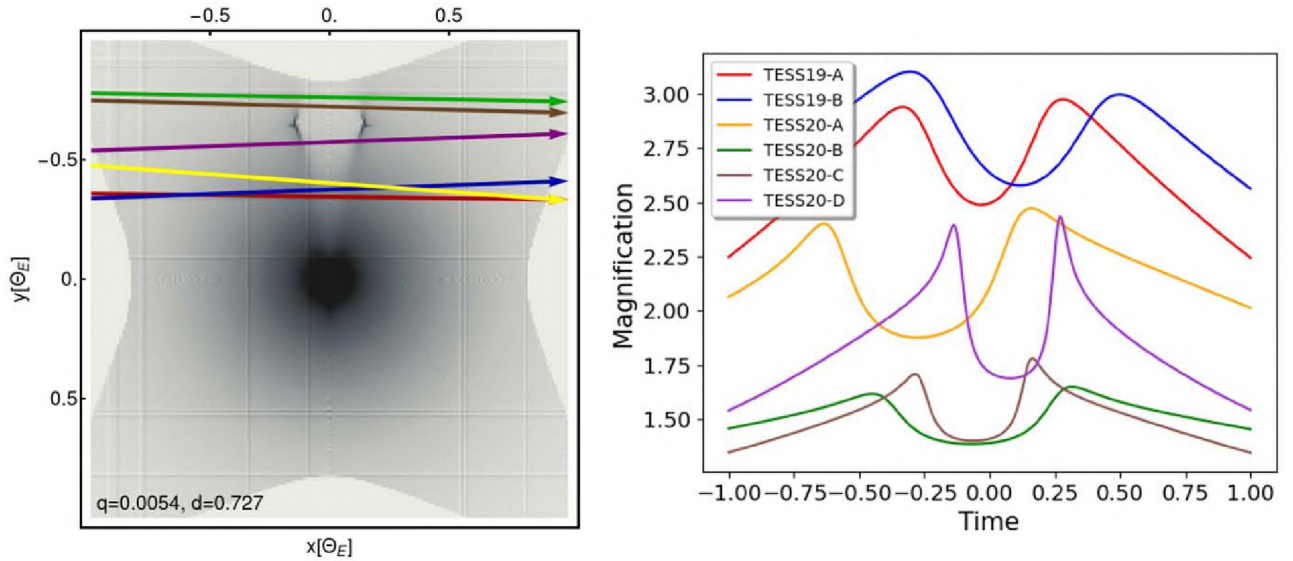


Figure 4. The magnification pattern along with the source trajectory (left panel) and the resulting source lightcurves (right panel), corresponding to the six MULENSMODEL fits presented in Figure 3.

which provides the strongest lower limit on the mass of the lens. In particular, the above Equation (8), together with the upper limit following from Equation (5), yield $10^6 < M/M_\odot < 10^8$ for the distance $D_{LS} \simeq 100$ pc and $10^4 < M/M_\odot < 10^6$ for $D_{LS} \simeq 10$ kpc.

All in all, our analysis of the optical lightcurves of S5 0716 +714 resulting from the TESS observations is consistent with the presence of a massive binary lens containing an IMBH, possibly even a supermassive one, and a much less massive companion (by a factor of $\lesssim 0.01$), located within the host galaxy of the blazar (most likely the central kpc regions).

Until now, no robust detections of any radiative or dynamical signatures for the presence of IMBHs have been reported although several models and scenarios are being discussed in the literature identifying IMBHs with ultraluminous X-ray sources, centers of dwarf galaxies, or centers of globular clusters (for reviews, see Mezcua 2017; Greene et al. 2020). Interestingly, gravitational lensing of background stars by globular clusters in our galaxy has been proposed as a viable method for detecting such IMBHs (Kains et al. 2016, 2018). Globular clusters are, moreover, expected to be more numerous in massive ellipticals than in galaxies such as our Milky Way (Harris et al. 2017; Lim et al. 2020).

We also note in this context recent N -body simulations of globular clusters, which not only reveal the presence of IMBHs but also the fact that most of such IMBHs are in binary systems (see, e.g., Konstantinidis et al. 2013; Leigh et al. 2013). According to our modeling, the IMBH lens in the host galaxy of S5 0716+714 should be accompanied by a much less massive (ratio 1:200) object. This mass ratio is within the range of the *intermediate mass-ratio inspirals*, mergers of which can be probed in the near future directly with the *Laser Interferometer Space Antenna* gravitational-wave detector (see the recent review by Amaro-Seoane et al. 2022), with the conservative estimate of the detection rate of about two per year for the events at high redshifts (Pestoni et al. 2021).

Currently, the only observational clue we have regarding the population of binary black holes approaching the IMBH range is a large number of black hole binaries discovered by the

Advanced LIGO and Advanced Virgo gravitational-wave detectors, with the merger rate density estimated at $23.9^{+14.9}_{-8.6}$ $\text{Gpc}^{-3} \text{yr}^{-1}$ (Abbott et al. 2021), albeit in a lower-mass range of $14\text{--}150 M_\odot$. Until now, no signal for merging of more massive systems has been found in the LIGO/Virgo data (Abbott et al. 2022), and the estimated rate of mergers for binaries with the total masses of $150 M_\odot$ is $\sim 0.08 \text{ Gpc}^{-3} \text{yr}^{-1}$.

We propose that lensing events of the sort we advocate for S5 0716+714 are in fact not uncommon among blazars, but the actual number of lensed sources and lensing events in a single source could be estimated only with a systematic optical monitoring of a larger number of targets with a sufficiently dense sampling. The underlying red noise-type intrinsic variability of blazar jets is the other major obstacle in a robust identification of microlensing events; finding a proper statistical test in this context is not straightforward and relates to a wider open problem of assigning a significance to recurring features in the periodograms of colored-noise time series. The events we have selected in the TESS lightcurve of S5 0716+714 are all very regular, symmetric in shape, and in addition, of a comparable duration. Moreover, the events are found during different flux levels of the source. While these improve our confidence in the lensing hypothesis, we caution the reader that the precise physical mechanisms that lead to intrinsic blazar optical variability have not yet been identified, which makes it difficult to conclusively establish lensing as the cause of the symmetric variations with the available data.

Lastly, we emphasize again an important aspect of the proposed scenario, namely, that we do expect the proposed lensing events in the optical lightcurve of S5 0716+714 to repeat, even though not necessarily in a periodic or even quasiperiodic fashion. Moreover, the orbital modulation of the lens expected to take place on the timescale of years/decades for the lens–source distances < 100 pc and large lens masses $M > 10^4 M_\odot$ offers an exciting possibility for observational tests of the model by detecting a systematic evolution in the magnification patterns for the lensing events (modulo systematic changes in the jet trajectory parameters due to the jet precession).

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Chapter 4

On the Origin of the X-ray Emission in Heavily Obscured Compact Radio Sources



On the Origin of the X-Ray Emission in Heavily Obscured Compact Radio Sources

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Abstract

X-ray continuum emission of active galactic nuclei (AGNs) may be reflected by circumnuclear dusty tori, producing prominent fluorescence iron lines at X-ray frequencies. Here, we discuss the broadband emission of three radio-loud AGNs belonging to the class of compact symmetric objects (CSOs), with detected narrow Fe K α lines. CSOs have newly born radio jets, forming compact radio lobes with projected linear sizes of the order of a few to hundreds of parsecs. We model the radio-to- γ -ray spectra of compact lobes in J1407+2827, J1511+0518, and J2022+6137, which are among the nearest and the youngest CSOs known to date, and are characterized by an intrinsic X-ray absorbing column density of $N_{\text{H}} > 10^{23} \text{ cm}^{-2}$. In addition to the archival data, we analyze the newly acquired Chandra X-ray Observatory and Submillimeter Array (SMA) observations, and also refine the γ -ray upper limits from Fermi Large Area Telescope monitoring. The new Chandra data exclude the presence of the extended X-ray emission components on scales larger than $1''5$. The SMA data unveil a correlation between the spectral index of the electron distribution in the lobes and N_{H} , which can explain the γ -ray quietness of heavily obscured CSOs. Based on our modeling, we argue that the inverse-Compton emission of compact radio lobes may account for the intrinsic X-ray continuum in all these sources. Furthermore, we propose that the observed iron lines may be produced by a reflection of the lobes' continuum from the surrounding cold dust.

Unified Astronomy Thesaurus concepts: Active galactic nuclei (16); Active galaxies (17); Radio lobes (1348); Jets (870); High energy astrophysics (739); Gamma-ray astronomy (628); X-ray astronomy (1810); X-ray active galactic nuclei (2035); Radio galaxies (1343)

1. Introduction

Hard X-ray emission of active galactic nuclei (AGNs) is typically associated with the presence of a corona, a hot plasma in the accretion disk vicinity, where optical/UV disk photons gain energy via the process of Comptonization (Haardt & Maraschi 1991, 1993; Svensson & Zdziarski 1994; Stern et al. 1995; Esin et al. 1997). Throughout the years, a variety of models have been discussed to describe coronas' location and structure, including a lamppost model (Esin et al. 1997; Yuan & Narayan 2004), a hot inner flow (Done & Zycki 1999; Yuan & Narayan 2014), or a jet/base of a jet (Markoff et al. 2005; Dauser et al. 2013). The X-ray continuum is often accompanied by fluorescent iron lines, which are now established as ubiquitous features in the X-ray spectra of both obscured and unobscured AGNs (e.g., Page et al. 2004; Nandra et al. 2007; Ricci et al. 2014). A narrow Fe K α line is produced when X-ray continuum emission reflects off mater located either on the line of sight (e.g., in the broad-line region clouds), or out of the line of sight, e.g., dusty torus or outer regions of an accretion disk. In the latter case, the iron line is accompanied by a Compton reflection component. (e.g., George & Fabian 1991; Matt et al. 1996; Nandra 2006; Bianchi et al. 2008; Shu et al. 2010; Gandhi et al. 2015).

In this paper, we complement the above scenarios by proposing young, compact lobes of radio galaxies and quasars as another source of AGN X-ray continuum emission. These lobes constitute a reservoir of a hot magnetized plasma backflowing from the jet termination shock, where the jet's bulk kinetic energy is converted to the internal energy of the jet particles (e.g., Begelman & Cioffi 1989). Unlike in the case of relativistic beamed jets, the lobes' emission is isotropic in the observers' rest frame. During the earliest phases of the jet's lifetime, the lobes have a size that is smaller or comparable to that of the dusty torus. Consequently, in objects where the primary X-ray continuum originates from radio lobes expanding inside or near a molecular torus, one should expect to observe X-ray reflection and fluorescence features, such as the Fe K α line.

With this picture in mind, we model the broadband spectral energy distributions (SEDs) of three objects: J1407+2827 (also known as OQ+208 or Mrk 668), J1511+0518, and J2022+6137 (2021+614), to quantitatively explore the validity of the above expectation. These objects were selected from a sample of compact symmetric objects (CSOs; a subset of AGNs known for their compact, young radio structures; see, e.g., O'Dea 1998; O'Dea & Saikia 2021; Kiehlmann et al. 2024a, 2024b; Readhead et al. 2024) considered in Sobolewska et al. 2019a based on their distinctive X-ray spectral features derived based on extensive observations by Chandra, XMM-Newton, and NuSTAR (see Sobolewska et al. 2019b, 2023, hereafter S19 and S23 respectively). Key features are a high intrinsic X-ray absorption (equivalent hydrogen column

densities, $N_{\text{H}} > 10^{23} \text{ cm}^{-2}$), as well as a narrow fluorescent Fe K α line and the broad Compton reflection component indicating a reflection from a toroidal obscurer. Such properties make the selected targets rare examples of *radio-loud/X-ray-obscured* AGNs (see Panessa et al. 2016; LaMassa et al. 2023).

The observed radio morphology of these sources, at the milliarcsecond scale, comprises a weak compact core and symmetric lobes, with the innermost projected linear sizes within the 7–25 pc range (An & Baan 2012). These sizes are comparable to the radii of dusty molecular tori in nearby Seyfert galaxies, as measured by the Atacama Large Millimeter Array (ALMA; Combes et al. 2019; García-Burillo et al. 2021). The integrated radio spectra display spectral turnovers at gigahertz frequencies, suggesting excess absorption in the lobes’ radio continuum emission (Wójtowicz et al. 2020, hereafter W20; Kiehlmann et al. 2024a). Additionally, the mid-infrared (IR) colors of these sources, as measured by the Wide-field Infrared Survey Explorer (WISE), place them in the AGN classification, meaning that their observed radiative output at the micron range is indeed dominated by the circumnuclear dust, typical for quasars and luminous Seyfert galaxies (see Kosmaczewski et al. 2020; Nascimento et al. 2022).

To investigate the potential of radio lobes as a source of X-ray continuum in the abovementioned CSOs, we constructed and then modeled the broadband SED using archival radio, IR, optical/UV, and X-ray data. Additionally, we incorporated recent submillimeter data collected with the Submillimeter Array (SMA). Finally, we revisit the Fermi Large Area Telescope (LAT) data, which indicate that our three CSOs are γ -ray quiet (Principe et al. 2020), and derive γ -ray upper limits based on 14 yr of LAT monitoring. Our modeling approach employs a simple dynamical description of the expanding lobes in young radio sources. The lobes’ X-ray emission is produced via inverse-Compton (IC) scattering of various soft photon fields—predominantly those emitted by accretion disks and dusty tori—by ultrarelativistic electrons injected by the jets into evolving compact lobes. All relevant energy losses of the radiating electrons are taken into account and integrated over the entire lifetime of the source (Stawarz et al. 2008, hereafter S08). We show that in the case of all three sources, X-ray emission can be successfully modeled within this scenario.

In addition, we present the new X-ray data collected with Chandra for J1511+0518. These new data enable us to perform an image analysis of the source and investigate the presence of an extended X-ray emission on scales larger than $1''5$ (~ 2.4 kpc). Let us note that, in contrast to our heavily X-ray-absorbed γ -ray quiet CSOs, the three γ -ray loud CSOs detected to date with Fermi-LAT are unobscured in X-rays. These include PKS 1718-649 (Migliori et al. 2016; Siemiginowska et al. 2016), TXS 0128+554 (Lister et al. 2020), and NGC 3894 (Principe et al. 2020; Balasubramaniam et al. 2021), with the observed (isotropic) γ -ray luminosities $L_{0.1-1000 \text{ GeV}} \simeq 1 \times 10^{42}$, 2×10^{43} , and $6 \times 10^{41} \text{ erg s}^{-1}$, respectively (Principe et al. 2021). We also note that the X-ray emission extended beyond the innermost $1''5$ nuclear region from the radio source position has been detected around two of three γ -ray loud CSOs (PKS 1718-649 and NGC 3894, see the details in Siemiginowska et al. 2016; Beuchert et al. 2018; Balasubramaniam et al. 2021, respectively). In both sources, this emission has been modeled as originating from a collisionally ionized plasma with temperatures of $\sim (0.7-0.8) \text{ keV}$, possibly with an additional

Table 1
Chandra Observations of J1511+0518

ObsID	Date	Exposure ^a (s)	Net Counts in Energy Range ^b	
			0.5–7 keV	0.5–1 keV
25457	2022 May 12	20,170	252 ± 16	12 ± 3
25863	2022 May 17	15,700	172 ± 13	11 ± 3
26417	2022 May 18	19,700	224 ± 15	11 ± 3

Notes.

^a ACIS-S3 exposure in seconds.

^b Net counts within the $5 \text{ px} \simeq 2''5$ circular source extraction region centered at the source coordinates.

photoionized plasma component in PKS 1718-649 (Beuchert et al. 2018).

To date, no extended X-ray emission has been spatially resolved around the other two heavily obscured CSOs, as discussed in S19 for J1407+2827 and in Siemiginowska et al. (2016) for J2022+6137. The low number of counts in the archival 2 ks Chandra observation of J1511+0518 did not allow for a meaningful spatial analysis. However, S23 pointed out that including a thermal plasma component with a temperature of $\sim 1 \text{ keV}$ could alleviate the residuals in the soft range of the XMM-Newton and NuSTAR joint spectrum. Since the spatial resolution of the XMM-Newton data was insufficient to determine whether this thermal emission originated within the unresolved AGN region or on scales exceeding several kiloparsecs, we reobserved J1511+0518 with Chandra to perform an image analysis of this source and investigate the presence of the extended emission. Here, we show that in the new 55 ks Chandra observation, we do not detect any extended emission.

This paper is organized as follows. In Section 2, we analyze the new Chandra data for J1511+0518 and show that the X-ray emission of this source is unresolved on arcsec scales. In Section 3, we describe the multiwavelength data used to build the broadband SEDs of J1407+2827, J1511+0518, and J2022+6137. In Section 4, we briefly summarize the applied radiative model for expanding radio lobes, and delineate the modeling procedure. We present our modeling results in Section 5. The main findings derived from the analysis are discussed further in Section 6 and summarized in Section 7.

Throughout the paper, we use modern flat Λ CDM cosmology with $H_0 = 69.3 \text{ km s}^{-1} \text{ Mpc}^{-1}$ and $\Omega_m = 0.287$ (Hinshaw et al. 2013). For J1407+2827, $z = 0.076$ (Stanghellini et al. 1993) gives $D_L \simeq 348 \text{ Mpc}$ and $1.5 \text{ kpc arcsec}^{-1}$. For J1511+0518, its redshift at $z = 0.084$ (Ahn et al. 2012) translates to the luminosity distance $D_L \simeq 388 \text{ Mpc}$ and the scale of $1.6 \text{ kpc arcsec}^{-1}$. Finally, for J2022+6137, $z = 0.227$ (Polatidis & Conway 2003) translates into $D_L \simeq 1,144 \text{ Mpc}$ and $3.7 \text{ kpc arcsec}^{-1}$.

2. Chandra Observations of J1511+0518

We analyzed the recent (2022 May) observations of J1511+0518 with Chandra’s Advanced CCD Imaging Spectrometer (ACIS), consisting of three separate pointings summarized in Table 1.

We used the standard procedures to reduce the data with CIAO v4.14 (Fruscione et al. 2006), starting with the reprocessing performed with the `chandra_repro` script.

The astrometry correction was carried out with respect to the longest observation, ObsID 25457, with the `wavdetect`, `wcs_match`, and `wcs_update` scripts. We modeled the Chandra spectra of J1511+0518 using `Sherpa` (Freeman et al. 2001), and the modeling details are described in Section 2.1. Next, we used the best-fitting spectral model as an input to the point-spread function (PSF) simulations of an on-axis pointlike source using the `ChART` (Carter et al. 2003) and `MARX` (Davis et al. 2012) tools. We compared the simulated PSF with that built based on our observations to assess the presence of extended X-ray emission on scales larger than $\sim 1''$ from the source coordinates (see Section 2.2).

2.1. Spectral Modeling

The source and background energy spectra were extracted separately for each observation of J1511+0518. We chose a circular source extraction region with a radius of 5 ACIS-S pixels⁷, corresponding to $\sim 2.5''$, centered at R.A. = $15^{\text{h}}11^{\text{m}}41.8^{\text{s}}$, decl. = $+5^{\circ}18'10''150$ (J2000). For the background region, we chose a concentric ring with $r_{\text{in}} = 10 \text{ px} \sim 5''$ and $r_{\text{out}} = 20 \text{ px} \sim 10''$, centered at the same coordinates as the source region. The number of net source counts for each observation can be found in Table 1. We binned the spectra requiring the signal-to-noise ratio of 3 in each bin. The maximum bin length was set to 0.2 keV. We performed a joint spectral fit of all three Chandra observations in the 0.5–7.0 keV energy band in order to inform the PSF simulations performed in Section 2.2, as part of the image analysis.

Our simplified spectral model applied to the new Chandra data was based on the joint modeling of the XMM-Newton and NuSTAR data of the source by S23. In particular, our continuum model consisted of a combination of a collisionally ionized plasma, `xsapec`, and a nonthermal emission approximated with a simple `powerlaw` model. We fixed the photon index and the plasma temperature at $\Gamma = 1.8$ and $kT = 1.0 \text{ keV}$ to avoid degeneracy in the fit parameters. Additionally, our model included a Gaussian line, `xsigma`, to account for the narrow Fe K α feature reported in S23, which is also present in the new Chandra data. The rest-frame energy and the width of the line were fixed at 6.4 and 0.02 keV, respectively. Moreover, we have included the Galactic absorption, `xsphabs`, with the equivalent hydrogen column density fixed at $N_{\text{H,Gal}} = 3.29 \times 10^{20} \text{ cm}^{-2}$ (Dickey & Lockman 1990) acting on all the spectral components, and an intrinsic absorption at the redshift of the source, `xsphabs`, with the column density $N_{\text{H,z}}$ acting exclusively on the power-law component.

The final model, `xsphabs*(xsphabs*powerlaw + xsapec + xsigma)`, had four free model parameters, which we linked across the three data sets (corresponding to the three separate pointings) during the fitting procedure. We obtained a satisfactory fit with the `chi2gehrrels`⁸ fit statistic of 43° for 97 degrees of freedom for the source-intrinsic-absorbing column density $N_{\text{H,z}} = (1.28 \pm 0.31) \times 10^{22} \text{ cm}^{-2}$, the normalization of the power-law component, $N_{\text{PL}} = (4.47 \pm 0.048) \times 10^{-5} \text{ ph keV}^{-1} \text{ cm}^{-2} \text{ s}^{-1}$, the normalization of the iron line $N_{\text{Fe}} = (1.43 \pm 0.69) \times 10^{-6} \text{ ph cm}^{-2} \text{ s}^{-1}$, and the normalization of the thermal plasma $N_{\text{APEC}} = (4.47 \pm 0.47) \times 10^{-19} (1+z)^2 \int dV n_e n_{\text{H}} / 4\pi D_L^2$, where dV is the emission volume element, and n_e

and n_{H} are the electron and hydrogen densities, respectively, all in cgs units. All the uncertainties represent 1σ confidence intervals.

The best-fit model returned the 0.5–7.0 keV de-absorbed flux of the power-law component $F_{0.5-7.0 \text{ keV}} \simeq (2.14 \pm 0.14) \times 10^{-13} \text{ erg cm}^{-2} \text{ s}^{-1}$, and the extrapolated 2–10 keV power-law luminosity of $L_{2-10 \text{ keV}} \sim 2.5 \times 10^{42} \text{ erg s}^{-1}$; the equivalent width of the iron line read as $\text{EW} \simeq 0.72_{-0.36}^{+1.12} \text{ keV}$. These values are all comparable within the errors to the corresponding parameter values obtained by S23. However, we note a difference in the complexity of the models applied therein and in our work due to a very limited energy range and photon statistics of the new Chandra data when compared with the joint XMM-Newton and NuSTAR data sets.

2.2. Image Analysis

We simulated the Chandra PSF for all three pointings using the `chart` tool and assuming the best-fit spectral model as presented in Section 2.1 above. We repeated the simulations 50 times for each of the three observations to account for the possibility of considerable differences in each realization of the PSF due to random photon fluctuations. We used the `marx` software to convert the obtained rays to the pseudo-event file with default ASPECT blur 0.07. Finally, we averaged the simulation results for each observation.

We extracted all the detected counts from concentric rings defined by the difference between the outer and inner radii of $\Delta r_{\text{int}} = (0.5 + n)''$, with $n = 0, \dots, 4$, and $\Delta r_{\text{ext}} = (0.5 + 2n)''$ with $n = 2, \dots, 7$. A comparison of the resulting source brightness profile for the merged three pointings and the simulated averaged PSF is presented in Figure 1. As shown, the observed brightness profile is consistent with the emission of a pointlike source within the ACIS spatial resolution, given a relatively large spread observed in various PSF realizations due to random photon fluctuations, especially at larger radii from the source coordinates (see the orange-shaded area in the figure) and the additional uncertainty in the astrometry correction between the three analyzed Chandra pointings.

The lack of the extended emission component in the new Chandra data can be significantly affected by Chandra’s ACIS-S chip degradation. As the extended emission component is expected to originate from thermal hot plasma radiation, it would predominantly appear in the Chandra soft band, which is most affected by the chip degradation. In this context, it is worth noting that the Chandra observations of CSO NGC 3894, which revealed the source’s extended emission, were conducted during Cycle 9. The effective area of the ACIS-S chip has significantly decreased between Cycles 9 and 23, the latter being when J1511+618 was observed. Namely, in soft X-ray band has decreased by ~ 2 orders of magnitude.⁹

3. Multiwavelength Data

3.1. SMA Observations

The SMA was utilized to observe young radio sources and flat spectrum radio quasars under the SMA programs 2018A-S042 and 2018B-S037. These programs were executed multiple times to accommodate a wide range of target locations, and thus J1407+2827 was observed twice (2018 August 8 and 2019 May 2), J1511+0518 was observed twice (2018 July 13

⁷ For ACIS-S, 1 px = $0.492''$.

⁸ <https://cxc.cfa.harvard.edu/sherpa/ahelp/chi2gehrrels.html>

⁹ For the comparison, see https://cxc.cfa.harvard.edu/cgi-bin/prop_viewer/build_viewer.cgi?ea.

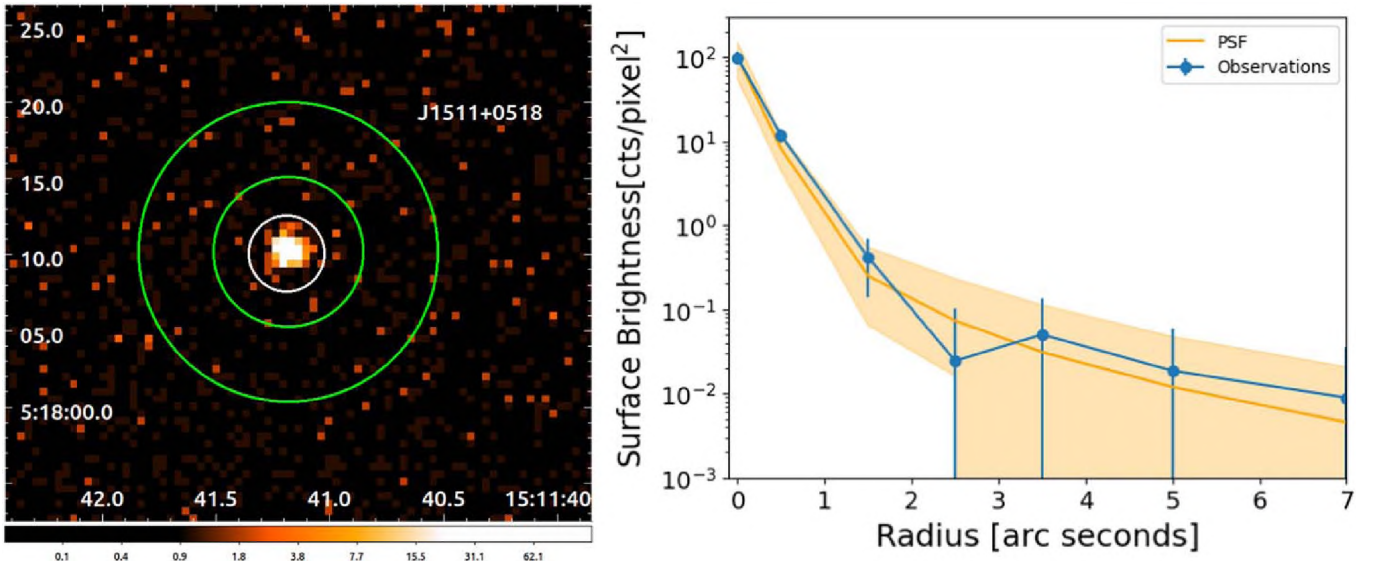


Figure 1. Left: Chandra The merged Chandra ACIS-S image of three J1511+0518 observations. The source and background extraction regions are denoted by a circle with the radius of $r \simeq 2''.5$ (white), and an annulus with $r \simeq (5\text{--}10)''$ (green), respectively. Right: brightness profile of the merged image (blue data points and lines), along with the simulated averaged PSF (orange curve). The orange-shaded region around the average PSF indicates the range (i.e., the minimum and maximum values) observed at a given radius in 150 random PSF realizations.

and August 8), and J2022+6137 was observed once (2018 July 13). For all the observations, complex-valued measurements (*visibilities*) of the sources were bracketed by measurements of nearby strong radio-loud AGNs (often blazars) as amplitude and phase gain calibrators. Data reduction was performed with the SMA inhouse reduction package MIR,¹⁰ including data flagging, initial system temperature calibration, removal of correlator-based amplitude spikes at specific channels, and generation of a broadband continuum channel (8 GHz wide, in each of two sidebands). Complex gain corrections, determined from the calibrator, were applied to the target visibilities, and the flux density scale was set using a solar system standard (Callisto) and a secondary standard, MWC 349a. The calibrated continuum-band visibility data were then vector averaged to determine the flux density of each source. The SMA observations and the resulting fluxes are listed in Appendix A.

3.2. Fermi-LAT Monitoring

The LAT is the main instrument on board the Fermi Gamma-ray Space telescope. It is a γ -ray telescope sensitive at energies ranging from 20 MeV to more than 300 GeV (Atwood et al. 2009).

In this work, we performed a dedicated analysis of the 14 yr of LAT observations (between 2008 August 4 and 2022 August 4) of the three CSOs, following the analysis technique of Principe et al. (2021). We selected P8R3 SOURCE class events (Bruel et al. 2018), in the energy range between 100 MeV and 1 TeV, from regions of interest (ROIs) of $15^\circ \times 15^\circ$ centered at the positions of each selected source. The value of the low-energy threshold is motivated by the large uncertainties in the arrival directions of the photons below 100 MeV (Principe et al. 2018), leading to possible confusion between pointlike sources and the Galactic diffuse component.

The analysis, consisting of model optimization, localization, spectrum, and variability study, was performed with *Fermipy* version 1.0.1¹¹ (Wood et al. 2017), a Python package that facilitates analysis of the LAT data with the Fermi Science Tools.

We created the count maps with a pixel size of $0''.1$, and we excluded all γ -rays with a zenith angle larger than 95° in order to limit the contamination from secondary γ -rays from Earth's limb (Abdo et al. 2009). Following the analysis technique reported in Principe et al. (2020, 2021), we made an even harder cut at low energies (<300 MeV) by reducing the maximum zenith angle ($<85^\circ$) and by excluding event types with the largest PSF (PSF0).¹²

We used the P8R3_SOURCE_V3 instrument response functions. The selected spectral model included all pointlike and extended LAT sources located at a distance $<20^\circ$ from the investigated CSOs, as listed in the Fourth Fermi-LAT Source Catalog (4FGL-DR2; Abdollahi et al. 2020), as well as the Galactic diffuse and the isotropic emission models¹³ adopted to compile the 4FGL-DR2. We first optimized the model for a given ROI, then we searched for additional new sources not included in 4FGL-DR2, and finally, if the source was significantly detected ($>3\sigma$ significance), we relocalized the source.

During the model fitting, we left the normalization of the isotropic and Galactic diffuse backgrounds, as well as the spectral parameters of the sources within 5° of our targets, free to vary. For sources at a distance between 5° and 10° , only the normalization was fitted; we fixed the parameters of all the sources within the ROIs at larger angular distances from our

¹¹ <http://fermipy.readthedocs.io/en/latest/>

¹² A measure of the quality of the direction reconstruction is used to assign events to four quartiles. γ -rays in Pass 8 data can be separated into four PSF event types: 0, 1, 2, and 3, where PSF0 has the largest PSF and PSF3 has the best.

¹³ <https://fermi.gsfc.nasa.gov/ssc/data/access/lat/BackgroundModels.html>, in particular the diffuse model `gll_iem_v07.fits` and the isotropic component `iso_P8R3_SOURCE_V3_v1.txt`.

¹⁰ <https://lweb.cfa.harvard.edu/~cqj/mircook.html>

Table 2
Parameters of the Expanding Radio Lobe Model

Parameter	Description	Values
LS	Radio linear size (half of the hotspot–hotspot distance)	Table 3
ν_{peak}	Radio peak frequency (spectral turnover frequency)	Table 3
ν_h	Jet advance velocity (half of the hotspot–hotspot separation velocity)	Table 3
ν_{IR}	Characteristic frequency of the torus emission	Table 3
L_{IR}	IR luminosity of the circumnuclear torus	Table 3
n_0	ISM number density in the central regions of the host galaxy	Table 4
ν_{UV}	Characteristic frequency of the disk UV emission	2.45×10^{15} Hz (fixed)
L_{UV}	Luminosity of the accretion disk UV emission	Table 5
η_e	Ratio of the radiating electrons' energy density and the total lobes' pressure	Table 5
η_B	Ratio of the magnetic field energy density and the total lobes' pressure	Table 5
L_j	Jet bulk kinetic power	Table 5
s_1	Low-energy slope of the injected electron population $Q_e(\gamma)$	Table 5
s_2	High-energy slope of the injected electron population $Q_e(\gamma)$	Table 5
γ_{min}	Minimum Lorentz factor of the electron energy distribution $Q_e(\gamma)$	3.0 (fixed)
γ_{br}	Break Lorentz factor of the electron energy distribution $Q_e(\gamma)$	Table 5
γ_{max}	Maximum Lorentz factor of the electron energy distribution $Q_e(\gamma)$	$100 m_p/m_e$ (fixed)

Note. The top part of the table contains parameters measured directly or estimated from the radio and IR data.

targets to their 4FGL values. Spectral fits were performed over the energy range from 100 MeV to 1 TeV.

None of the selected targets were significantly detected ($>3\sigma$ significance) in the high-energy γ -ray range with LAT. In particular, all three sources present a source significance compatible with 0σ . In Table A1 of Appendix A, we provide the resulting upper limits.

3.3. Archival Data

The radio, IR, and optical/UV data for J1407+2827, J1511+0518, and J2022+6137 used for SED modeling were acquired from the NASA/IPAC Extragalactic Database (NED). The relevant references are listed in Table A3 of Appendix A. For the data within the near-IR-to-UV range, dereddening was performed using the extinction law $A_\lambda = A_V a(\lambda) + (A_B - A_V) b(\lambda)$, with the coefficients $a(\lambda)$ and $b(\lambda)$ as given in Cardelli et al. (1989), and the B -band and V -band Galactic extinction values A_B and A_V as given in the NED following Schlafly & Finkbeiner (2011). Source-intrinsic reddening was not taken into account.

In addition, we also collected de-absorbed fluxes corresponding exclusively to the power-law emission components, obtained via spectral modeling of the broadband X-ray data collected previously with Chandra, XMM-Newton, and NuSTAR by S19 (for J1407+2827) and S23 (for J1511+0518 and J2022+6137). These power-law components represent the direct intrinsic X-ray continuum emission of our sources.

4. Broadband SED Modeling

4.1. The “Expanding Lobes” Model

We conducted a modeling analysis of the broadband SEDs for the three selected CSOs, to probe the origin of their X-ray continuum *power-law* emission. Based on the modeling results, we also investigated the origin of the X-ray reflection components, as well as the apparent γ -ray quietness of the sources. For this purpose, we explored the model of expanding young radio lobes in AGN, put forward by S08. In this section, we briefly summarize the model assumptions and model

parameters (see Table 2). We refer the readers to S08 for further details (see also Ostorero et al. 2010, hereafter O10).

The model evolves a self-consistent set of equations that describe both the expansion of light relativistic jets and their lobes in a dense ambient medium (see Begelman & Cioffi 1989), and the evolution of ultrarelativistic electrons injected into the lobes by the jets, taking into account adiabatic energy losses and radiative cooling. In this framework, a pair of relativistic jets with total kinetic energy L_j , propagates through the inner segment of the host galaxy, characterized by a constant density of $\rho \simeq m_p n_0$. The jets' thrust is balanced by the ram pressure of the ambient medium spread over some area larger than the jet terminal cross-section (due to the effects of a jet precession), so that the jet advance velocity, ν_h , is at most mildly relativistic. Moreover, the model assumes that all the bulk kinetic energy transported by the jets, is converted at the jet termination shock into the internal energy of the jet particles and the jet magnetic field, amounting together to the lobes' internal pressure p . The resulting pressure-driven expansion of the lobes in the direction perpendicular to the jets' axis, is supersonic, even though subrelativistic.

Ultrarelativistic electrons and positrons (hereafter referred to as *electrons* for simplicity) injected into the lobes, are described by the initial energy distribution $Q(\gamma, t)$, which in our case is assumed to be constant in time and a broken power-law function of electron energy $E_e \equiv \gamma m_e c^2$, namely, $Q_e(\gamma) \propto \gamma^{-s_1}$ for $\gamma_{\text{min}} < \gamma < \gamma_{\text{br}}$, and $Q_e(\gamma) \propto \gamma^{-s_2}$ for $\gamma_{\text{br}} < \gamma < \gamma_{\text{max}}$. This distribution is evolved in expanding lobes and integrated over the entire source lifetime, self-consistently taking into account adiabatic energy losses, synchrotron emission, and the IC scattering on different photon populations. Those seed photon populations include photons originating from the lobes' synchrotron emission, IR photons originating from the obscuring torus, and UV photons associated with the radiation of the accretion disk. Here, we do not consider Comptonization of the starlight emission of host galaxies, noting that this process is of minor relevance in the case of the youngest and most compact (LS < 100 pc) CSOs (see the related discussions in S08 and O10). Likewise, the synchrotron self-Compton process is in all the analyzed cases negligible as well.

Table 3

Parameter Values Assumed in Our Modeling Based on the Radio and IR Data

Source	LS ^a (pc)	ν_{peak}^a (GHz)	ν_h^b (c)	ν_{IR}^c (10^{13} Hz)	L_{IR}^c (10^{45} erg s ⁻¹)
J1407+2827	9.9	4.0	0.13	1.5	0.7
J1511+0518	7.3	9.8 ^d	0.08	3.0	0.8
J2022+6137	24.5	8.4	0.18	1.0	1.3

Notes.^a References to the high-resolution radio data: Stanghellini et al. (1998), Tschager et al. (2000), Orienti & Dallacasa (2008); Wu et al. (2013).^b See An & Baan (2012) and references therein.^c For references to the IR data, see Table A3.^d Mean of ν_{peak} for both lobes.

We assume that the IR and UV spectra of the soft photons are monochromatic, with luminosities (L_{IR} and L_{UV}) and frequencies (ν_{IR} and ν_{UV}), corresponding to the bolometric luminosities and characteristic frequencies of the blackbodies, which match the observed data points in the respective bands. We note that such an approximation does not significantly affect the calculated high-energy IC emission of the lobes, given the broad energy range of the lobes' electron energy distribution, $Q_e(\gamma)$.

The radio continuum is modeled as synchrotron emission of the lobes, assuming that the low-frequency break in the CSO radio spectrum around the observed peak frequencies ν_{peak} , is due to the free-free absorption on the interstellar medium (ISM) clouds engulfed by the expanding lobes, and photo-ionized by the nuclear emission (Begelman 1999).

The lobes' magnetic field and electron energy densities are parameterized as $U_B = \eta_B p$ and $U_e = \eta_e p$, respectively, where p is the total lobes pressure, and the proportionality factors η_B , $\eta_e < 1$ are assumed to be constant during the source evolution (see the discussion in S08). The linear size of the radio structure, measured from the radio core position to the jet termination shock (*hotspot*), is denoted below as LS.

4.2. The Modeling Procedure

Among the 16 modeling parameters introduced in the previous section and summarized in Table 2, three could be measured directly from high-resolution radio monitoring data available for all the analyzed sources. These included radio linear sizes, LS, peak frequencies, ν_{peak} , and jet advance velocities, ν_h , all as given in Table 3 with the corresponding references. Two other modeling parameters, namely, the IR torus luminosity L_{IR} and its characteristic frequency ν_{IR} , could be estimated based on the integrated IR data, also as summarized in Table 3.

We have constrained ranges of the ambient medium density based on the intrinsic equivalent hydrogen column density values, N_{H} , provided for each source based on the detailed analysis of the X-ray spectrum (see Table 4). In particular, we estimated lower limits for the ambient medium density utilizing the fact that all three CSOs appeared pointlike in the Chandra images, and assuming that the X-ray absorbing matter is distributed in a circumnuclear region with a characteristic spatial scale corresponding to $\ell = 1''.5$ at the redshift of a source, a typical radius of a Chandra source extraction region; this gave $n_{\text{low}} \simeq N_{\text{H}}/\ell$. Upper limits for the ambient medium density, on the other hand, were found assuming that the characteristic spatial scales of the X-ray absorbers were

Table 4Adopted Range of the Ambient Medium Density, n_0

Source	N_{H} (10^{23} cm ⁻²)	n_{low} (cm ⁻³)	n_{int} (cm ⁻³)	n_{high} (cm ⁻³)
J1407+2827	4.4	70	2000	20,000
J1511+0518	0.9	15	400	4000
J2022+6137	3.5	25	200	4000

Note. References for N_{H} estimates: S19, S23.

comparable to the sources' radio linear sizes, namely, $n_{\text{high}} \simeq N_{\text{H}}/\text{LS}$. Because the range spanned by the n_{low} and n_{high} values is relatively wide, ~ 2.5 dex, we have also considered intermediate (in orders of magnitude) densities, n_{int} , all as given in Table 4.

Since the observed IR-to-UV data points were not corrected for the source-intrinsic reddening, the UV emission of an accretion disk was treated as a model-free parameter, except that we fixed its comoving characteristic frequency at $\nu_{\text{UV}} = 2.45 \times 10^{15}$ Hz, following O10. Generally, broadband quasar observations suggest that $L_{\text{IR}} \simeq \eta_{\text{h}} L_{\text{UV}}$, with $\eta_{\text{h}} \gtrsim 0.5$ (see in this context, e.g., Rałowski et al. 2024), albeit with a rather wide spread. Therefore, in this work, we considered a range of L_{UV} luminosities corresponding to $\eta_{\text{h}} = 0.1$ – 1 . From this range, we chose the values of L_{UV} that resulted in a good match of the high-energy spectra with the given X-ray and γ -ray observational constraints.

Regarding the pressure content of the lobes, in our modeling, we investigated two sets of the η_B and η_e parameters. The first set, $\eta_B = 0.3$ and $\eta_e = 3.0$, corresponds to the scenario in which ultrarelativistic electrons account for the bulk of the lobes' pressure, $p_e \equiv U_e/3 \sim p$, where p stands for the total pressure, while the electron-to-magnetic field energy density ratio is of the order of 10, $U_e/U_B \simeq 10$ (a typical value emerging from modeling relativistic jets in blazar sources; e.g., Ghisellini et al. 2010). The second set, $\eta_B = 0.75$ and $\eta_e = 0.75$, corresponds to the energy equipartition between lobes' ultrarelativistic electrons and magnetic field, $U_e/U_B \simeq 1$, in which case $p_e \sim p/4$.

For a given combination of n_0 , L_{UV} , η_e , and η_B , the model synchrotron (radio) and the IC (X-ray and γ -ray) emission continua are set by the jet bulk kinetic power, L_j , and the spectral shape of the injection function, $Q_e(\gamma)$. The L_j parameter predominantly controls the normalization of the emission continua. However, it also affects the spectral shape. For example, for a given source size, LS, and advance velocity, ν_h , the larger the L_j parameter, the higher the energy density of the lobes' magnetic field, which leads to an enhanced synchrotron cooling and therefore to softer high-energy spectra of the evolved electron population. Thus, in our modeling procedure, for each considered combination of the n_0 , L_{UV} , η_e , and η_B parameters, we adjust the corresponding L_j and $Q_e(\gamma)$ values until we obtain a satisfactory match to the radio data points, and at the same time, the X-ray power-law continua, paying attention also to the upper limits in the high-energy γ -ray frequency range.

To reduce the number of the free model parameters, we fixed the minimum and maximum electron Lorentz factors as $\gamma_{\text{min}} = 3.0$, and $\gamma_{\text{max}} = 100 m_p/m_e$, respectively. This choice does not affect the model spectra as long as $s_1 < 2.0$ and $s_2 \gg 2.0$, which, in fact, is the case for all three targets (see Table 2 for a description of the parameters).

5. Results of SED Modeling

We successfully employed the expanding radio lobe emission scenario to model the broadband SED of three CSOs: J1407+2827, J1511+0518, and J2022+6137. For the relevant parameter space of the model we employ, the X-ray emission continuum is produced through the IC scattering of both the torus (IR) and disk (UV) photons, while the γ -ray emission is generated exclusively through the IC scattering of the UV emission. Since the torus luminosity is constrained by the observed fluxes at IR wavelengths, for a given set of model parameters, a particular L_{IR} value sets a strict lower limit on the X-ray emission of the lobes. The disk luminosity, on the other hand, is a free parameter of the model. For large enough values of L_{UV} , the IC scattering of this photon field may dominate over the IC scattering of the torus IR photons, increasing the overall level of the lobes' X-ray continuum. Yet, since high-energy γ -rays are produced via the IC scattering of the UV photons, the Fermi-LAT upper limits set, in turn, an upper limit on the allowed value of L_{UV} . The high-energy emission output is additionally moderated by the ISM density and the U_B/U_e ratio. Therefore, the X-ray and γ -ray emission is controlled by the interplay between L_{UV} , n_0 , and U_B/U_e , allowing different sets of their values to reproduce observed X-ray radiation and follow the γ -ray flux upper limits.

The free model parameters related to the spectral shape of the electron injection function, $Q_e(\gamma)$, are tightly constrained by the observed shape of the radio continua, with the data in the submillimeter range being of particular importance. Specifically, the shapes of the radio continua require a broken power-law electron injection function with $s_1 < s_2$ with s_1 in the ~ 1.3 – 1.6 range. The steep high-frequency segments of the radio continua require that the high-energy slope of the electron energy distribution, s_2 , is in the range of ~ 4.7 – 5.8 , which is significantly softer than $s_2 = 2$ and corresponds to a standard diffusive shock acceleration. We revisit this finding in the discussion Section 6 in the context of the γ -ray quietness of the analyzed targets.

When it comes to the low-frequency segments of the observed radio continua of the discussed CSOs, we note that the applied scenario of *free-free absorption on engulfed clouds* performs extremely well for J1407+2827 and J1511+0518, but fails in reproducing the radio spectrum of J2022+6137. In particular, in J2022+6137 we observe an excess low-frequency radio emission over the model spectra for all combinations of the model-free parameters.

5.1. J1407+2827

In the case of J1407+2827, a wide range of the free model parameters under the equipartition condition $U_e/U_B \simeq 1$ could account for the broadband SED of the source. For sub-equipartition models with $U_e/U_B \simeq 10$ (i.e., $\eta_e = 3.0$ and $\eta_B = 0.3$), the X-ray emission was generally exceeded for any model with the density values below n_{int} . The corresponding model solutions are presented in Figure 2, and summarized in Table 5.

The broadband SED of J1407+2827 was previously discussed in the framework of the expanding radio lobe scenario by O10. However, they explored only the sub-equipartition models with $U_e/U_B \simeq 10$, and with a relatively low-ambient medium density of $n_0 \simeq 0.1 \text{ cm}^{-3}$. Moreover, the X-ray spectrum considered in O10, based on the analysis of the

initial XMM-Newton observations (Guainazzi et al. 2004), was subject to large uncertainties (~ 2 orders of magnitude at $\sim 0.3 \text{ keV}$, 1 order of magnitude at $\sim 10 \text{ keV}$), preventing robust constraints on the model parameters. In our analysis, we utilized the X-ray continuum and X-ray absorption constraints following from the modeling of an updated broadband X-ray data set (S19). As a result, we probed significantly higher ISM density values, 70 – $2 \times 10^4 \text{ cm}^{-3}$. Finally, we added the new submillimeter and γ -ray constraints. The main difference between our modeling and the one presented in O10 is in the jet kinetic luminosity L_j , which we found to be lower by 1–2 orders of magnitude, while the parameters of the electron injection function and the soft photon fields were in rough agreement. In particular, the value $s_2 = 5.8$ is in agreement with the value reported in O10 ($s_2 = 5.6$).

5.2. J1511+0518

The SED and best model are shown in Figure 3. Models that could match the observed SED in both the radio and the X-ray domains were the equipartition models $U_e/U_B \simeq 1$ (i.e., $\eta_e = \eta_B = 0.75$), with relatively low UV luminosities, $L_{\text{UV}} \sim (1.0$ – $3.5) \times L_{\text{IR}}$, and a relatively high ISM density, within the n_{int} – n_{high} range, as presented in Figure 3, and summarized in Table 5. Models with lower ISM densities, $n_0 < n_{\text{int}}$, or departing from the equipartition condition, $U_e > U_B$, generally overpredicted the X-ray flux, unless the disk UV luminosity was set to be very low, $L_{\text{UV}} \ll L_{\text{IR}}$, the possibility of which we deem as unlikely.

5.3. J2022+6137

J2022+6137 is characterized by an enhanced level of the intrinsic X-ray continuum emission when compared to J1407+2827 and J1511+0518. The 0.3 – 40 keV luminosity of the source, $\sim 10^{44} \text{ erg s}^{-1}$, could not be obtained within the framework of the applied scenario for high values of the ambient density, regardless of the equipartition ratio U_e/U_B , without invoking an additional, phenomenological X-ray component. On the other hand, the ISM density at the level approaching n_{low} and an equipartition parameter of $1 \lesssim U_e/U_B \lesssim 10$, allowed us to match the intrinsic X-ray power-law continuum with the lobes' IC emission without violating the Fermi-LAT upper limits. Interestingly, for J2022+6137, the γ -ray upper limits were more restrictive than the X-ray constraints. An increase in L_{UV} or a decrease in n_0 violated the γ -ray limits before the X-ray limits. The results of the modeling for this source are presented in Figure 4, and summarized in Table 5.

6. Discussion and Conclusions

In this work, we have studied the broadband emission of three highly X-ray obscured active galaxies, with a CSO radio classification, J1407+2827, J1511+0518, and J2022+6137, each featuring young, compact radio lobes with linear sizes up to 25 pc . We demonstrated that the IC scattering of the IR and UV photons (supplied by a dusty torus and an accretion disk, respectively) within the expanding radio lobes provides a plausible source of intrinsic X-ray continuum emission on the levels consistent with the observational constraints.

For the two analyzed CSOs, J1407+2827 and J1511+0518, certain models could be ruled out because they overestimated the intrinsic X-ray emission, while none of the probed models

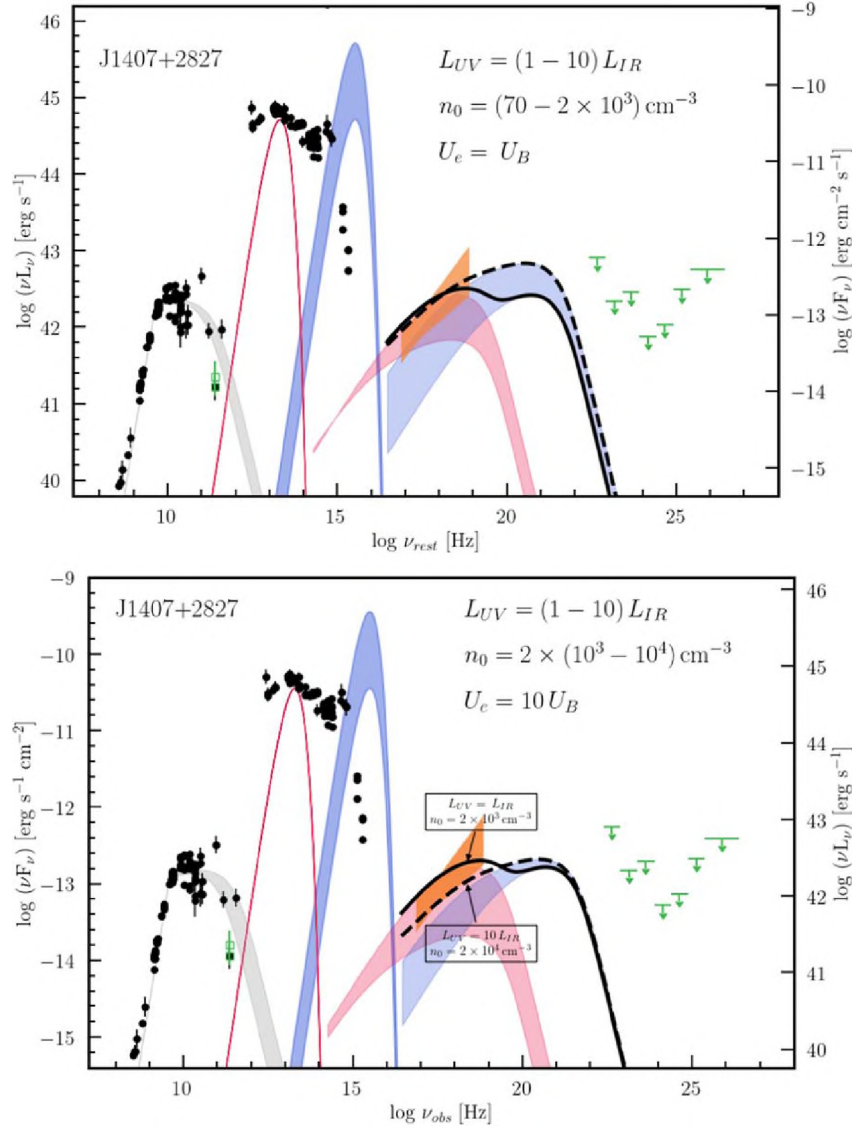


Figure 2. The SED and the best-matching models for J1407+2827, corresponding to the model-free parameter values/ranges as summarized in Table 5 and the figure legend. The archival radio and radio-to-UV data are denoted in the figure by black circles. The new SMA data are plotted as open green squares. The orange bow tie represents the intrinsic de-absorbed X-ray power-law constraints. Finally, the green arrows illustrate the Fermi-LAT upper limits. The model components include the lobes' synchrotron emission (gray-shaded area), the torus IR emissions (red curve), the disk UV emission (dark-blue shaded area), the IC emission off the disk UV photons (light-blue-shaded area), and the IC emission off the torus IR photons (light-red-shaded area). Upper panel: radio lobes in equipartition, lower panel: $U_e = 10U_B$.

Table 5
Modeling Results

Source	U_e/U_B	n_0	L_{UV} (10^{45} erg s $^{-1}$)	L_j (10^{42} erg s $^{-1}$)	s_1	s_2	γ_{br} (m_p/m_e)
J1407+2827	1	$n_{low}-n_{high}$	0.7–7.0	17.5–5.0	1.6	5.8	1.0–0.8
	10	$n_{int}-n_{high}$	0.7–7.0	4.0–5.0	1.6	5.8	0.8
J1511+0518	1	$n_{int}-n_{high}$	1.0–3.5	8.0–6.0	1.6	5.2	1.0
J2022+6137	1	$n_{low}-n_{high}$	13	200–40	1.3	4.7–4.8	1.0
	10	$n_{low}-n_{high}$	1.3–13	100–27	1.3	4.7	1.0

were found to underestimate the intrinsic X-ray emission or to violate the γ -ray upper limits. Models accounting for the comparatively high 0.3–40 keV luminosity of J2022+6137 ($\sim 10^{44}$ erg s $^{-1}$), and at the same time complying with its γ -ray constraints, were characterized by a low-ambient medium density and high-ultraviolet luminosity, $L_{UV} = 10L_{IR}$, in the

case of lobes in the equipartition. A range of n_0 and L_{UV} values, consistent with the lower end of the range depicted in Figure 4, was allowed for the case of the lobes in sub-equipartition. Since the high-density models underestimated the X-ray emission of J2022+6137, we note that an additional phenomenological X-ray continuum component may be present in the SED of

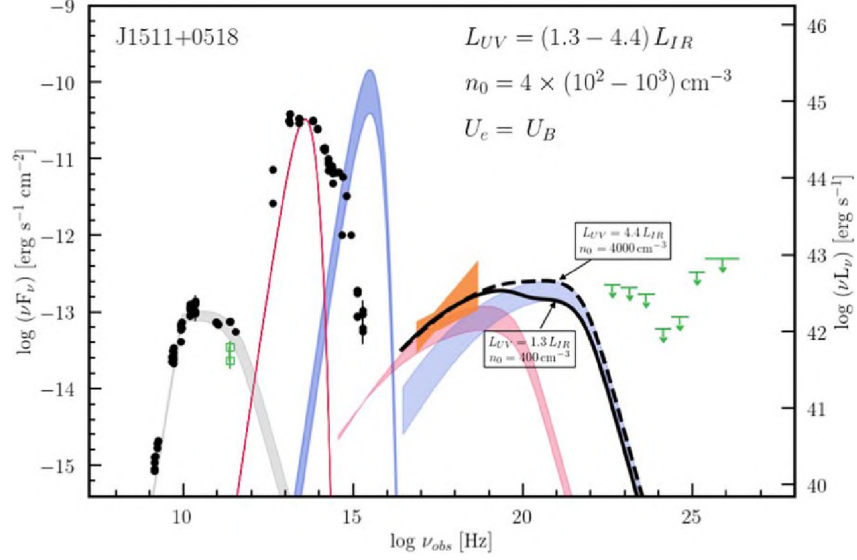


Figure 3. The SED and the best-matching models for J1511+0518 with radio lobes in equipartition. Data and model components are marked the same as in Figure 2.

J2022+6137. In Figure 5, we show an example of a model consisting of the emission of the expanding radio lobes with $n_0 = 4 \times 10^3 \text{ cm}^{-3}$, and a phenomenological cutoff power-law component with a photon index of 1.45 (the best-fit photon index of the intrinsic de-absorbed X-ray power law in S19). The photon index of 1.45 is within the range observed in AGNs and X-ray binaries (Yang et al. 2015). This additional component, with the 2–10 keV luminosity of $\sim 10^{44} \text{ erg s}^{-1}$, could originate, e.g., due to a compact lamppost corona very close to the black hole or due to a complex structure of the jetted outflow (see Migliori et al. 2014 and Sobolewska et al. 2022 in the context of young radio sources).

Apart from its high X-ray luminosity, J2022+6137 also stands out in our sample because of an excess low-frequency radio emission compared to the model spectrum for any combination of the free model parameters. This could signal the presence of an additional, extended (on arcsec scales) radio component (see in this context in Kharb et al. 2010). It is also possible that the scenario of the free-free absorption on engulfed clouds is not applicable in the case of J2022+6137.

For all three sources, we compared the resulting L_{UV} and L_j values to those estimated in W20: $31 \times 10^{42} \text{ erg s}^{-1}$ (J1407+2827), $L_j = 10.5 \times 10^{42} \text{ erg s}^{-1}$ (J1511+0518), $220 \times 10^{42} \text{ erg s}^{-1}$ (J2022+6137); and $L_{UV} = 4.3 \times 10^{45} \text{ erg s}^{-1}$ (J1407+2827), $1.5 \times 10^{45} \text{ erg s}^{-1}$ (J1511+0518), and $1.5 \times 10^{45} \text{ erg s}^{-1}$ (J2022+6137). Due to the high absorption of these CSOs, W20 estimated L_{UV} through the SED for J1511+0518, (following the Trichas et al. 2013 method), based on the $12 \mu\text{m}$ and O III lines, respectively, for J1407+2827 and J2022+6137 (Wu 2009; Kosmaczewski et al. 2020). The inferred L_{UV} values from our analysis are in agreement with those presented by W20. The jet power resulting from our modeling is also roughly in agreement with the values estimated by W20. The differences, which are up to a factor of a few, could result from the different assumptions about lobe geometry and the strict energy equipartition assumed in W20. Note in this context that the emerging total jet power for J1407+2827 and J1511+0518 are $\leq 10^{44} \text{ erg s}^{-1}$, below the threshold for effective confinement of the lobes within the ISM. Therefore, they are expected to be trapped for a relatively long

time before they expand outside of the host galaxies (Mukherjee et al. 2016, 2017), whereas the total jet power of J2022+6137 for some models is above this threshold.

Modeling the radio continua, including newly acquired SMA data, indicated that the electron energy spectrum has to be of a broken power-law form, with $s_1 \leq 2$ below $\gamma \sim m_p/m_e$, and $s_2 \gg 2$ at higher energies. Fermi-LAT observations reinforced this conclusion because with no sharp break in the electron energy distribution, $s_2 - s_1 > 2$, the current LAT upper limits for the source fluxes would be exceeded. Moreover, we emphasize that this broken power-law form of the energy distribution corresponds to the electron injection spectrum. Thus, it described the electron spectrum formed at the acceleration site, namely, the jet termination shock, and not the spectrum modified by the subsequent evolution within the expanding and radiatively cooling lobes. This finding aligns with the scenario proposed by Stawarz et al. (2007), where a broken power-law electron energy distribution with a critical break energy around m_p/m_e represents two distinct regimes of electron acceleration at mildly relativistic, proton-dominated shocks with perpendicular magnetic field configurations. The main hypothesis is that in a proton-mediated shock, only electrons with energies exceeding those of cold protons ($\gamma > m_p/m_e$) can undergo diffusive shock acceleration. This process is expected to yield steep electron spectra (energy indices $\gg 2$) at mildly relativistic perpendicular shocks. Electrons with lower energies ($\gamma < m_p/m_e$) are likely accelerated by different mechanisms, possibly in the near upstream of the shock (as discussed and referenced in Stawarz et al. 2007).

In Figure 6, we compare the s_2 parameter values emerging from our modeling of the three heavily X-ray obscured CSOs, with those obtained for X-ray unobscured CSOs by O10 (excluding J1407+2827 and Sobolewska et al. 2022, PKS 1718-649) using the same expanding radio lobe model. We plot s_2 as a function of the projected linear radio size, LS, and as a function of the source-intrinsic-absorbing column density, N_H . It can be seen that the s_2 index correlates with N_H , and that the most compact CSOs, with LS $\lesssim 25 \text{ pc}$, have softer s_2 indices than CSOs with LS $\gtrsim 100 \text{ pc}$.

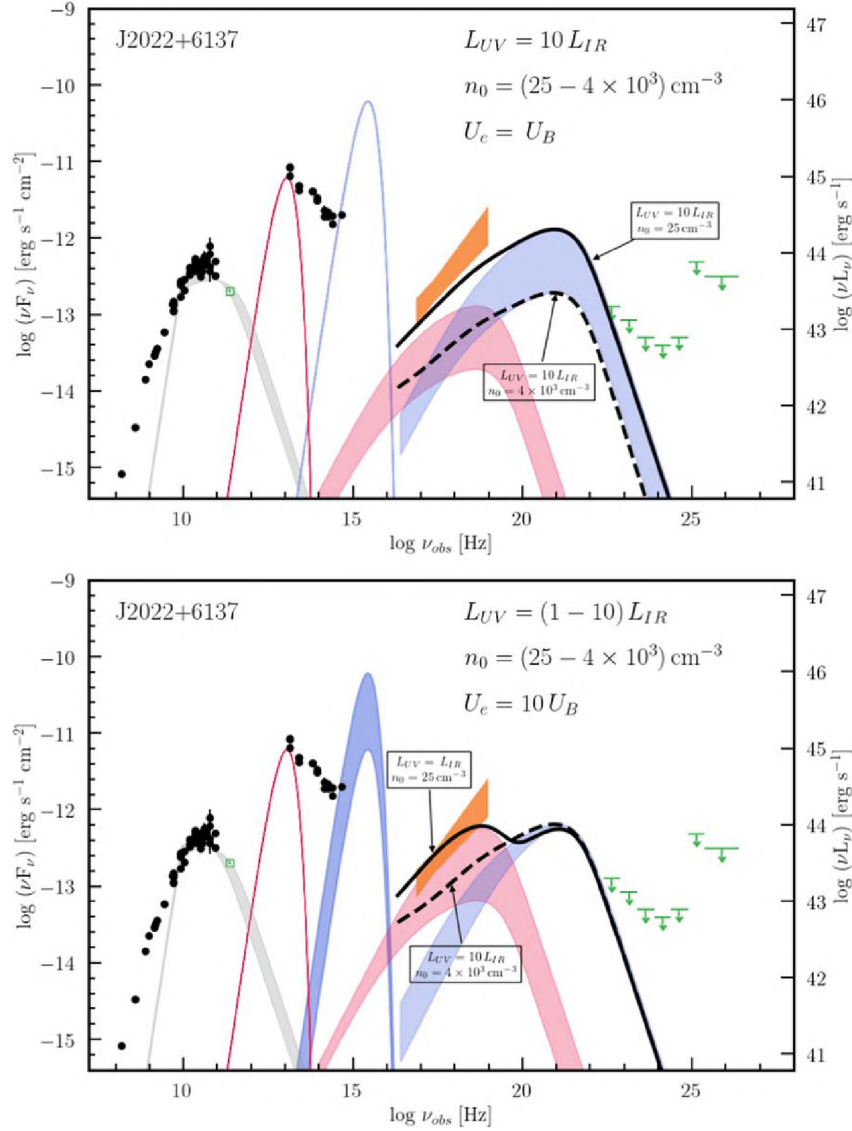


Figure 4. The SED and the best-matching models for J2022+6137. Upper panel: radio lobes in equipartition, lower panel: $U_e = 10U_B$. Data and model components are marked the same as in Figure 3.

In our modeling, the absorbing column density guides the values of the ambient medium density into which the newly born CSO jets propagate, $N_H \propto n_0$, the s_2 – N_H correlation observed in our small sample agrees very well with the s_2 – n_0 trend reported in the model data for the evolved radio galaxies by Wójtowicz et al. (2021). This allows us to draw the same conclusion as Wójtowicz et al., namely that the efficiency of the electron acceleration at mildly relativistic termination shocks of AGN jets decreases with increasing density of the ambient medium. Moreover, very soft values of s_2 may explain the apparent quietness of compact and heavily obscured CSOs in the γ -ray range since it is the high-energy slope of the electron energy distribution that effectively suppresses the lobes' IC emission at GeV photon energies for a given total jet kinetic power and a given disk UV luminosity (compare to the case of the γ -ray detected PKS 1718-649 discussed in Sobolewska et al. 2022).

The ISM number densities resulting from our modeling are relatively high, orders of magnitude higher than the fiducial

value of $n_0 = 0.1 \text{ cm}^{-3}$ considered in the analogous modeling of CSOs by O10. However, all the sources analyzed in O10, except J1407+2827, have linear radio sizes exceeding 100 pc, and one should expect lower-ambient medium densities for them when compared to the three particularly compact/young targets discussed here. Indeed, assuming a general scaling of n_0 with the square of the distance from the black hole, as appropriate for the winds from black hole accretion flows (see, e.g., Cui et al. 2020), gas number densities at the level of $\sim 10^3 \text{ cm}^{-3}$ at 10 pc scale, would correspond to $\sim 0.1 \text{ cm}^{-3}$ at the distance of 1 kpc.

We can also compare our estimates of the ISM density with the estimates obtained for galaxies with a radio-loud TDE. Cendes et al. (2023) presented a collection of ISM density profiles in TDE hosts, constructed assuming an external shock model (see, however, in this context, Pasham & van Velzen 2018). For the distance scale corresponding to 1–25 pc (linear radio size of the lobes in our sources), expressed in Schwarzschild radii (10^5 – 10^6 Schwarzschild radii

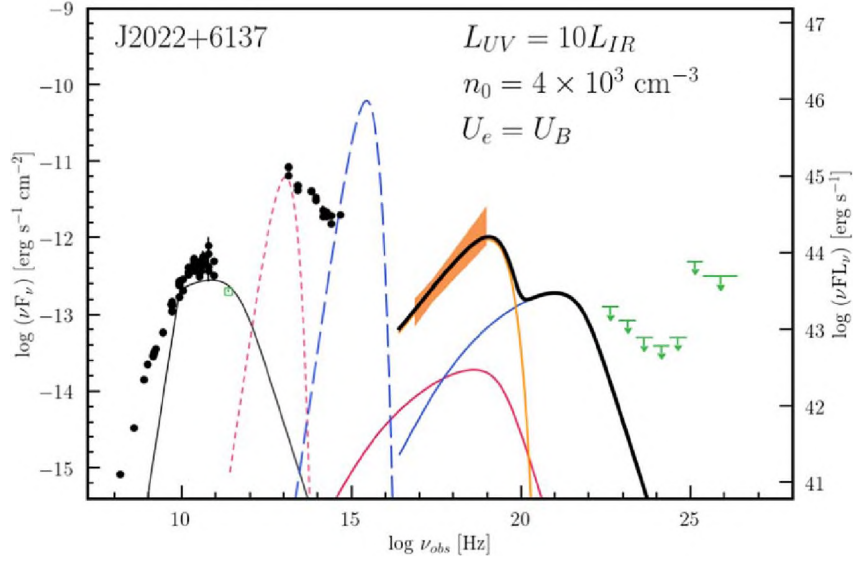


Figure 5. An example of the expanding radio lobe model solution applied to the SED of J2022+6137, which underestimates the most the intrinsic X-ray emission (see the plot label and Table 5 for model parameters). Data, intrinsic X-ray emission and Fermi-LAT upper limits are plotted the same as in Figure 4. An additional, phenomenological X-ray component with a photon index of 1.45 and a 2–10 keV luminosity of $\sim 9.95 \times 10^{43} \text{ erg s}^{-1}$ is required (orange line) to match the intrinsic X-ray emission, and it dominates the X-ray band. The model components include the lobes' synchrotron emission (gray-solid line), torus IR emissions (red-short-dashed line), disk UV emission (blue-long-dashed line), IC emission of the disk seed photons (blue solid line), IC emission of the torus seed photons (red solid line). The thick solid black line corresponds to the sum of the IC components from the lobes and the additional X-ray component.

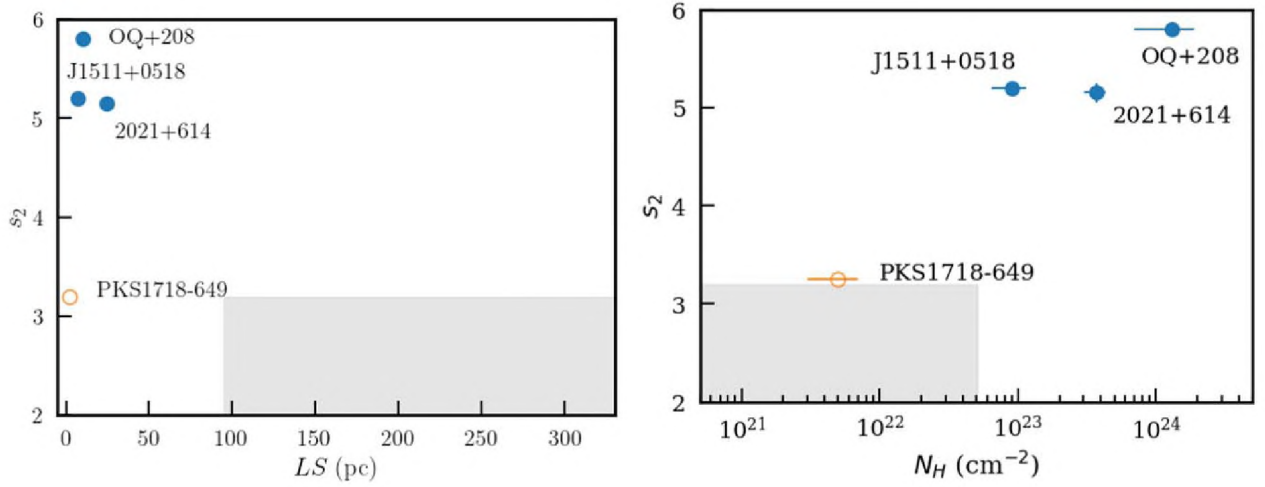


Figure 6. Comparison of the high-energy slope of the injection electron distribution, s_2 , with the radio linear size LS (left panel), and the source-intrinsic-absorbing column density N_H , for the three obscured CSOs analyzed in this paper, J1407+2827, J1511+0518 and J2022+6137, the X-ray unobscured PKS 1718-649 modeled in Sobolewska et al. (2022), and other CSOs with low N_H modeled by O10 and depicted in the figure by the gray-shaded areas.

for black hole masses of $\sim 10^8$ – $10^9 M_\odot$; W20), the ISM density estimates of Cendes et al. (2023) cover a density range, ~ 10 – 10^3 cm^{-3} , similar to that resulting from our work.

X-ray reflection continua, including the fluorescent narrow Iron K α line at 6.4 keV, have been detected in all three CSOs (S19; S23). This raises a question regarding the origin of the X-ray reflection component in these objects. Interestingly, these galaxies fall within the quasar/Seyfert region on the WISE color diagram and their mid-IR colors align with the dominant emission from circumnuclear dust (Kosmaczewski et al. 2020; Nascimento et al. 2022). Additionally, the linear sizes of their radio lobes are of the same order as those of dusty tori in numerous Seyfert Galaxies, as measured by ALMA

(e.g., Combes et al. 2019; García-Burillo et al. 2021). Therefore, we propose that in all three cases, the reflection components may in fact be related to the isotropic IC emission from the lobes reflecting off the tori.

In Appendix B, we provide the energetic considerations, concluding that the lobes' high-energy (X-ray and γ -ray) emission may indeed dominate over the disk corona radiative output in the case of the youngest and the most powerful CSOs. We emphasize in this context that, in both cases, we are dealing with the accretion power (or, to be more precise, a combination of the accretion power and the black hole rotational energy), but, nonetheless, the mechanism for extracting a fraction of this power to produce the high-energy emission is quite distinct. In

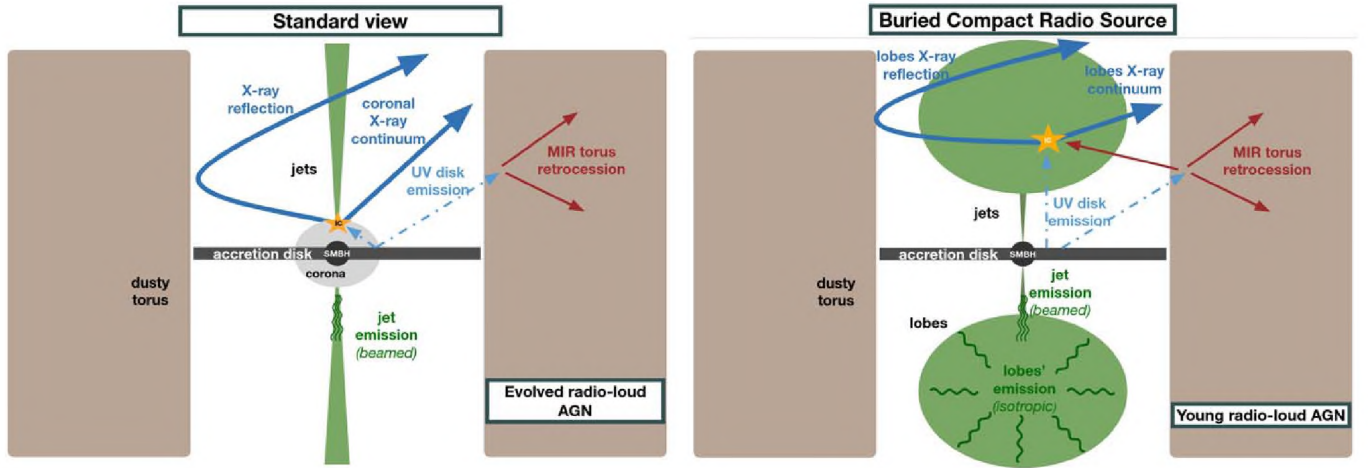


Figure 7. Left panel: schematic view of an evolved radio-loud AGN, depicting the UV disk emission Comptonized within an ultracompact disk hard X-ray emitting region, and subsequently reflected by a dusty torus to form the reflection X-ray component. Right panel: schematic view of a young radio source, where the X-ray continuum emission is produced within compact radio lobes, formed by newly born relativistic jets propagating through the circumnuclear environment.

particular, the energy radiated away in the high-energy domain is either at the expense of a disk magnetic energy dissipated within disk corona to hot electrons (see, e.g., Sridhar et al. 2021), or a bulk kinetic energy of relativistic jets converted at the jet termination shock to the internal energy of the lobes' ultrarelativistic electrons.

In Figure 7, we present a sketch illustrating the proposed scenario. In the left panel of the figure, we present a schematic *standard* view of an evolved radio-loud AGN, depicting, in particular, how the UV disk emission is IC scattered within an ultracompact region filled with hot electrons (e.g., a lamppost corona above the accretion disk, a hot inner flow, or a base of a jet), giving rise to the X-ray continuum emission, which is subsequently reflected from a dusty torus to form the X-ray reflection component. The right panel of the figure illustrates a young radio source, where the X-ray continuum emission originates within the compact radio lobes, inflated by the newly born relativistic jets propagating through the circumnuclear environment. Such a scenario predicts that the X-ray emission of CSOs, whose radio lobes expanded beyond the scale corresponding to the size of a dusty torus should not show strong absorption and reflection features. So far, the existing X-ray observations of the CSO sources support this idea in that they have not identified X-ray absorbed CSOs or CSOs showing a 6.4 keV Iron fluorescence among the sources with radio sizes $\gtrsim 40\text{--}50$ pc (e.g., Sobolewska et al. 2019a).

Interestingly, due to the rather different geometries of the systems shown in the left and right panels of Figure 7, as well as different spatial scales and energies of the radiating electrons involved, variability patterns and polarization properties of the X-ray continuum emission produced in the framework of the standard model and our CSO scenario should be quite distinct. In particular, in the CSO scenario, we expect a rather steady X-ray continuum (power-law) emission, involving possibly only a gradual monotonic decrease on the timescale of years, as opposed to a stochastic red-noise variability expected in the case of an ultracompact X-ray source (Uttley et al. 2014). An in-depth spectral modeling of the reflection component in the framework of our model, along with an analysis of the X-ray polarization properties, keeping in mind the capabilities of current and future X-ray polarimetry missions, will be presented in a forthcoming paper.

7. Summary

We studied the broadband SED of three galaxies with CSO radio classification: J1407+2827, J1511+0518, and J2022+6137. The radio morphologies of these galaxies are characterized by small projected linear sizes (up to 25 pc). Additionally, in X-rays, they are heavily absorbed, $N_{\text{H}} > 10^{23} \text{ cm}^{-2}$, and display fluorescent Fe K α lines in their spectra. The main findings from our modeling of their radio-to- γ -ray emission are summarized as follows:

- (1) The high-energy emission of these three sources can be explained within the framework of an expanding radio lobe model. Namely, the X-ray continuum originates from radio lobes that are still embedded within the dusty tori. These lobes constitute a reservoir of relativistic electrons producing high-energy emission through Comptonization. This emission is then reflected from the torus, adding fluorescent Fe K α lines and the reflection component.
- (2) The proposed scenario may result in different polarization properties than the polarization from a standard compact corona. These properties will be considered in future work.
- (3) Given the energetics of the system, detailed in Appendix B, the high-energy emission of the most compact radio lobes might dominate over the disk corona radiative output.
- (4) The high-energy slope of the broken power-law electrons' energy distribution at injection has to be notably soft ($s_2 \simeq 4.7\text{--}5.8$) to account for the spectral shape of the radio emission in all three sources. Relative to other CSOs, this index appears to correlate with the X-ray absorbing column density (the higher N_{H} , the higher s_2) and to anticorrelate with source size (the larger the object, the smaller s_2). The softness of the high-energy index of the electron energy spectrum at the injection suggests that the efficiency of particle acceleration decreases with the density of an ambient medium. This anti-correlation between acceleration efficiency and n_0 might be responsible for the γ -ray quietness of these sources. Additionally, it might be responsible for the γ -ray quietness of these sources. In general, the obtained ISM density values

are rather high, ranging from ~ 10 to even 10^4 cm^{-3} . Interestingly, the ISM densities obtained through our modeling agree with estimates obtained from the multi-epoch radio follow-ups of TDE events in several galaxies, at distances corresponding to the radio lobes size in our sources.

- (5) The broadband spectral properties of J2022+6137 differ from those of J1407+2827 and J1511+0518 in several aspects. The ratio of the intrinsic (de-absorbed) X-ray power-law flux to the upper limits on the γ -ray flux is significantly higher in J2022+6137 than in the case of the other two CSOs. Consequently, for some models, the data of J2021+624 require an additional X-ray emission component (as in the case of PKS 1718-649; Sobolewska et al. 2022). This additional component can be associated with, e.g., a compact corona, a base of a jet, or a hot inner flow. The emerging total jet power of J1407+2827 and J1511+0518 is $\lesssim 10^{44} \text{ erg s}^{-1}$. However, in J2022+6137, the jet power of the lowest-density models reaches and even exceeds this threshold, which suggests that jet expansion beyond the host galaxy is more likely in the case of J2022+6137 than in the case of the other two considered sources (Mukherjee et al. 2016, 2017). Finally, our model underpredicts the radio emission of J2022+6137 below the spectral break, signaling either the presence of an additional, extended radio component or indicating that it cannot be described within the framework of the free-free absorption on engulfed clouds.

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This paper employs a list of Chandra data sets obtained by the Chandra X-ray Observatory.¹⁴

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Software: CIAO v.4.14 (Fruscione et al. 2006), Sherpa (Freeman et al. 2001).

Appendix A Observational Data

Here, we present the observational data used in our analysis. The Fermi-LAT upper limits on the CSOs fluxes within seven narrow energy bins are listed in Table A1. The SMA data are presented in Table A2. For details on the data processing, see Sections 3.1 and 3.2. A list of relevant references for the multiwavelength data acquired from NED is provided in Table A3.

Table A1
Fermi-LAT Upper Limits

Object	Energy Band (GeV)	Upper Limit ($10^{-13} \text{ erg cm}^{-2} \text{ s}^{-1}$)
J1407+2827	0.100–0.316	5.53
	0.316–1.000	1.51
	1.000–3.162	1.97
	3.162–10.00	0.52
	10.00–31.62	0.74
	31.62–100.0	2.14
	100.0–1,000	3.88
J1511+0518	0.100–0.316	2.26
	0.316–1.000	2.07
	1.000–3.162	1.69
	3.162–10.00	0.60
	10.00–31.62	0.86
	31.62–100.0	3.33
	100.0–1,000	4.92
J2022+6137	0.100–0.316	1.25
	0.316–1.000	0.85
	1.000–3.162	0.49
	3.162–10.00	0.39
	10.00–31.62	0.50
	31.62–100.0	4.89
	100.0–1,000	3.11

¹⁴ doi:10.25574/cdc.213

Table A2
Details of the SMA Observations

Source	Redshift	R.A. (J2000)	Decl. (J2000)	Date	ν (GHz)	Flux (mJy)
J1407+2827	0.076	14 ^h 07 ^m 0 ^s .4	+28°27′14″.7	2018 Aug 08	227.225	6.96 ± 3.20
J1511+0518	0.084	15 ^h 11 ^m 41 ^s .3	+05°18′09″.3	2018 Jul 13	225.561	10.22 ± 2.55
				2018 Aug 08	225.558	15.46 ± 2.98
				2019 May 02	225.558	5.01 ± 1.35
J2022+6137	0.227	20 ^h 22 ^m 06 ^s .7	+61°36′58″.8	2018 Jul 13	225.561	90.11 ± 2.89

Table A3
References to the Archival Radio, IR, and Optical/UV Data

Energy Band	Source	References
Radio	J1407+2827	Adgie et al. (1972), Pauliny-Toth et al. (1972), Kellermann & Pauliny-Toth (1973), Owen & Mufson (1977), Owen et al. (1978), Witzel et al. (1979), Kojanian et al. (1980), Owen et al. (1980), Geldzahler & Witzel (1981), Kuehr et al. (1981), Condon et al. (1983), Chini et al. (1989), Langston et al. (1990), Becker et al. (1991), White & Becker (1992), Wiren et al. (1992), Bica et al. (1995), Douglas et al. (1996), Stanghellini et al. (1998), Dallacasa et al. (2000), Antón et al. (2004), Horiuchi et al. (2004), Kellermann et al. (2004), Kovalev et al. (2005), Stanghellini et al. (2005), Tinti et al. (2005), Rickett et al. (2006), Xiang et al. (2006), Healey et al. (2007), Helmboldt et al. (2007), Liu et al. (2007), Orienti et al. (2007), Vollmer et al. (2008), Richards et al. (2011), Mingaliev et al. (2014), Richards et al. (2014), Pasetto et al. (2016), Lee et al. (2017)
	J1511+0518	Becker et al. (1991), Griffith et al. (1995), Dallacasa et al. (2000), Tinti et al. (2005), Orienti et al. (2006), Xiang et al. (2006), Healey et al. (2007), Orienti et al. (2007), Richards et al. (2011), Mingaliev et al. (2014), Richards et al. (2014), Lee et al. (2017), Bennett et al. 1986, Gregory & Condon (1991), Condon et al. (1998), Ching et al. (2017)
	J2022+6137	Dallacasa et al. (2000), Bennett et al. (2003), Cara & Lister (2008), Arshakian et al. (2010), Chen & Wright (2009), Condon et al. (1998), Douglas et al. (1996), Giommi et al. (2009), Gold et al. (2011), Gregory & Condon (1991), Healey et al. (2007), Horiuchi et al. (2004), Kellermann et al. (2004), Kovalev et al. (2005), Kuehr et al. (1981), Lee et al. (2017), Lister & Homan (2005), Lister et al. (2011), Lister et al. (2015), Lowe et al. (2007), Massardi et al. (2009), Owen et al. (1980), Planck Collaboration et al. (2011), Stanghellini et al. (1998), Patnaik et al. (1992), Pauliny-Toth et al. (1972), Richards et al. (2011), Richards et al. (2014), Vollmer et al. (2008), White & Becker (1992), Wiren et al. (1992), Witzel et al. (1978), Witzel et al. (1979), Wright et al. (2009), Zavala & Taylor (2004)
IR	J1407+2827	Cutri et al. (2013), Hanish et al. (2015), Terashima et al. (2015)
	J1511+0518	Cutri et al. (2013), Hanish et al. (2015)
	J2022+6137	Hanish et al. (2015), Chen & Wright (2009), Agudo et al. (2014), Lee et al. (2008), Wright et al. (2009), Gold et al. (2011), Bennett et al. (2003)
Optical/UV	J1407+2827	Petrosian et al. (2007), Petrosian et al. (2008), de Vaucouleurs & Longo (1988), Healey et al. (2008), Skrutskie et al. (2006), Rudy et al. (1982), Heckman et al. (1983), Fanti et al. (2011), Seibert et al. (2012)
	J1511+0518	Adelman-McCarthy et al. (2008), Makarov et al. (2017), Skrutskie et al. (2006), Fanti et al. (2011), Seibert et al. (2012)
	J2022+6137	Healey et al. (2008), Skrutskie et al. (2006)

Appendix B

Basic Energetic Considerations

The accretion disk UV luminosity can be written as

$$L_{UV} \simeq \eta_d \dot{M}_{acc} c^2, \quad (B1)$$

where η_d is the disk radiative efficiency, and $\dot{M}_{acc}$ is the mass accretion rate. For standard geometrically thin/optically thick disks, one expects $\eta_d \sim 0.1$ (depending, however, on the black hole spin; see, e.g., Abramowicz & Fragile 2013, and references therein). Some fraction of this emission may be reprocessed within the disk corona, and reemitted in the X-ray range, so the resulting corona luminosity is $L_X \simeq \eta_c L_{UV}$. The efficiency factor here, η_c , may obviously change from source to source, and may, in particular, scale with the disk luminosity. A relatively tight X-ray/UV luminosity correlation established for quasars (e.g., Lusso & Risaliti 2016) nevertheless indicates average values of $\eta_c \sim 0.1$ for lower-luminosity sources ($L_{UV} \lesssim 10^{45} \text{ erg s}^{-1}$), decreasing down to ~ 0.01 for high-luminosity ones ($L_{UV} \gtrsim 10^{47} \text{ erg s}^{-1}$). The mid-IR luminosity of a dusty torus is also expected to constitute some fraction of a time-averaged disk emission, $L_{IR} \simeq \eta_h L_{UV}$, where $\eta_h \gtrsim 0.5$ as

implied by the analysis of the broadband SEDs of quasars (see, e.g., Rąłowski et al. 2024).

The energy density of the nuclear emission components, including both the torus IR and the disk UV emission, at distances larger than the dust sublimation radius, i.e., effectively on the scales of the order of the CSO linear size $LS > 1 \text{ pc}$ (defined in this paper as half of the hotspot-hotspot distance) is

$$U_0 = \frac{(1 + \eta_h) L_{UV}}{4\pi LS^2 c}. \quad (B2)$$

The total IC luminosity of the lobes is therefore

$$L_{IC} = \frac{4c\sigma_T}{3m_e c^2} f_\gamma V U_0 U_e, \quad (B3)$$

where V is the total volume of the lobes,

$$f_\gamma \equiv \langle \gamma^2 \rangle / \langle \gamma \rangle = \int d\gamma \gamma^2 N_d(\gamma) / \int d\gamma \gamma N_d(\gamma), \quad (B4)$$

$N_e(\gamma)$ is the evolved electron energy distribution within the lobes, and U_e is the corresponding electron energy density,

$$U_e \equiv \int d\gamma \gamma m_e c^2 N_d(\gamma). \quad (B5)$$

For various electron injection spectra and other CSO parameters explored in S08, f_γ turned out to typically be of the order of the electron cooling energy, $\gamma_{\text{cr}} \sim 10^2 - 10^3$.

In spite of the electron energy distribution $N_e(\gamma)$ evolving with time in the expanding lobes, due to the adiabatic and radiative cooling, S08 argued that at every moment of the CSO lifetime, U_e constitutes approximately the same fixed fraction of the lobes' total pressure p , namely, $U_e \simeq \eta_e p$. In the case of radiating electrons dominating the lobes' pressure, one has, for example, $\eta_e \simeq 3$. And since the total energy of the plasma filling the lobes is deposited by a pair of jets over the source lifetime, one may further write

$$3pV = 2L_j \tau_j, \quad (\text{B6})$$

where L_j is the jet kinetic power, and $\tau_j = \text{LS}/v_h$ is the source lifetime for the observed (and constant, by assumption) hotspots' advance velocity, v_h (defined here as half of the hotspot-hotspot separation velocity). Note that the jet total power and the mass accretion rate can be related by introducing the jet production efficiency parameter η_j , namely,

$$2L_j = \eta_j \dot{M}_{\text{acc}} c^2. \quad (\text{B7})$$

In principle, in the most favorable conditions provided by the magnetically arrested disks around maximally spinning black holes, η_j may even exceed unity, as the jet power is in this case extracted at the expense of the accretion power and the black hole rotational energy (Tchekhovskoy et al. 2011). However, for the sample of known CSOs with measured parameters of the central engine, W20 observed in general values of the jet production efficiency parameter much lower than unity, with a significant spread around from $\eta_j \gtrsim 0.01$ to < 0.1 .

Taking into account the above considerations, the ratio of the lobes' total IC luminosity and the disk corona luminosity, reads as

$$\begin{aligned} \frac{L_{\text{IC}}}{L_X} &\simeq \frac{4 G m_p}{9 m_e c} \frac{(1 + \eta_e) \eta_e \eta_j}{\eta_e \eta_d} \frac{f_\gamma \lambda_{\text{Edd}} M_{\text{BH}}}{\text{LS } v_h} \\ &\sim \frac{(1 + \eta_e) \eta_e \eta_j}{\eta_e \eta_d} \left(\frac{f_\gamma}{10^3} \right) \left(\frac{M_{\text{BH}}}{10^9 M_\odot} \right) \\ &\times \left(\frac{\lambda_{\text{Edd}}}{0.01} \right) \left(\frac{\text{LS}}{10 \text{ pc}} \right)^{-1} \left(\frac{v_h}{0.03 c} \right)^{-1}, \end{aligned} \quad (\text{B8})$$

where we introduced the Eddington ratio

$$\lambda_{\text{Edd}} \equiv \frac{L_{\text{UV}}}{L_{\text{Edd}}}, \quad (\text{B9})$$

for the Eddington luminosity $L_{\text{Edd}} = 4\pi G M_{\text{BH}} m_p c / \sigma_T$, corresponding to the black hole mass M_{BH} . The luminosity ratio given in Equation (B8) may be larger or smaller than unity, depending on a particular combination of the source parameters. In general, however, one may expect that the lobes' IC emission dominates over the disk corona radiative output in the case of the youngest and the most powerful CSOs, i.e., sources characterized by high accretion rates, $\lambda_{\text{Edd}} > 0.01$, and compact lobes with linear sizes of $\text{LS} \lesssim 10 \text{ pc}$, preferentially located in a dense environment so that the jets' advance velocities are low, $v_h < 0.1c$.

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Chapter 5

Magnetization of Relativistic Current-carrying Jets with Radial Velocity Shear



Magnetization of Relativistic Current-carrying Jets with Radial Velocity Shear

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Abstract

Astrophysical jets, launched from the immediate vicinity of accreting black holes, carry away large amounts of power in a form of bulk kinetic energy of jet particles and electromagnetic flux. Here we consider a simple analytical model for relativistic jets at larger distances from their launching sites, assuming a cylindrical axisymmetric geometry with a radial velocity shear, and purely toroidal magnetic field. We argue that as long as the jet plasma is in magnetohydrostatic equilibrium, such outflows tend to be particle dominated, i.e., the ratio of the electromagnetic to particle energy flux, integrated over the jet cross-sectional area, is typically below unity, $\sigma < 1$. At the same time, for particular magnetic and radial velocity profiles, magnetic pressure may still dominate over particle pressure for certain ranges of the jet radius, i.e., the local jet plasma parameter $\beta_{\text{pl}} < 1$, and this may be relevant in the context of particle acceleration and production of high-energy emission in such systems. The jet magnetization parameter can be elevated up to the modest values of $\sigma \lesssim \mathcal{O}(10)$ only in the case of extreme gradients or discontinuities in the gaseous pressure, and a significantly suppressed velocity shear. Such configurations, which consist of a narrow, unmagnetized jet spine surrounded by an extended, force-free layer, may require an additional poloidal field component to stabilize them against current-driven oscillations, but even this will not substantially elevate their σ parameter.

Unified Astronomy Thesaurus concepts: Jets (870); Plasma jets (1263); Magnetohydrodynamics (1964)

1. Introduction

Relativistic jets found in various types of astrophysical sources of high-energy radiation, such as active galactic nuclei (AGNs), microquasars, and gamma-ray bursts, are believed to be formed via the efficient extraction of energy and angular momentum, in the form of Poynting flux, from the rotating black hole/accretion disk system (Blandford & Znajek 1977; for an updated compendium see Meier 2012, also Komissarov & Porth 2021 for a summary of recent developments in numerical simulations of jet launching). At the initial stages of their evolution, such structures are magnetically dominated, rapidly expanding, and only mildly relativistic. Thereafter, collimation and acceleration start to proceed in accord, gradually converting the outflow to a fully formed plasma jet, at the expense of the magnetic energy. However, in the framework of the ideal magnetohydrodynamic (MHD) description, and in the relativistic regime, such a conversion cannot be efficient, in the sense that the collimation and acceleration of the flow due to the magnetic tension and magnetic pressure gradient, respectively, are limited by the increasing electric force, which counterbalances the magnetic force. Efficient collimation and acceleration may, however, be reinforced by the external pressure, for example, related to the star's interior in the case of gamma-ray bursts, or accretion disk winds in the case of AGN (e.g., Lyubarsky 2009, and references therein).

Considering a power-law external pressure profile, Lyubarsky (2010) showed in particular that for various initial magnetic field configurations or external pressure profiles jets could possibly cease to be Poynting-flux dominated only at *logarithmically large distances*. It is worth pointing out that the author defines a *matter-dominated* jet through the condition $\sigma < 0.1$, where σ is the ratio of the jet magnetic energy flux to the particles' kinetic energy flux; the reason for this is that in the intermediate regime, $0.1 < \sigma < 1$, formation of strong shocks within the outflow—and hence energy dissipation as well as bulk deceleration of the flow due to interactions with the ambient medium—proceed rather differently than in the purely hydrodynamical case (e.g., Komissarov 1999; Kirk et al. 2000).

A contradictory conclusion follows from modeling of the observational data for relativistic jets in AGN, which often point toward very weak magnetization of the outflows, even at relatively close distances from the launching site. For example, various approaches to the spectral fitting of broadband blazar emission typically give the estimate $\sigma \ll 1$ for distances $\lesssim 10^4 r_g$, where $r_g \simeq 10^{14} M_9 \text{ cm}$ is the gravitational radius corresponding to the black hole mass $M_9 \equiv M/10^9 M_\odot$ (e.g., Sikora et al. 2009; Ghisellini et al. 2010; Rueda-Becerril et al. 2014; Saito et al. 2015; but see also in this context Sobacchi & Lyubarsky 2019). This would be consistent with the maximum efficiency of Poynting-to-matter energy flux conversion achieved relatively close to the jet base, followed by a steady and basically dissipation-free propagation of a relativistic matter-dominated outflow. Beyond the ideal MHD approximation, such a maximum efficiency could possibly be achieved with the help of dissipation processes, in particular,



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magnetic reconnection, or the nonlinear development of various MHD instabilities, rapidly converting the jet magnetic energy into kinetic energies of the jet particles (e.g., Sikora et al. 2005; Giannios & Spruit 2006; Chatterjee et al. 2019).

Keeping in mind the results and findings summarized above, here we consider the simplest analytical model for a relativistic, current-carrying jet at a *large distance* from the launching site, where the outflow can be considered as fully collimated and accelerated to terminal bulk velocities. In particular, we assume a cylindrical axisymmetric geometry for a nonrotating jet, with a radial velocity shear; moreover, we assume that the jet magnetic field has only a toroidal component, the jet is in *magnetohydrostatic equilibrium*, and finally that the jet particles obey an ultrarelativistic equation of state.

In the framework of this approximation, introduced in more detail in Section 2, we argue that the outflows tend to be matter dominated, in the sense that the ratio of the jet Poynting and particle energy fluxes, integrated over the jet cross-sectional area, i.e., the jet magnetization σ parameter, is less than unity, as long as the radial pressure profiles are smooth with no extreme gradients or discontinuities. In Section 3, we provide simple analytical proof for this statement, valid however for only a certain class of magnetic field profiles, namely, those for which the rest-frame magnetic pressure attains its global maximum at the jet boundary. In Section 4, we explore further numerical solutions corresponding to various parameterizations of the generalized magnetic and velocity profiles, as well as various boundary conditions. There we observe that, even in the case of matter-dominated outflows with $\sigma < 1$, it is possible for the magnetic pressure to exceed particle pressure at certain ranges of the jet radii, or in other words for the local plasma beta parameter to alternate between $\beta_{\text{pl}} > 1$ and $\beta_{\text{pl}} < 1$ depending on the distance from the jet axis. We also explore cases with tangential discontinuities in the jet gaseous pressure, arguing that, for such, one may formally obtain magnetic-dominated outflows with $\sigma \lesssim \mathcal{O}(10)$, however, at the price of huge pressure jumps by orders of magnitude. In Section 5, we discuss our findings in the general context of particle acceleration in astrophysical jets, commenting also on the stability of the analyzed structures against current-driven oscillations, and on a role of an additional component of a poloidal magnetic field, in particular, when confined to the narrow spine of the jet.

2. Jet Model

Let us consider a nonrotating, cylindrical axisymmetric jet in magnetohydrostatic equilibrium, for which the comoving magnetic field $\mathbf{B}'$ is decomposed into the poloidal and toroidal components,

$$\mathbf{B}'_P = B'_r \hat{r} + B'_z \hat{z} \quad \text{and} \quad \mathbf{B}'_T = B'_\phi \hat{\phi}. \quad (1)$$

In this paper, we analyze a purely toroidal configuration, $B'_r = B'_z = 0$, due to the fact that at large distances from the jet base, any poloidal magnetic field is typically expected to be negligible (see Begelman et al. 1984). The current $\mathbf{I}$ associated with such a jet is therefore parallel or antiparallel to the jet axis, depending on the helicity of the jet magnetic field. Note that this would not be valid for a conical jet, for which there could also be a radial current component. However, at large distances from the jet base, effectively beyond the host galaxy, the background pressure should not be that of the interstellar

medium, but instead of the hot, over-pressured jet cocoon, and as such should be independent of z , justifying our model assumption regarding a cylindrical jet geometry (see in this context Begelman & Cioffi 1989). The only nonzero component of the current density, $\mathbf{J}' = J' \hat{z}$, is then

$$J' = \frac{c}{4\pi r'} \partial_{r'}(r' B'_\phi), \quad (2)$$

so that the magnetohydrostatic equilibrium condition $c \partial_{r'} P = -J' B'_\phi$, where P denotes the comoving particle pressure, gives the radial profile

$$\partial_r P = -\frac{1}{8\pi r^2} \partial_r \left(\frac{r^2 B_\phi^2}{\Gamma^2} \right), \quad (3)$$

where $B_\phi = B'_\phi \Gamma$ and $\Gamma = (1 - \beta^2)^{-1/2}$ is the bulk Lorentz factor of the portion of the fluid with the velocity β (see, e.g., Appl & Camenzind 1992). The jet current, $I = \frac{1}{2} c r B_\phi(R_j)$, where R_j is the jet radius, has to be compensated by a return current I_{ret} at larger scales where the outflow terminates.

Let us denote the proper enthalpy and the proper internal energy density of the jet particles by w and u , respectively. In the case of ultrarelativistic jet particles, i.e., for $P = u/3$, one has $w = P + u = 4P$, the case we analyze exclusively hereafter. With such, the energy flux associated with the jet particles is

$$L_p = 2\pi c \int dr r \Gamma^2 \beta w = 8\pi c \int dr r \beta \Gamma^2 P, \quad (4)$$

while the jet Poynting flux is

$$L_B = 2\pi c \int dr r \frac{E_r B_\phi}{4\pi} = \frac{1}{2} c \int dr r \beta B_\phi^2, \quad (5)$$

where $E_r = \beta B_\phi$ is the radial component of the jet electric field. The jet magnetization parameter σ is then

$$\sigma \equiv \frac{L_B}{L_p} = \frac{1}{16\pi} \frac{\int dr r \beta B_\phi^2}{\int dr r \beta \Gamma^2 P}. \quad (6)$$

It is convenient to introduce the dimensionless parameter $x \equiv r/R_j$, and the normalized profiles

$$b(x) \equiv \frac{B_\phi(x)}{B_\phi(1)}: \quad b(0) = 0, \quad b(1) \equiv 1, \quad \text{and} \quad (7)$$

$$p(x) \equiv \frac{P(x)}{P(1)}: \quad p(0) > 0, \quad p(1) \equiv 1. \quad (8)$$

For now, we do not make any assumption regarding the particular monotonicity of $b(x)$ or $p(x)$ profiles, but only that the magnetic field has a particular helicity all along the outflow, and that it does not vanish at the jet boundary, $B_\phi(1) > 0$; note however that for a purely toroidal configuration, magnetic field always drops to zero at the jet axis, $B_\phi(0) = 0$. As for the velocity profile, we consider only the cases with the jet bulk Lorentz factor $\Gamma(x)$ decreasing monotonically along the jet radius starting from $\Gamma(0) \equiv \Gamma_0 > 1$ down to $\Gamma(1) = 1$, excluding anomalous shear layers (see in this context Alloy & Rezzolla 2006; Mizuno et al. 2008).

Next we introduce the parameter

$$q \equiv \frac{B_\phi^2(1)/8\pi}{P(1)} \equiv \frac{P_B(1)}{P(1)} \equiv \beta_{\text{pl}}^{-1}(1), \quad (9)$$

which is the ratio of the magnetic and gas pressures at the jet boundary, i.e., the inverse of the plasma beta parameter at $x = 1$. To simplify the notation, we further introduce

$$f(x) \equiv \frac{b^2(x)}{\Gamma^2(x)}: \quad f(0) = 0, \quad f(1) = 1, \quad (10)$$

which is basically the normalized rest-frame magnetic pressure, $f(x) \propto P_B(x) \equiv B_\phi^2/8\pi$.

With the above definitions, for a given set of profiles $b(x)$ and $\Gamma(x)$, or equivalently $f(x)$ and $\Gamma(x)$, and with a given boundary condition q , the magnetohydrostatic equilibrium condition, given in Equation (3), determines the particle pressure profile

$$p(x) = 1 + q - q f(x) + 2q \int_x^1 ds \frac{f(s)}{s}. \quad (11)$$

Importantly, the requirement $\forall_{x \in [0,1]}: p(x) > 0$, translates to the constraint

$$1 + q^{-1} > f(x) - 2 \int_x^1 ds \frac{f(s)}{s}. \quad (12)$$

3. Jet σ Parameter

It was noted by Lisanti & Blandford (2007) that for the jet model setup as considered here, numerical solutions to the jet magnetization parameter always return $\sigma < 1$. Below we draft simple analytical proof for this statement, valid however for only a certain class of magnetic profiles. In particular, in this section we consider only those magnetic profiles, for which the comoving magnetic field pressure is continuous and attains its maximum at the jet boundary, or in other words for which $f(x) < 1$ for every $x < 1$, even though $f(x)$ is not necessarily monotonic within the entire range $x \in [0, 1]$.

First, let us introduce the function

$$h(x) \equiv x \beta(x) \Gamma^2(x), \quad (13)$$

such that $h(0) = h(1) = 0$, which has exactly one maximum at a given $x_0 \in (0,1)$ for the assumed monotonically decreasing $\Gamma(x)$. With such, the jet σ parameter can be rewritten as

$$\sigma = \frac{q}{2} \frac{\int_0^1 dx h(x) f(x)}{\int_0^1 dx h(x) p(x)}. \quad (14)$$

We also define

$$H(x) \equiv \frac{\int_0^x dy h(y)}{\int_0^1 dz h(z)}, \quad (15)$$

which, by definition, is monotonically increasing from $H(0) = 0$ to $H(1) = 1$. Note that

$$\frac{x H'(x)}{H(x)} \equiv \frac{d \ln H(x)}{d \ln x} = \frac{x h(x)}{\int_0^x dy h(y)}. \quad (16)$$

Moreover, since $\beta(x) \Gamma^2(x)$ is a continuous, monotonically decreasing function, from the mean value theorem (MVT) we know that for a given fixed x , there is always $\xi \in [0, x]$ such that

$$\int_0^x dy y \beta(y) \Gamma^2(y) = \beta(\xi) \Gamma^2(\xi) \int_0^x dy y = \frac{1}{2} x^2 \beta(\xi) \Gamma^2(\xi),$$

and therefore

$$\frac{x H'(x)}{H(x)} = \frac{x^2 \beta(x) \Gamma^2(x)}{\frac{1}{2} x^2 \beta(\xi) \Gamma^2(\xi)} \leq 2. \quad (17)$$

In fact, $x H'(x)/H(x)$ is always monotonically decreasing from 2 at $x = 0$ down to 0 at $x = 1$.

Using the form of the σ parameter as given in Equation (14), along with the solution for the pressure profile, Equation (11), and the condition for a positive gas pressure, Equation (12), we obtain

$$\frac{1}{2\sigma} > \frac{f(x) - 2 \int_x^1 ds \frac{f(s)}{s} - \langle f \rangle + \left\langle 2 \int_y^1 ds \frac{f(s)}{s} \right\rangle}{\langle f \rangle}, \quad (18)$$

where the averaging is over the $h(x)$ distribution, namely,

$$\langle X \rangle \equiv \int_0^1 dH X \equiv \frac{\int_0^1 dy h(y) X}{\int_0^1 dy h(y)}, \quad (19)$$

with $\int dH \equiv \int dx H'$. Hence, if the right-hand side of Equation (18) is larger than 1/2 for every x , then $\sigma < 1$.

In the case of $f(x)$ attaining its maximum at the jet boundary, as considered in this section, the condition, Equation (18), at $x = 1$ reads as

$$\frac{1}{2\sigma} > \frac{1 - \langle f \rangle + \left\langle 2 \int_y^1 ds \frac{f(s)}{s} \right\rangle}{\langle f \rangle}. \quad (20)$$

Meanwhile, from the MVT we know that there is $\eta \in (0, 1)$ such that $\langle f \rangle = f(\eta) < 1$. Hence,

$$\begin{aligned} \left\langle 2 \int_y^1 ds \frac{f(s)}{s} \right\rangle &= 2 \int_0^1 dy H' \int_y^1 ds \frac{f(s)}{s} \\ &= 2 \int_0^1 dy H(y) \frac{f(y)}{y} \\ &= 2 \frac{H}{yH'} \bigg|_{y=\zeta} \int_0^1 dH f = 2 \frac{H}{yH'} \bigg|_{y=\zeta} \langle f \rangle \geq \langle f \rangle, \end{aligned} \quad (21)$$

where we used again the MVT with a certain $\zeta \in (0,1)$, and the fact that $H/(yH') \leq 1/2$ for every $y \in (0,1)$ (Equation (17)). All in all, we therefore have

$$\frac{1}{2\sigma} > \frac{1}{f(\eta)} > 1, \quad (22)$$

resulting in $\sigma < \frac{1}{2}$.

4. Jet Radial Profiles

In the framework of our simple model for current-carrying jets in magnetohydrostatic equilibrium, outlined in Section 2, the jet magnetization parameter σ is always less than unity, as long as the comoving toroidal magnetic field pressure is continuous and attains its maximum at the jet boundary; simple proof of this statement has been presented in Section 3. More complex cases, with pressure discontinuities and sharp maxima of the magnetic pressure located well within the jet body, can be investigated numerically, by adopting various parameterizations of the $b(x)$ and $\Gamma(x)$ profiles.

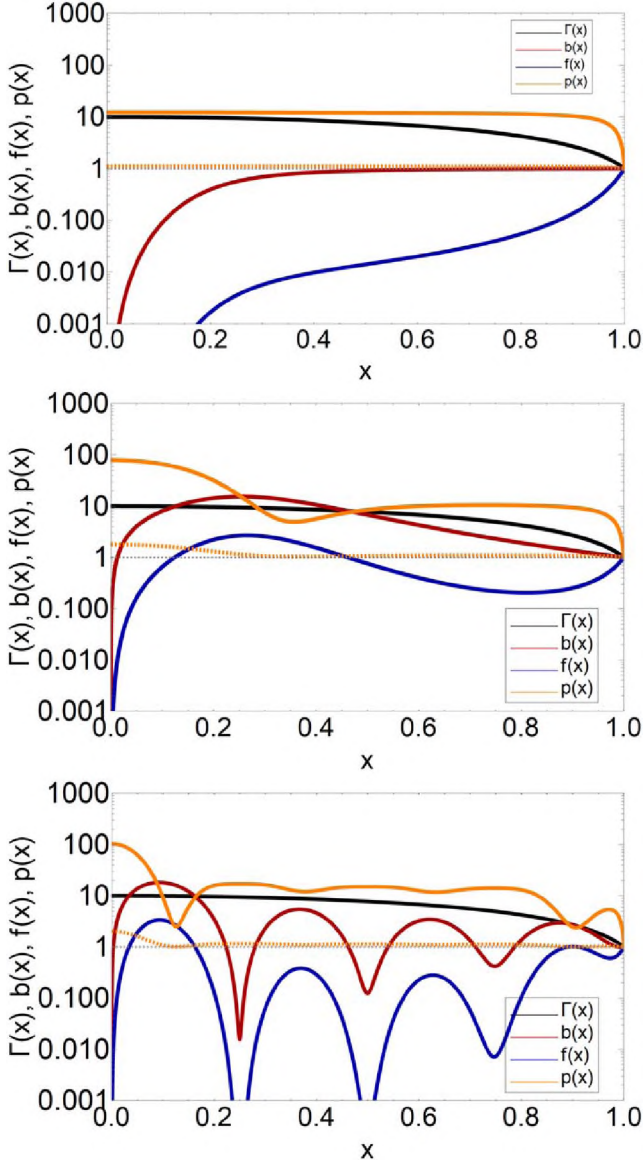


Figure 1. Three exemplary sets of $b(x)$ and $\Gamma(x)$ profiles, along with the corresponding $f(x)$ profiles, and the resulting pressure profiles $p(x)$ calculated for the boundary conditions $q = 0.1$ (dashed curves) and $q = 10$ (solid curves).

4.1. Continuous Profiles

As noted in Section 2, the anticipated magnetohydrostatic equilibrium condition determines the jet particle pressure profile, $p(x)$, for a given magnetic pressure and jet bulk Lorentz factor profiles $b(x)$ and $\Gamma(x)$, as well as a given boundary condition $q \equiv \beta_{\text{pl}}^{-1}(1)$. Let us therefore consider the three illustrative cases of certain $b(x)$ and $\Gamma(x)$ parameterizations, corresponding to the three different behaviors of the function $f(x)$, namely, (a) the case with monotonically increasing $f(x)$, (b) the case with one pronounced global maximum of $f(x)$ at $x < 1$, and (c) the case with multiple local maxima of $f(x)$ throughout the entire range of x .

In the following numerical analysis, for clarity we follow one particularly simple parameterization of the jet bulk Lorentz, namely,

$$\Gamma(x) = 1 + (\Gamma_0 - 1) \times (1 - x^k), \quad (23)$$

with $\Gamma_0 = 10$ and $k = 2$. For the magnetic profile $b(x)$, on the other hand, we consider various analytical prescriptions, selected to ensure positive particle pressure at all jet radii. In Figure 1 we present the three exemplary sets of $b(x)$ and $\Gamma(x)$ profiles, along with the corresponding $f(x)$ profiles and the resulting pressure profiles $p(x)$ calculated for $q = 0.1$ and $q = 10$.

As shown, the particle pressure is maximized at the jet axis, and for monotonically increasing $f(x)$, it decreases monotonically toward the jet boundary. In the case of more complex magnetic pressure profiles, local maxima in the $b(x)$ profile correspond to local minima in $p(x)$. In general, sharp gradients in the comoving magnetic pressure with $f(x)$ locally exceeding unity, often generate exponential drops in particle pressure, leading to unphysical solutions with $p(x) < 0$ for some ranges of the jet radius.

Figure 1 also indicates how the boundary condition q impacts the resulting jet particle pressure profiles, and hence the overall jet magnetization. Figure 2 illustrates the latter effect explicitly, by providing the exact σ values corresponding to different values of $q \in (0.01, 100)$. As shown, starting from small q , i.e., negligible magnetic pressure at the jet boundary (with respect to the particle pressure), the jet magnetization σ increases monotonically with increasing q , but only until $q \gtrsim 1$; from that point, the σ parameter saturates and remains effectively constant (always below the unity), regardless of the ever-increasing $q > 10$.

Magnetohydrostatic equilibrium, assumed here for the jet, implies that any spatial changes in the magnetic pressure across the jet must be counter-balanced by the changes in the particle pressure and the magnetic tension. As a result, the rest-frame plasma beta parameter

$$\beta_{\text{pl}}^{-1}(x) \equiv \frac{P_B(x)}{P(x)} = \frac{q f(x)}{p(x)} \quad (24)$$

may change quite dramatically along the jet radius. Similarly, the *local* jet magnetization σ parameter calculated for different *layers* of a sheared outflow with thickness Δ , namely,

$$\overline{\sigma(x \pm \Delta)} \equiv \frac{\langle f(x) \rangle_{\Delta}}{\langle p(x) \rangle_{\Delta}} \equiv \frac{q}{2} \frac{\int_{x-\Delta}^{x+\Delta} dy h(y) f(y)}{\int_{x-\Delta}^{x+\Delta} dy h(y) p(y)}. \quad (25)$$

may also vary across the jet. Those changes, corresponding to the three sets of the $b(x)$ and $\Gamma(x)$ profiles considered in this subsection, are shown in Figure 3, for the two selected boundary condition values $q = 0.1$ and 10 , with midpoints of averaging $x = \{0.1, 0.3, 0.5, 0.7, 0.9\}$ and $\Delta = 0.1$.

In the case of a monotonic $f(x)$, both $\beta_{\text{pl}}^{-1}(x)$ and $\sigma(x)$ increase monotonically from the jet spine toward the jet boundary. Models for which the comoving magnetic pressure attains a local maximum with $f(x) \gtrsim 1$ are, on the other hand, characterized by $\beta_{\text{pl}}^{-1}(x)$ exceeding unity at the approximately corresponding jet radii; the local $\sigma(x)$ parameter may then also approach unity for certain ranges of x , depending on the considered layer thickness Δ . In other words, for such complex models with alternating currents, we observe the presence of extended domains within the outflow in which the comoving jet magnetic pressure dominates over the gaseous pressure, while the jet Poynting energy flux is comparable to the kinetic energy flux of the jet particles, even though the global σ value (i.e., the value obtained after integrating over the entire cross-sectional area of the jet) is significantly below unity.

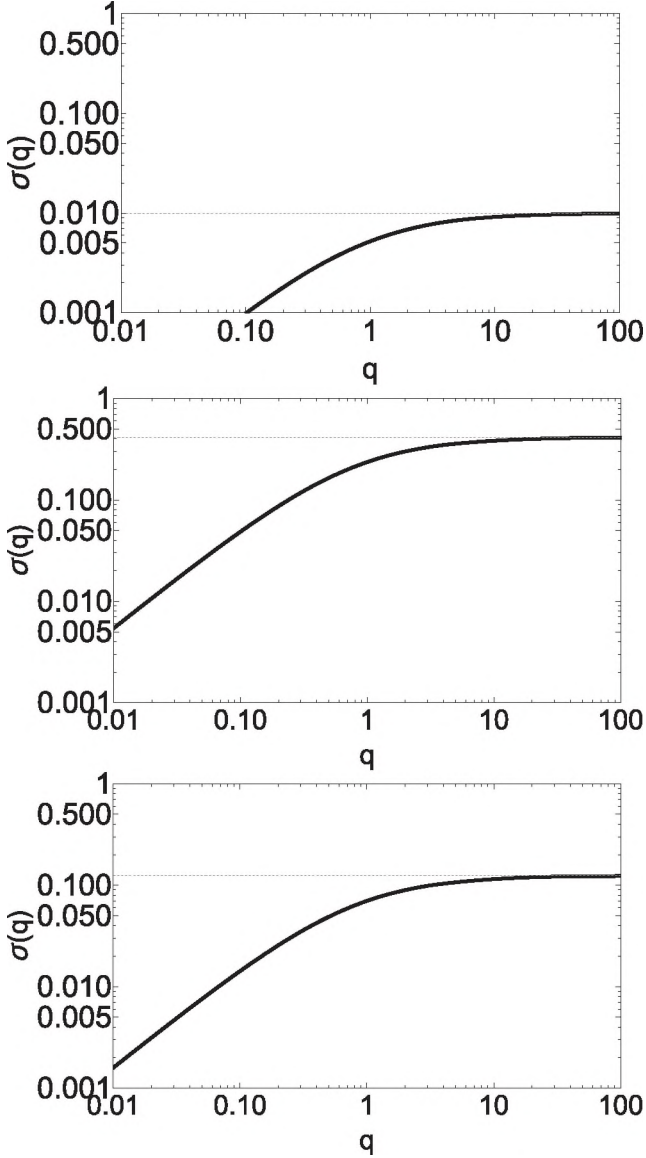


Figure 2. Jet magnetization parameter σ calculated for different boundary condition values $q \equiv \beta_{\text{pl}}^{-1}(1)$; the three panels from top to bottom correspond to the three illustrative sets of $b(x)$ and $\Gamma(x)$ profiles shown in Figure 1.

Interestingly, one can give examples of particular magnetic pressure profiles with multiple maxima across the jet, for which the plasma beta parameter alternates between >1 and <1 , resulting in a substantial jet radial stratification with respect to not only the bulk velocity of the jet plasma, but also the jet plasma magnetization. The question is what relative amount of the total energy of the outflow is carried by the layers with such vastly different magnetizations.

In order to investigate this issue, for the same three illustrative magnetic and velocity profiles, and the boundary conditions $q = 0.1$ and $q = 10$, we also calculate the *cumulative* $\sigma_c(x)$, defined as

$$\sigma_c(x) \equiv \frac{q}{2} \frac{\int_0^x dy h(y) f(y)}{\int_0^x dy h(y) p(y)}, \quad (26)$$

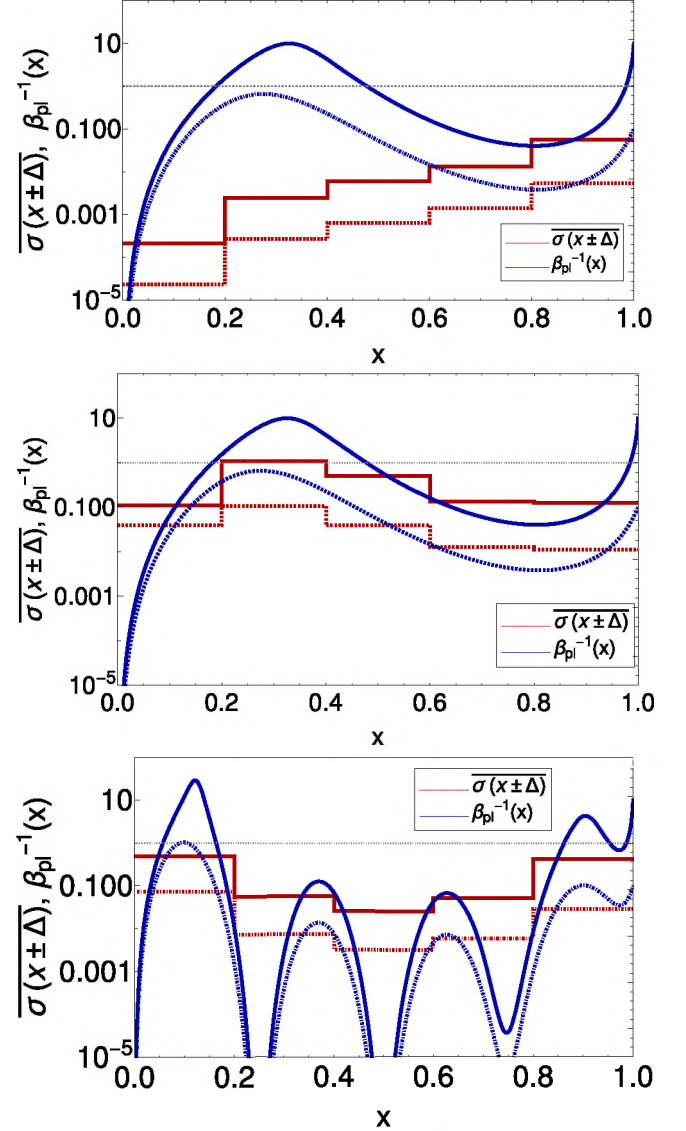


Figure 3. Jet plasma magnetization parameters $\beta_{\text{pl}}^{-1}(x)$ and $\overline{\sigma(x \pm \Delta)}$ with midpoints of averaging $x = \{0.1, 0.3, 0.5, 0.7, 0.9\}$ and $\Delta = 0.1$, calculated for the boundary condition values $q = 0.1$ (dashed lines) and $q = 10$ (solid lines); the three panels from top to bottom correspond to the three illustrative sets of $b(x)$ and $\Gamma(x)$ profiles shown in Figure 1.

as well as

$$\epsilon(x) \equiv \frac{L_p(x) + L_B(x)}{L_p(1) + L_B(1)}, \quad (27)$$

which represents the fraction of the total energy flux integrated from 0 up to a given x . Those are shown in Figure 4. By comparing the $\sigma_c(x)$ and $\epsilon(x)$ profiles with the plasma beta parameter profiles we conclude that there potentially *is* significant energy carried in the highly magnetized jet layers (either boundaries or internal regions), in the sense that the layers with $\beta_{\text{pl}}^{-1} > 1$ may carry a significant amount of the jet total energy flux, from a few up to even tens of a percent.

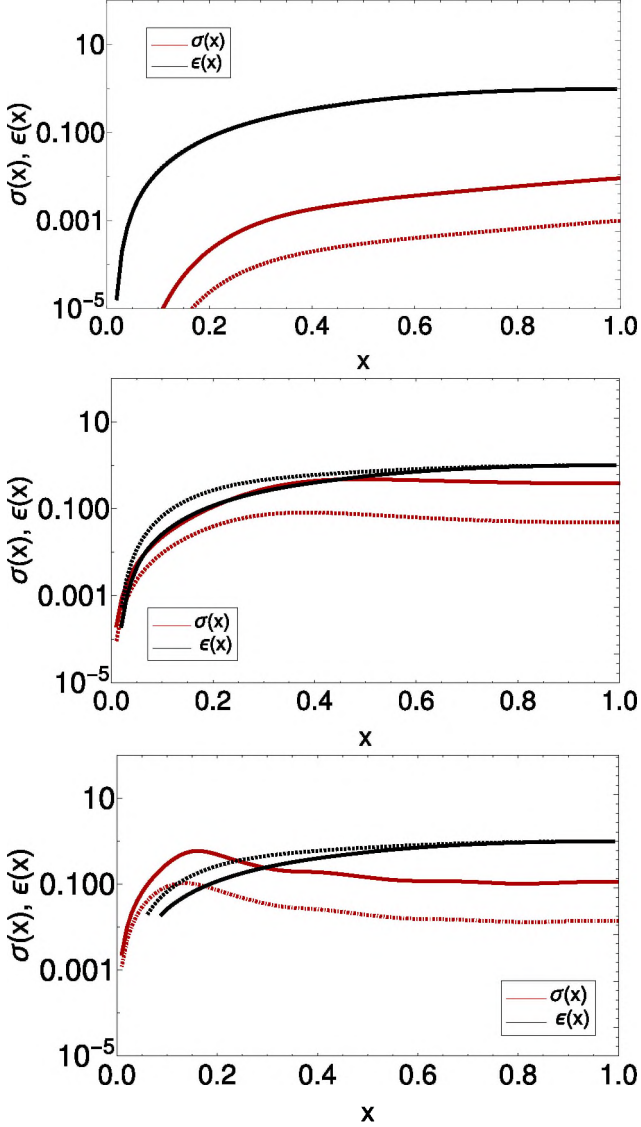


Figure 4. Cumulative $\sigma_c(x)$ parameter and the share of a total energy flux $\epsilon(x)$, calculated for the boundary condition values $q = 0.1$ (dashed lines) and $q = 10$ (solid lines); the three panels from top to bottom correspond to the three illustrative sets of $b(x)$ and $\Gamma(x)$ profiles shown in Figure 1.

4.2. Tangential Discontinuities

In our analysis presented above, we have assumed that the jet magnetic and gaseous radial profiles are continuous. For such, we noted that sharp gradients in the comoving magnetic pressure $f(x)$ often generate $p(x) < 0$. On the other hand, in the framework of the MHD approximation, tangential pressure discontinuities may be present, and those may elevate the overall jet σ parameter without unphysical negative particle pressure.

The simplest parameterization of the $f(x)$ profile allowing for a sudden jump in the comoving magnetic pressure is

$$f(x) = A x^{-\varepsilon} \times H[x - x_0], \quad (28)$$

with $\varepsilon > 0$, where $H[x - x_0]$ is the Heaviside step function; note that the normalization $f(1) = 1$ implies $A \equiv 1$. The magnetohydrostatic equilibrium condition, along with the normalization p

(1) = 1, then gives

$$p(x) = 1 + \frac{2q}{\varepsilon} x_0^{-\varepsilon} - \frac{q(2 - \varepsilon)}{\varepsilon} - \frac{2q}{\varepsilon} x_0^{-\varepsilon} H[x - x_0] + \frac{q(2 - \varepsilon)}{\varepsilon} x^{-\varepsilon} H[x - x_0], \quad (29)$$

with $\varepsilon < 2$. The discontinuity in the comoving magnetic field profile at $x = x_0$ therefore corresponds to the discontinuity in the gaseous pressure. For example, $\varepsilon = 1$ corresponds to $f(x) = 0$ with $p(x) = 1 - q(1 - 2x_0^{-1})$ for $x < x_0$, and $f(x) = x^{-1}$ with $p(x) = 1 - q(1 - x^{-1})$ for $x > x_0$; at the same time, the total jet pressure $P_{\text{tot}} \propto p(x) + qf(x)$ is continuous across the jet.

Let us moreover assume a uniform bulk velocity across the jet, $\Gamma(x) = \text{const}$, so that $h(x) \propto x$. This assumption is handy for two different reasons. First, it significantly simplifies calculations, and second, a negligible velocity shear in fact always maximizes the overall jet σ value (since the toroidal magnetic field pressure is typically larger in the outer layers of a jet). With such, we obtain

$$\begin{aligned} \sigma &= \frac{q}{2} \frac{\int_0^1 dx h(x) f(x)}{\int_0^1 dx h(x) p(x)} = \frac{q}{1+q} \frac{1 - x_0^{2-\varepsilon}}{2 - \varepsilon} \\ &\equiv \Delta p \times \frac{x_0^\varepsilon - x_0^2}{(1+q)(2-\varepsilon)}, \end{aligned} \quad (30)$$

where the gaseous pressure jump at $x = x_0$ is

$$\Delta p \equiv \lim_{x \rightarrow x_0^-} p(x) - \lim_{x \rightarrow x_0^+} p(x) = q x_0^{-\varepsilon}. \quad (31)$$

Clearly, in the above the σ parameter is always smaller than unity for $\varepsilon \leq 1$. In the range $\varepsilon \in (1, 2)$, for a given set of the q and x_0 values, σ increases with ε approaching 2, and may occasionally exceed unity. Considering, however, that $\lim_{\varepsilon \rightarrow 2} \sigma = -\frac{q}{1+q} \ln x_0$, and also $\frac{q}{1+q} \sim 1$ for large $q > 1$, we see that the jet magnetization increases only very slowly (logarithmically) with decreasing x_0 . At the same time, the jump in the gaseous pressure Δp increases very rapidly with decreasing x_0 , ε approaching 2, and finally with increasing q , as illustrated in Figures 5 and 6. For example, with $\varepsilon = 1.9$, $q = 10$, and $x_0 = 0.01$, we have $\sigma \simeq 3.3$ and $\Delta p \sim 3 \times 10^3$; lowering x_0 by one order of magnitude results in only slightly higher $\sigma \simeq 4.5$, but dramatically larger $\Delta p \sim 3 \times 10^5$.

All in all, we conclude that by allowing tangential discontinuities in the gaseous and magnetic pressure profiles, and at the same time diminishing radial velocity shear, one can elevate the jet magnetization parameter above unity. However, values of $\sigma \lesssim \mathcal{O}(10)$ require pressure jumps of several orders of magnitude, which could hardly be considered realistic.⁷

⁷ In fact, there is an intermediate class of profiles, for which the jet gaseous and magnetic pressures, albeit continuous, change with radius by orders of magnitude, while the jet magnetization parameter approaches/exceeds unity. In particular, consider a simple parameterization $f(x) = (1+a)x^2/(1+ax^4)$ and a constant jet bulk Lorentz factor across the outflow. For such, $p(x) = 1 + q + (1+a)q/\sqrt{a} [\tan^{-1}(\sqrt{a}) - \tan^{-1}(x^2\sqrt{a})] - qf(x)$ and $\sigma = \frac{q}{1+q} (1+a)(4a)^{-1} \ln(1+a)$, so that with $q \rightarrow \infty$ and $a > 50$ one can indeed obtain $\sigma \gtrsim 1$; for example, with $a = 3000$ and $q = 10,000$, one has $\sigma \simeq 2$ and $p(0) \simeq 10^6$.

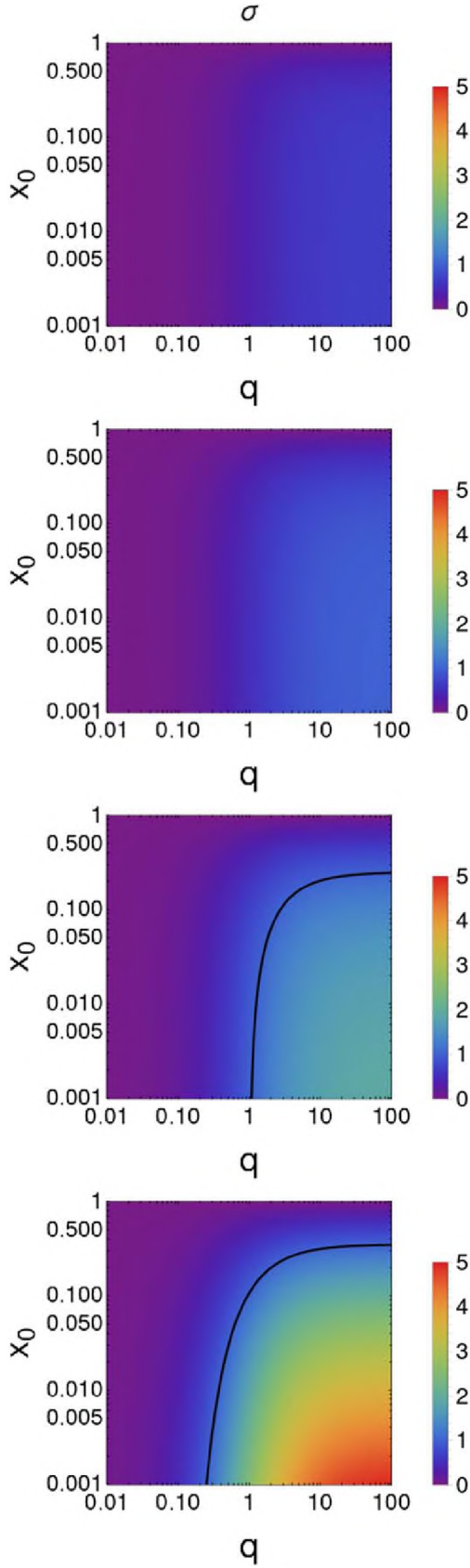


Figure 5. Density plots of σ given in Equation (30) as functions of x_0 and q , for the ε values of 0.5, 1.0, 1.5, and 1.9 (from top to bottom, respectively). Black contours correspond to $\sigma = 1$.

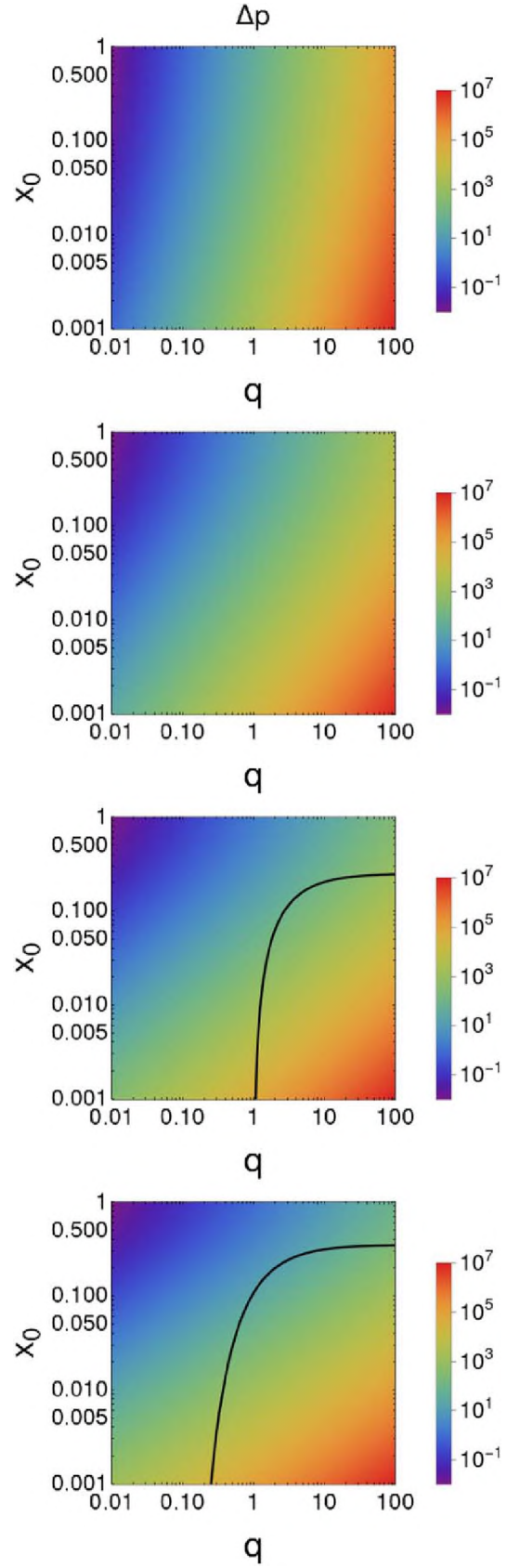


Figure 6. Density plots of Δp given in Equation (31) as functions of x_0 and q , for the ε values of 0.5, 1.0, 1.5, and 1.9 (from top to bottom, respectively). Black contours correspond to $\sigma = 1$.

5. Discussion and Summary

In this paper we analyze the magnetization of a relativistic jet at large distances from the launching site, where it can be considered as fully formed, i.e., accelerated to terminal bulk velocity. In our simple model, we consider perfectly cylindrical jet geometry, purely toroidal configurations for the jet magnetic field, monotonic radial bulk velocity shear, and ultrarelativistic equation of state for the jet particles. We show analytically that, as long as the jet plasma is in magnetohydrostatic equilibrium, and the pressure radial profiles are continuous with the comoving magnetic pressure attaining its maximum at the jet boundary, the ratio of the electromagnetic to particle energy fluxes, both integrated over the jet cross-section area, has to be below unity, $\sigma < 1$.

More complex cases, in particular those with global maxima of the magnetic pressure located well within the jet body, could only be explored numerically. For such, we found that sharp gradients in the comoving magnetic pressure often lead to unphysical solutions with negative particle pressure for some ranges of the jet radius. But when the particle pressure is positive everywhere, the condition $\sigma < 1$ tends to hold anyway. At the same time, however, for certain magnetic and bulk velocity profiles, magnetic pressure may still dominate over particle pressure for certain ranges of cylindrical radius within the jet. In other words, even though a current-carrying outflow as a whole is dominated by the kinetic energy flux of the jet particles, there may be extended domains carrying a significant fraction of the total jet energy, in which the comoving jet magnetic pressure dominates over the gaseous pressure, while the jet Poynting flux is comparable to the particle energy flux. This finding may be relevant in the context of jet particle acceleration processes, since energy dissipation in relativistic plasma proceeds rather differently depending on the plasma magnetization.

Recent particle-in-cell simulations find that magnetic reconnection may proceed very efficiently in the regime of high magnetization, $\beta_{\text{pl}} < 1$ and $\sigma > 1$, where shocks are expected to be weakened by strong magnetic forces. In both pair and electron-ion plasmas, these simulations exhibit efficient nonthermal particle acceleration, with maximum energies and power-law indices that depend on the particular values of the plasma magnetization parameters (e.g., Sironi & Spitkovsky 2014; Guo et al. 2015, 2016; Werner et al. 2016, 2018; Petropoulou et al. 2019).

Our analysis therefore identifies an interesting possibility for astrophysical jets to be characterized by radial stratification with respect to both the plasma bulk velocity and the plasma magnetization. This includes not only the case of a particle-dominated jet spine surrounded by a magnetically dominated boundary layer, but also the possibility of alternating-current jets consisting of layers with low and high values of σ and β_{pl} . In such outflows, shocks and magnetic reconnection may dominate alternately, resulting in the formation of highly inhomogeneous distributions of radiating particles (with respect to maximum particle energies and particle spectral indices). When combined with relativistic beaming effects related to the radial velocity shear (see, e.g., Komissarov 1990; Stawarz & Ostrowski 2002; Aloy & Mimica 2008), this may lead to jets having vastly diverse appearances to an observer, depending on the jet viewing angle.

In the framework of the analyzed simple jet models with purely toroidal magnetic field, we have found that the jet

magnetization parameter can be elevated up to relatively modest values $\sigma \lesssim \mathcal{O}(10)$ only in the case of extreme gradients or discontinuities in the gas pressure, and a significantly suppressed velocity shear. Such cases would therefore correspond to a narrow, unmagnetized jet spine, surrounded by an extended, essentially force-free layer with a toroidal field, both characterized by comparable bulk Lorentz factors. However, in the absence of velocity shear, relativistic outflows with strong toroidal magnetic field (and no poloidal component) are known to be susceptible to current-driven kink instability (e.g., Mizuno et al. 2012; Nalewajko & Begelman 2012; Martí et al. 2016; Kim et al. 2018; Das & Begelman 2019; Sobacchi & Lyubarsky 2019 and references therein; for a recent review on the topic see also Perucho 2019).

Let us therefore comment in this context on the role played by an additional poloidal magnetic field, represented here for simplicity by a purely vertical component B_z . If this component is uniform across and along the jet, $\partial_\phi B_z = \partial_z B_z = \partial_r B_z = 0$, the gaseous pressure profiles following from magnetohydrostatic equilibrium remain unchanged with respect to the ones we have calculated and discussed above, since no additional current component $\mathbf{J}' \propto \nabla' \times \mathbf{B}'$ is associated with such an additional field, and so the magnetohydrostatic equilibrium $\nabla' P \propto \mathbf{J}' \times \mathbf{B}'$ is not affected either. Moreover, the relative increase of the jet Poynting flux is then only in the ϕ direction, so it represents energy circulation around the jet and does not add to the original Poynting flux along the z direction. Hence, there is no net increase in the σ parameter. That is to say, a small amount of a uniform vertical field should stabilize the jet against the current-driven oscillations (see Mizuno et al. 2012; Das & Begelman 2019), but should not affect our conclusions presented above regarding the jet magnetization.

On the other hand, in the presence of pronounced radial gradients in the vertical field, $\partial_r B_z \neq 0$ (but still $\partial_\phi B_z = \partial_z B_z = 0$), the situation may change. That is because in such a case the particle energy flux $L_p \propto \int dr r \beta \Gamma^2 P$ should be affected as the gaseous pressure profile has to adjust in response to the additional azimuthal current component, $4\pi J'_\phi = -c \partial_r B'_z$, following the altered magnetohydrostatic equilibrium, namely, $\partial_r P \propto B_z \partial_r B_z - \frac{1}{2} r^{-2} \partial_r (r^2 B_\phi^2 \Gamma^{-2})$.

We note, however, that the configuration with the vertical field confined within a narrow jet spine, and therefore with strong radial gradients in B_z , is to be expected rather only for fully collimated and accelerated electromagnetic outflows, i.e., the ones for which the Poynting and particle energy fluxes are in a rough equipartition $\sigma \sim \mathcal{O}(1)$ anyway (see Beskin & Nokhrina 2009; Lyubarsky 2009). Moreover, as discussed in Mizuno et al. (2012), relativistic outflows consisting of a poloidal field concentrated toward the jet axis, and a dynamically relevant toroidal field in the outer layer, are highly kink unstable, and as such subjected to efficient dissipation of the magnetic energy. Hence, we conclude that large values of the jet magnetization parameter σ should not be realistically expected in such cases anyway.

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Chapter 6

Summary and Future Prospects

Black holes exhibiting varying levels and forms of activity can be found in various astrophysical systems, which they shape and co-evolve with. In the majority of cases, an exact description of the associated phenomena is impossible due to the complexity of the physical theories involved and, as is often the case in astronomy, limited access to the source. To understand this complex picture, a holistic approach integrating various formalisms, approximations, and techniques is essential. The research projects presented in this Ph.D. dissertation delve into the physics of black holes and their relativistic jets, utilizing appropriate tools across different astrophysical systems. This variety constitutes an excellent starting point for other projects, some of which I am already pursuing.

In particular, we are currently building upon our analysis of relativistic current-carrying jets, by studying the additional drift currents that can be formed by protons accelerated within the velocity shears. By doing so, we hope to introduce a model that will explain the observed helicity preference in the transverse Faraday rotation gradients observed in radio galaxies (Gabuzda, Nagle, and Roche, 2018). Also, our previous analysis of potential micro-lensing events in the blazar light curves, presented some challenges. In the future, we hope to revisit the analysis by investigating the probability of such events based on the number density of potential binary lenses in host galaxies, and the exact kinematics of precessing jets. Secondly, the stochastic nature of the underlying AGN variability presents an obstacle in robustly determining the extrinsic nature of the variability. We hope to approach this issue in a more well-rounded way, by fitting the source light curves in the time domain, in the framework of the continuous auto-regressive model. Finally, through my previous collaboration, I am still involved in GR MHD simulations of accreting black holes and relativistic jets, addressing such problems as, among others, the description of self-gravity in the collapsar model (see Janiuk, Shahamat Dehsorkh, and Król, 2023).

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