Two-temperature, Magnetically Arrested Disc simulations of the jet from the supermassive black hole in M87

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ABSTRACT
We present two-temperature, radiative general relativistic magnetohydrodynamic simulations of Magnetically Arrested Discs (MAD) that launch powerful relativistic jets. The mass accretion rates of our simulations are scaled to match the luminosity of the accretion flow around the supermassive black hole in M87. We consider two sub-grid prescriptions for electron heating: one based on a Landau-damped turbulent cascade, and one based on heating from trans-relativistic magnetic reconnection. The simulations produce jets with power on the order of the observed value for M87. Both simulations produce spectra that are consistent with observations of M87 in the radio, millimetre, and submillimetre. Furthermore, the predicted image core-shifts in both models at frequencies between 15 GHz and 86 GHz are consistent with observations. At 43 and 86 GHz, both simulations produce wide opening angle jets consistent with VLBI images. Both models produce 230 GHz images with distinct black hole shadows that are resolvable by the Event Horizon telescope (EHT), though at a viewing angle of 17°, the 230 GHz images are too large to match EHT observations from 2009 and 2012. The 230 GHz images from the simulations are dynamic on time-scales of months to years, suggesting that repeated EHT observations may be able to detect the motion of rotating magnetic fields at the event horizon.

Key words: accretion, accretion discs – black hole physics – relativistic processes – methods: numerical – galaxies: jets – galaxies: nuclei

1 INTRODUCTION
The core of the giant elliptical galaxy Messier 87 (M87) launches a relativistic jet that extends for several kiloparsecs (Curtis 1918). The jet is dynamic, and it has been observed in wavelengths from radio γ-ray (Abramowski et al. 2012). Very-long-baseline interferometry (VLBI) observations at radio and millimetre wavelengths (e.g. Palmer et al. 1967; Reid et al. 1982; Junor et al. 1999; Kovalev et al. 2007; Ly et al. 2007; Asada & Nakamura 2012; Hada et al. 2016; Mertens et al. 2016; Walker et al. 2018; Kim et al. 2018) show that the jet remains collimated on sub-parsec scales into the heart of the galaxy, where it terminates in a bright radio core. The radio core has a frequency-dependent position, indicative of optically thick emission from the jet that becomes increasingly transparent at higher frequencies (Blandford & Königl 1979). At frequencies higher than ~ 86 GHz, the radio core is coincident with a supermassive black hole (SMBH: Hada et al. 2011). The black hole mass has been measured to be 6.6 × 10⁹ M⊙ from stellar dynamics in the surrounding galaxy nucleus (for a distance of D = 17.9 Mpc; Gebhardt et al. 2011). The mass measured from gas dynamics is a factor of two smaller (Walsh et al. 2013).

Making images of the horizon-scale structure around the black holes in M87 and the Galactic Centre (Sgr A*) is a primary objective of the Event Horizon Telescope (EHT), a global millimetre VLBI array with an imaging resolution of ~ 15 μas at 230 GHz (Doeleman et al. 2009). The EHT has already constrained the compact structure in the core of M87 to the scale of a few Schwarzschild radii (Doeleman et al. 2012; Akiyama et al. 2015). Assuming a distance to the black hole of D = 16.7 Mpc (Mei et al. 2007) and a black hole mass of 6.2 × 10⁹ M⊙ (Gebhardt et al. 2011, scaled for this distance), M87’s black hole “shadow” (i.e., the lensed photon orbit; Bardeen et al. 1972) should have a diameter of ~ 40 μas, making it accessible to imaging by the full EHT (Lu et al. 2014; Akiyama et al. 2017; Chael et al. 2018b).

At the core of M87 is a Low Luminosity Active Galactic Nucleus (LLAGN), with a luminosity many orders of magnitude below the Eddington limit. LLAGN like M87 constitute the bulk of the SMBH population in the local universe (Greene & Ho 2007; Ho 2008). Their low luminosity is natu-
rally produced by an Advection-Dominated Accretion Flow (ADAF: Ichimaru 1977; Rees et al. 1982; Abramowicz et al. 1995; Narayan & Yi 1995a,b; Mahadevan 1998; Blandford & Begelman 1999). In contrast to the optically thick, geometrically thin accretion discs found in bright AGN (Shakura & Sunyaev 1973), ADAFs in LLAGN are optically thin and geometrically thick (see Yuan & Narayan 2014 for a review). Jets like that in M87 are likely powered by the black hole’s rotational energy, which is extracted by ordered magnetic fields threading the black hole event horizon (Blandford & Znajek 1977; Tchekhovskoy et al. 2011; Zamaninasab et al. 2014).

The dynamics of hot accretion flows and the formation of jets have been extensively investigated using grid-based general relativistic magnetohydrodynamic (GRMHD) and radiative magnetohydrodynamic (GRRMHD) codes (e.g. Komissarov 1999; De Villiers et al. 2003; Gammie et al. 2003; Tchekhovskoy et al. 2010, 2011; Sadowski et al. 2013a, 2014; McKinney et al. 2014; Ryan et al. 2015; White et al. 2016; Chandra et al. 2017). These simulations have demonstrated that jets powered by the black hole spin can be launched from thick discs and accelerated to high Lorentz factors (McKinney 2006; Komissarov et al. 2007; McKinney & Blandford 2009; Liska et al. 2018). In the specific case of M87, Dexter et al. (2012), Moscibrodzka et al. (2016, 2017), and Ryan et al. (2018) have investigated the spectra and 230 GHz images predicted from various GRRMHD simulations.

In hot, low-density, low-accretion rate flows like that around M87 (and Sgr A*), Coulomb coupling between electrons and ions is inefficient (Mahadevan 1998). Electrons and ions will not be in thermal equilibrium with each other, and in most cases the temperature of the single fluid evolved by a GRMHD or GRRMHD simulation will be dominated by the ion temperature, leaving the temperature of the emitting electrons unconstrained. In a simulation, the electron-to-gas temperature ratio is typically set manually during postprocessing (e.g. Shcherbakov et al. 2012; Moscibrodzka et al. 2014; Chan et al. 2015 for Sgr A* and Dexter et al. 2012; Moscibrodzka et al. 2016 for M87). In this approach, the electron temperature is usually adjusted to find the best fit to the measured spectrum. Nonthermal particle distributions can also be added (e.g. Dexter et al. 2012; Ball et al. 2016; Davelaar et al. 2018).

Another approach is to self-consistently evolve ions and electrons, each with its own thermodynamics and interactions. This approach has been pursued with several different codes (Ressler et al. 2015; Sadowski et al. 2017; Ryan et al. 2017), all of which evolve a thermal electron population alongside the other standard GRMHD or GRRMHD fluid variables. These methods have been used in recent simulations of Sgr A* (Ressler et al. 2017; Chael et al. 2018a) and M87 (Ryan et al. 2018).

A substantial source of uncertainty in these simulations arises from the fact that the physical processes that govern viscous dissipation occur on scales much smaller than the finest scales resolved by the grid. The range of physical processes that govern dissipation at these smallest scales in hot accretion flows is still relatively unconstrained. In Chael et al. (2018a), we investigated two candidates for the origin of viscous dissipation in simulations of Sgr A*: the Landau-damped turbulent cascade heating prescription of Howes (2010) and a prescription for magnetic reconnection heating obtained from particle-in-cell (PIC) simulation data (Rowan et al. 2017). We showed that predictions for the image morphology at 230 GHz and longer wavelengths depend strongly on the choice of heating prescription. For Sgr A*, the turbulent cascade prescription results in a natural “disc-jet” image morphology at longer wavelengths, while the reconnection prescription produces isotropic, disc-dominated emission (Chael et al. 2018a; Issaoun et al. 2018).

Recently, Ryan et al. (2018) carried out axisymmetric simulations of M87 with two-temperature evolution and frequency-dependent radiative transport. They found that, because M87 is more radiatively efficient than Sgr A*, including radiation in the simulation along with temperature evolution of the electrons is critical. They performed simulations at both low and high values of the SMBH mass and found that the accretion flow in the high mass, high spin case produced a 230 GHz image consistent with published EHT observations. However, their simulations were performed with weak values of magnetic flux threading the horizon. As a result, the jets in their simulations had a narrower opening angle than that observed in VLBI images of M87, and the jet power was lower than the measured value by several orders of magnitude.

Magnetically Arrested Discs (MADs; Bisnovatyi-Kogan & Ruzmaikin 1976; Narayan et al. 2003) represent the opposite limit to the weak magnetic flux mode of black hole accretion explored in Ryan et al. (2018). In these systems, coherent magnetic flux builds up on the black hole, both launching powerful jets and limiting accretion via magnetic pressure. In GRRMHD simulations (Igumenshchev et al. 2003; Tchekhovskoy et al. 2011; McKinney et al. 2012; Narayan et al. 2012; Sadowski et al. 2013b), MADs launch jets powered by the black hole spin with wide opening angles and large jet powers. The measured jet power of M87 is large ($\sim 10^{43} - 10^{44}$ erg s$^{-1}$; Reynolds et al. 1996; Owen et al. 2000; Stawarz et al. 2006; de Gasperin et al. 2012), and the jet is launched with a wide opening angle ($\sim 55^\circ$ at 43 GHz; Walker et al. 2018). Thus, there is reason to think that M87 has a magnetically arrested disc at its core.

In this paper we present the results of two fully 3D, two-temperature GRRMHD simulations of Magnetically Arrested Discs around the black hole in M87 performed using the code KORAL (Sadowski et al. 2013a, 2014, 2017). In these simulations, we again compare the Landau-damped turbulent cascade heating prescription from Howes (2010) with the magnetic reconnection prescription from Rowan et al. (2017); Chael et al. (2018a). Both simulations are performed assuming a SMBH mass of $6.2 \times 10^8 M_\odot$ (Gebhardt et al. 2011, scaled for a distance of 16.7 Mpc; Mei et al. 2007) and a dimensionless spin of $a = 0.9375$. These simulations are the first Magnetically Arrested Discs evolved with two-temperature electron-ion thermodynamics.

In Section 2, we briefly review the method used in Sadowski et al. (2017) and Chael et al. (2018a) for evolving thermal electron and ion entropies, as well as the two heating prescriptions considered in this paper, and we describe the setup of our simulations in Section 3. We present our results in Section 4. In Section 4.1, we first discuss the time-averaged characteristics of our models, including the accretion rate, jet power, and distributions of temperature, magnetic field, and radiation. We then present spectra for our simulations in Section 4.2 and discuss variability in the
2 EQUATIONS

2.1 Evolution Equations

In this section, we briefly review the method used in the GRRMHD code KORAL to a two-temperature fluid (Sadowski et al. 2017, also summarized in Chael et al. 2018a).

We consider two fluids in spacetime, electrons and ions (assumed for the remainder of this work to be entirely ionized Hydrogen). Charge neutrality demands these fluids have the same number density $n$ and four velocity $u^\mu$ everywhere, but lacking efficient processes to bring them into equilibrium $T_e \neq T_i$.

Together, electrons and ions form a single mixed fluid, which is characterized by a mass density dominated by the ions, $\rho = m_p n$, and a total internal energy density $u = u_e + u_i$. The total pressure $p = p_e + p_i$ can then be expressed with an effective adiabatic index; $\Gamma_{gas} = (\Gamma_{gas} - 1)u$. In a hot accretion flow, which typically has temperatures $>10^{13}$ K in the inner regions, electrons become relativistic and their adiabatic index can decrease from 5/3 towards 4/3. In the innermost regions where $T_i > 10^{13}$ K, even the ions become quasi-relativistic. Thus, the effective gas adiabatic index $\Gamma_{gas}$ takes on values in the range $4/3 \leq \Gamma_{gas} \leq 5/3$ depending on the local temperatures and energy densities of the two component species (see Sadowski et al. 2017 for the form of the adiabatic indices as a function of temperature used in KORAL).

The MHD stress-energy tensor $T^\mu_\nu$ consists of contributions from the fluid variables as well as the magnetic field four-vector $b^\mu$:

$$T^\mu_\nu = (p + u + p_b^2) u^\mu u^\nu + \left( p + \frac{1}{2} b^2 \right) \delta^\mu_\nu - b^\mu b^\nu.$$  \hspace{1cm} (1)

KORAL treats the frequency-integrated radiation field $R^\mu_\nu$, as a second perfect fluid. The radiation field is described by its rest frame energy density $\bar{E}$ and its four-velocity $R^\mu_{\nu} \neq u^\nu$:

$$R^\mu_\nu = \frac{4}{3} \bar{E} n^\mu R^\nu_\nu + \frac{1}{3} \bar{E} b^\mu b^\nu.$$  \hspace{1cm} (2)

In this formulation (M1 closure), radiation is described at each spacetime point by four bolometric quantities: $\bar{E}$ and the three $u^\mu$. We also track a fifth quantity, the photon number density $\hat{n}_R$, which encodes information about the mean photon frequency $E/\hbar \omega_R$. In contrast to this frequency-integrated approach, Ryan et al. (2015, 2017, 2018) use a Monte-Carlo approach for including radiation in their simulations. Their approach represents the radiation field with many individual particle “superphotons” with different frequencies that are emitted and absorbed in between the fluid evolution timesteps.

The set of GRRMHD equations for evolving the total fluid, the magnetic field, the frequency-integrated radiation field, and the photon number are (Gammie et al. 2003; Sadowski et al. 2014; Sadowski & Narayan 2015)

$$(p u^\mu)_\nu = 0,$$  \hspace{1cm} (3)

$$T^\mu_{\nu,\mu} = G_{\nu},$$  \hspace{1cm} (4)

$$R^\mu_{\nu,\mu} = -G_{\nu},$$  \hspace{1cm} (5)

$$(\hat{n}_R u^\mu_R)_\nu = \hat{n}_R,$$  \hspace{1cm} (6)

$$F^{\mu}_{\nu,\mu} = 0,$$  \hspace{1cm} (7)

where $F^{\mu\nu} = b^\mu u^\nu - b^\nu u^\mu$ is the dual of the MHD Maxwell tensor, $G_{\nu}$ is the four-force density that couples the radiation and gas (see Sadowski et al. 2017 for the precise form), and $\hat{n}_R$ is the frame-invariant photon production rate (see Sadowski & Narayan 2015).

In evolving electrons and ions, we consider the entropy per particle of each, $s_e$ and $s_i$. The temperatures $T_e$, $T_i$ and energy densities $u_e$, $u_i$ are functions of the species entropy and number density (see Sadowski et al. 2017 and the Appendix of Chael et al. 2018a). The species entropies are evolved using the first law of thermodynamics:

$$T_e (n_i u_i u^\mu)_{\nu,\mu} = \hat{\delta} q^\gamma + G^\beta,$$  \hspace{1cm} (8)

$$T_i (n_i u_i u^\mu)_{\nu,\mu} = (1 - \delta_i) q^\gamma - q^\beta,$$  \hspace{1cm} (9)

where $q^\gamma$ is the dissipative heating rate, $\delta_i$ is the fraction of the dissipative heating that goes into electrons, $q^\beta$ is the energy exchange rate from ions to electrons due to Coulomb coupling (Stepney & Guilbert 1983), and $G^\beta$ is the radiative cooling rate.

The physical processes that produce dissipation occur at scales far smaller than the simulation grid. We identify the total dissipative heating $q^\gamma$ numerically by evolving the thermal entropies adiabatically over a timestep $\Delta \tau$. We then compare the sum of the adiabatically evolved energy densities, $u_{i,\text{adiab}}$ and $u_{e,\text{adiab}}$, to the separately-evolved total gas energy $u$, thereby estimating the dissipative heating in the total fluid:

$$q^\gamma = \frac{1}{\Delta \tau} [u - u_{i,\text{adiab}} - u_{e,\text{adiab}}].$$  \hspace{1cm} (10)

The fraction $\delta_i$ of the heating that goes into the electrons, however, must be determined by a sub-grid prescription.

2.2 Electron Heating Prescriptions

In this work, we again consider the two sub-grid prescriptions that we previously applied to simulations of Sgr A* (Chael et al. 2018a). These prescriptions depend on three parameters: the “plasma-beta” $\beta$, the magnetization $\sigma$, and the temperature ratio $T_e/T_i$.

The plasma-beta parameter $\beta$ is the ratio of the thermal ion pressure to the magnetic pressure:

$$\beta_i = \frac{8 \pi n_i k T_i}{|B|^2},$$  \hspace{1cm} (11)

and the magnetization $\sigma_i$ compares the magnetic energy...
density to the rest-mass energy density of the fluid:

\[ \sigma_i = \frac{|B|^2}{4\pi n_i m_i c^2}. \]  \hspace{1cm} (12)

While in general \( \sigma_i \ll 1 \) in the disc, in our MAD simulations \( \sigma_i \) exceeds unity in the jet as well as close to the black hole.

The two heating prescriptions we consider are plotted as a function of different plasma parameters in Fig. 1. The first prescription for \( \delta_e \) is taken from calculations of the Landau damping of a MHD turbulent cascade in a weakly collisional plasma (Howes 2010). This prescription is based on nonrelativistic calculations with \( \sigma_i \ll 1 \) (Howes 2008a,b), and while it matches solar wind measurements (Howes 2011), it may not be well-adapted to relativistic systems like the M87 accretion flow. Recently, however, Kawazura et al. (2018) performed numerical simulations of the turbulent damping process that indicate the qualitative behavior of the Howes (2010) prescription holds even in relativistic plasmas.

The Howes (2010) turbulent cascade prescription is a function of the temperature ratio and \( \beta_i \). It most strongly depends on \( \beta_i \), sharply transitioning from dumping most of the turbulent heating in electrons (\( \delta_e \approx 1 \)) at low \( \beta_i \), to primarily heating ions (\( \delta_e \approx 0 \)) at high \( \beta_i \). While in general we expect radiation to cool electrons to lower temperatures than ions, in Chael et al. (2018a) we saw that using this prescription results in an electron temperature that can greatly exceed the ion temperature in the jet region (where \( \beta_i \ll 1 \)).

Our second prescription was obtained by measuring electron heating in particle-in-cell (PIC) simulations of trans-relativistic magnetic reconnection (Rowan et al. 2017; Chael et al. 2018a). Magnetic reconnection heating arises from reconnection events at small scales of the cascade of MHD turbulence (Carbone et al. 1990; Boldyrev 2006; Boldyrev & Loureiro 2017; Loureiro & Boldyrev 2017). Rowan et al. (2017) explored magnetic reconnection with PIC simulations in plasmas over a trans-relativistic range of temperatures and magnetic field strengths (see also Werner et al. 2018 for a similar study). A fitting function to \( \delta_e \) as measured in the reconnection PIC simulations of Rowan et al. (2017) was presented in Chael et al. (2018a). Note that the simulations in Rowan et al. (2017) had no guide field perpendicular to the reconnection; the efficiency of electron heating in the strong guide field regime is yet to be studied in detail and may be qualitatively different than in the zero guide field case.

The reconnection prescription differs substantially from the turbulent cascade prescription. At a fixed temperature, decreasing \( \beta_i \) results in energy being shared equally between the species downstream from the reconnection region. Thus, \( \delta_e \) approaches a maximum of 0.5 in the most magnetically dominated plasmas. We should thus never see \( T_e > T_i \) when using this prescription (Chael et al. 2018a). At the other limit of large \( \beta_i \), the heating fraction approaches a constant value that depends only on \( \sigma_w \) (defined with respect to the fluid enthalpy \( w \))\(^1\), which is nonzero even for \( \sigma_w < 10^{-3} \). This behavior is qualitatively different from the Howes (2010) prescription, which sends \( \delta_e \rightarrow 0 \) for similarly large values of \( \beta_i \). In the Sgr A* simulations of Chael et al. (2018a), the floor on \( \delta_e \) in the reconnection prescription resulted in hotter discs than in the simulations that used the turbulent prescription.

\(^1\) The PIC simulations of Rowan et al. (2017) define \( \sigma \) with respect to the fluid enthalpy \( w \) instead of only including the rest mass in the denominator, as in Eq. 12. See Section 2.2 of Chael et al. (2018a) for a discussion of the difficulties of simply interpreting \( \sigma_w \) in terms of \( \sigma_i \) in relativistic plasmas.
3 NUMERICAL SIMULATIONS

3.1 Units

In both simulations presented in this work, we take the distance to M87 as $D = 16.7$ Mpc (Mei et al. 2007) and fix the black hole mass to $6.2 \times 10^8 M_\odot$ (Gebhardt et al. 2011, scaled for this distance).

For this mass, the gravitational length scale of M87 is $r_g = GM/c^2 = 9.2 \times 10^{14}$ cm = 61 AU. The corresponding angular scale is $r_g/D = 3.7 \mu$mas. The gravitational time-scale is $t_g = r_g/c = 3 \times 10^5$ s = 8.5 hr.

M87’s Eddington luminosity is $L_{\text{Edd}} = 7.8 \times 10^{47}$ erg s$^{-1}$. The Eddington accretion rate is $\dot{M}_{\text{Edd}} = L_{\text{Edd}}/\eta c^2 = 77 M_\odot$ yr$^{-1}$, where for our chosen value of spin, we set the efficiency $\eta = 0.18$, as expected for a thin accretion disc with $a = 0.9375$ (Novikov & Thorne 1973).

3.2 Simulation Setup

Our simulations were performed in the Kerr metric using a modified Kerr-Schild coordinate grid that is exponential in radius and concentrates grid cells near the equator (see the Appendix of Chael et al. 2018a for the transformation between our coordinates and standard Kerr-Schild coordinates). We used a resolution of $288 \times 224 \times 128$ cells in the $r$, $\theta$, and $\phi$ directions, respectively, which well-resolves the magnetorotational instability (MRI) and enabling accretion. To capture the evolution of the jet at large radii, we set the outer boundary of the simulation box at $10^5 r_g$.

We set up initial equilibrium gas torii using the model of Penna et al. (2013). To build up magnetic field to the point where the disc reaches the saturation value of magnetic flux to become magnetically arrested, we initialized the torus with a single weak ($\beta_{\text{max}} = 100$) magnetic field loop centered around $r \approx 50 r_g$. The initial energy in electrons was set at one per cent of the total gas energy, with the remainder in ions.

KORAL solves Eqs. (8) and (9) for the electron and ion thermodynamics in parallel with the conservation Eqs. (3) – (6) for the matter and radiation fluids and the magnetic field induction equation Eq. 7. The advection of quantities across cell walls is computed explicitly by reconstructing the appropriate fluxes at the cell walls using the second-order piecewise parabolic method (PPM). Source and coupling terms in the evolution equations are then applied implicitly at each cell center using a Newton-Raphson solver (Sadowski et al. 2013a, 2014, 2017).

In the jet region, high fluid velocities rapidly evacuate the funnel and cause the fluid density to drop without bound. In order to ensure the numerical stability of our simulations, we put a global ceiling on the magnetization $\sigma_r$, as measured in the zero angular momentum observer (ZAMO) frame (McKinney et al. 2012). In this frame, the fluid density is increased to bring the magnetization back to our chosen limit, $\sigma_r^{\text{max}} = 100$. In nature, pair-production of electrons and positrons may populate the jet (Moscibrodzka et al. 2011; Broderick & Tchekhovskoy 2015). In the nearest two cells to the polar axis, we control numerical instability from fluid flow across the poles by replacing the value of $u^\theta$ with the value from the third cell at the end of each timestep.

We evolved one simulation using the Howes (2010) prescription for dividing viscous dissipation between electrons and ions; we refer to this simulation as H10. We evolved the other simulation using the magnetic reconnection prescription of Rowan et al. (2017); Chael et al. (2018a); we refer to this simulation as R17. We first ran the two simulations for $10^6 t_g$ in 3D. During this this time, both simulations formed a thick disc at small radii and accumulated magnetic flux on the black hole horizon that exceeds the MAD threshold of $\sim 50 \sqrt{M_{\odot} r_g}$ (Tchekhovskoy et al. 2011; McKinney et al. 2012). At this point, we rescaled the gas density and magnetic field (keeping the temperatures and magnetization fixed), so that the 230 GHz flux from the models was approximately equal to $0.98 \pm 0.04$ Jy of compact flux density measured by the EHT in 2009 and 2012 (Doeleman et al. 2012; Akiyama et al. 2015). We then ran the simulations from the rescaling point for another $5000 t_g$. For both models, over the final $5000 t_g$ period, the inflow equilibrium region in the disc (Narayan et al. 2012) extends to $\sim 40 r_g$. In the fast moving jet, the region of outflow equilibrium extends to nearly $10^4 r_g$.

3.3 Radiative Transfer

We produced spectra, images, and lightcurves from our simulations using two post-processing codes. For computing...
spectra, we used HEROIC, (Zhu et al. 2015; Narayan et al. 2016), a code that solves for the spectrum and angular distribution of radiation at each grid position self-consistently. Inverse Compton scattering is included along with free-free (from both e- e and e- i interactions) and synchrotron emission and absorption. At millimetre and radio wavelengths, synchrotron radiation dominates the emission. To produce high-resolution images and lightcurves in the millimetre and lower frequencies, we used the ray tracing and emission and absorption. At millimetre and radio wave-lengths, synchrotron radiation dominates the emission. We use the “fast-light” approximation, where radiation is produced from each simulation snapshot independently, ignoring the evolution of the fluid as the photons propagate.

The jet inclination angle of M87 is constrained from observed “super-luminal” motion of jet components in VLBI images (Heinz & Begelman 1997). In the present paper, we use an inclination angle of 17° (Mertens et al. 2016; Walker et al. 2018). We choose to measure this angle up from the lower pole, so that the sense of rotation of the accretion disc and black hole spin is clockwise on the sky. This is the preferred orientation of the jet angular momentum vector as determined by the differential brightening and pattern velocities of the jet limbs in VLBI images (Walker et al. 2018). To match the orientation of the M87 jet on the sky at ~72° east of north (Reid et al. 1982), we rotate our computed images 108° counterclockwise.

When calculating spectra and images, we eliminate all regions of the simulation that have $\sigma_i > \sigma_{\text{cut}}$, where we choose to set $\sigma_{\text{cut}} = 25$. The reason for this cut can be seen in Fig. 2, which shows profiles of $\rho$ and $\sigma_i$ versus polar angle $\theta$ in the time- and azimuth-averaged simulation data. The gas density levels off at a floor value inside the polar region. In model $H10$, $\delta_e \approx 1$ everywhere inside the jet region defined by the $Be = 0.05$ contour. Outside the highly magnetized funnel, $\delta_e$ drops to nearly zero in the outer regions of the disc, $r \gtrsim 40 r_g$. In contrast, in model $R17$, $\delta_e$ reaches its limit of equipartition of thermal energy ($\delta_e = 0.5$) inside the $\sigma_i = 1$ contour. In the less magnetized disc outside $r \gtrsim 15 r_g$, $\delta_e$ falls to a small but nonzero value $\delta_e \approx 0.2$.

While the temperature distribution of the combined gas is similar in the two models (the second column of Fig. 3), the different asymptotic behaviors of the two heating prescriptions introduced in Sec. 2.2 are evident in the average values of $\delta_e$ in the jet region. In model $H10$, $\delta_e \approx 1$ everywhere inside the jet region defined by the $Be = 0.05$ contour. Outside the highly magnetized funnel, $\delta_e$ drops to nearly zero in the outer regions of the disc, $r \gtrsim 40 r_g$. In contrast, in model $R17$, $\delta_e$ reaches its limit of equipartition of thermal energy ($\delta_e = 0.5$) inside the $\sigma_i = 1$ contour. In the less magnetized disc outside $r \gtrsim 15 r_g$, $\delta_e$ falls to a small but nonzero value $\delta_e \approx 0.2$.

In the leftmost panels of Fig. 3, the different asymptotic behaviors of the two heating prescriptions introduced in Sec. 2.2 are evident in the average values of $\delta_e$ in the jet region. In model $H10$, $\delta_e \approx 1$ everywhere inside the jet region defined by the $Be = 0.05$ contour. Outside the highly magnetized funnel, $\delta_e$ drops to nearly zero in the outer regions of the disc, $r \gtrsim 40 r_g$. In contrast, in model $R17$, $\delta_e$ reaches its limit of equipartition of thermal energy ($\delta_e = 0.5$) inside the $\sigma_i = 1$ contour. In the less magnetized disc outside $r \gtrsim 15 r_g$, $\delta_e$ falls to a small but nonzero value $\delta_e \approx 0.2$.

While the temperature distribution of the combined gas is similar in the two models (the second column of Fig. 3), the different heating prescriptions result in different electron temperatures and temperature ratios in the inner disc and jet. In model $H10$, the deposit of nearly all thermal energy into electrons inside the jet results in high electron temperatures $T_e \sim 10^{11}$ K near the black hole that climb to $10^{12}$ K in the jet around $50 r_g$. While the temperature ratio is less than unity in the regions closest to the black hole, it rises above unity by $20 r_g$ along the jet.

In contrast, in the magnetic reconnection heated model $R17$, $T_e/T_i < 1$ everywhere. In the jet around $30 r_g$ from the black hole, $T_e/T_i \approx 0.3$, and the ratio increases with radius, reaching 0.75 around $1000 r_g$. In the disc, while the value of $\delta_e$ is higher than in the turbulent cascade heating simulation $H10$, the disc temperature ratio is not substantially different, taking an average value of $\sim 0.1$ around $30 r_g$.

This was not the case in the Sgr A* simulations presented in Chael et al. (2018a), where the turbulent cascade heated simulations had lower electron temperatures in the disc than the simulations heated by magnetic reconnection. The similarity of the outer disc electron temperatures in the present models may arise from the increased importance of Coulomb coupling in the denser regions of these higher accretion rate simulations (Ryan et al. 2018).

Fig. 3 shows that, as in the simulations of Sgr A* (Chael et al. 2017), the choice of electron heating prescription has a

4 RESULTS

4.1 Accretion Flow Properties

In Figs. 3 and 4, we show quantities averaged in azimuth and time over the time period $t = 10,000 - 15,000 r_g$ after rescaling the density, internal energy, and magnetic field to match the 230 GHz flux density measured by the EHT.

Fig. 3 shows properties related to the thermodynamics of the accretion flow: the electron heating fraction $\delta_e$, the gas temperature $T_{gas}$, the electron temperature $T_e$, the ion temperature $T_i$, and the temperature ratio $T_e/T_i$. Fig. 4 displays the mass density $\rho$, the bulk Lorentz factor $\gamma = u^{\nu}/\sqrt{-g^{00}}$, the magnetization $\sigma_i$, the ratio of ion thermal pressure to magnetic pressure $\beta_i$, and the ratio of radiation pressure to gas pressure in the fluid frame $\beta_R = E/3p$.

In each profile in Figs. 3 and 4, the solid white contour shows the $\sigma_i = 1$ surface, while the dotted black contour shows the surface where the Bernoulli number $Be = 0.05$. Expressing $T_e$ in Boyer-Lindquist coordinates, the Bernoulli number is (Narayan et al. 2012; Sadowski et al. 2013b)

$$Be = \frac{T_e^4 + R_i^4}{\rho u^2} - 1. \quad (13)$$

For a cold unmagnetized fluid, $Be = 0.05$ corresponds to a flow velocity of $\approx 0.3c$ at infinity.

From the leftmost panels of Fig. 3, the different asymptotic behaviors of the two heating prescriptions introduced in Sec. 2.2 are evident in the average values of $\delta_e$ in the jet region. In model $H10$, $\delta_e \approx 1$ everywhere inside the jet region defined by the $Be = 0.05$ contour. Outside the highly magnetized funnel, $\delta_e$ drops to nearly zero in the outer regions of the disc, $r \gtrsim 40 r_g$. In contrast, in model $R17$, $\delta_e$ reaches its limit of equipartition of thermal energy ($\delta_e = 0.5$) inside the $\sigma_i = 1$ contour. In the less magnetized disc outside $r \gtrsim 15 r_g$, $\delta_e$ falls to a small but nonzero value $\delta_e \approx 0.2$.
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Figure 3. Time- and azimuth-averaged thermodynamic quantities of the two simulations over the period $t = 10,000 - 15,000 t_g$. The top row shows quantities for the model H10 heated by the turbulent cascade prescription, and the bottom row shows quantities for the model R17 heated by magnetic reconnection. Snapshot quantities were averaged in azimuth and then time-averaged for $5,000 t_g$. The resulting averages were symmetrized over the equatorial plane. From left to right, the quantities shown are the electron heating fraction $\delta_e$, the combined gas temperature $T_{\text{gas}}$ in K, the electron temperature $T_e$, the ion temperature $T_i$, and the electron-to-ion temperature ratio $T_e/T_i$. The solid white contour in each panel denotes the surface where $\sigma_i = 1$, and the dashed black contour shows the surface where the Bernoulli parameter (Eq. 13) $B_e = 0.05$, which we take as the definition of the jet-disc boundary.

Figure 4. Additional time- and azimuth-averaged properties of the two simulations. From left to right, the quantities displayed are the density $\rho$ in g cm$^{-3}$, the bulk Lorentz factor $\gamma$, the plasma magnetization $\sigma_i$, the ratio of ion thermal pressure to magnetic pressure $\beta_i$, and the ratio of radiation pressure to thermal pressure $\beta_R$. In the first column, white contours show the poloidal magnetic field in the averaged data. In the remaining columns, the solid white contour denotes the $\sigma_i = 1$ surface. The dashed black contour shows the $B_e = 0.05$ surface defining the jet boundary. The dashed white contour in the third panel shows the $\sigma_i = 25$ surface; this is the maximum $\sigma_i$ included in the radiative transfer (see Section 3.3).

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Table 1. Time-averaged quantities for both simulations. From left to right, the quantities presented are the model name, the spin and heating prescription used, the average mass accretion rate through the black hole horizon measured in Eddington units, the magnetic flux threading the horizon measured in natural units, the mechanical jet power $P_J$ measured at a radius of $100 \, r_g$ (Eq. 14), the corresponding jet efficiency $\epsilon_J = \langle P_J \rangle / \langle M \rangle c^2$, the jet power including radiation $P_{J,rad}$ (Eq. 15), and the corresponding efficiency, both measured at the same radius.

| Model | Spin | Heating            | $\langle M/M_{\text{Edd}} \rangle$ | $\langle \Phi_{\text{BH}}/(Mc)^{1/2} \rangle$ | $\langle P_J(100) \rangle$ | $\epsilon_J$ | $\langle P_{J,rad}(100) \rangle$ | $\epsilon_{J,rad}$ |
|-------|------|--------------------|-------------------------------------|---------------------------------------------|-----------------|--------------|-----------------------------|-----------------|
| H10   | 0.9375 | Turb. Cascade     | $3.5 \times 10^{-6}$                | 54                                          | $6.6 \times 10^{43}$  | 0.5          | $1.2 \times 10^{43}$       | 0.9             |
| R17   | 0.9375 | Mag. Reconnection | $2.3 \times 10^{-6}$                | 63                                          | $1.4 \times 10^{43}$  | 1.4          | $1.4 \times 10^{43}$       | 1.6             |

noticeable effect on the electron-ion thermodynamics of the system. Unlike in the low-magnetic-flux simulations considered in that work, however, the two MAD simulations considered here have notably different gas and radiation kinematics as well, arising from the choice of heating prescription, even though both models produce thick ($H/R \approx 0.41$ at $10 r_g$), highly magnetized discs.

Table 1 summarizes important time-averaged quantities from our simulations, and Fig. 5 shows the accretion rate, magnetic flux through the horizon, and 230 GHz flux density as a function of time. Both simulations have a low accretion rate $\sim 10^{-6} M_{\text{Edd}}$, and both simulations reach the MAD state, with the magnetic flux threading the black hole $> 50 \sqrt{M_{\text{Edd}}}$, though R17 is slightly more magnetized. In both simulations, while the averaged value of the accretion rate is stable, there are larger excursions with time than in the low-magnetic-flux simulations of Chael et al. (2018a) due to the interaction of the accretion flow and the magnetic flux. This is particularly apparent in R17, which is more magnetized. Furthermore, R17 shows a slight but steady increase both in the magnetization and 230 GHz flux density over the entire 5000 $t_g$ period. It is possible this indicates this simulation, unlike H10, requires more time to completely restore equilibrium after the density rescaling at $t = 10^4 t_g$.

The high jet electron temperatures arising from the turbulent cascade heated model in simulation H10 produce a much more intense radiation field from synchrotron and inverse Compton scattering in the jet region than in R17. This is easily seen in the last column of Fig. 4, which shows the ratio of the radiation pressure to gas pressure as measured in KGRAL. While this quantity is $< 1$ almost everywhere in the R17 simulation, it approaches values $\gtrsim 100$ in the jet in model H10.

The conversion of much of model H10’s energy and momentum to radiation in the inner jet has a significant impact on its mechanical properties. While both simulations...
launch relativistic jets. R10’s jet is weaker, with Lorentz factor $\gamma \approx 1.5$ at $50 \, r_g$ compared to $\gamma \approx 2$ for R17 (Fig. 2, second column). At larger radii $r \sim 1000 \, r_g$, R10 reaches a Lorentz factor of $\approx 3$, a factor of two less than $\gamma \approx 6$ for R17. The conversion of fluid and magnetic energy into radiation also likely accounts for H10’s smaller value of horizon flux ($\Phi_{\text{BHY}}$) and correspondingly larger mass accretion rate $\dot{M}$ (Table 1).

We measure the thermal, magnetic, and jet mechanical power in both simulations as a function of radius using the definition (Tchekhovskoy et al. 2011; Ryan et al. 2018)

$$P_T = -\int \left( T'_r + \rho u'_r \right) \sqrt{-g} \, d\tau d\phi,$$

where the integral at a fixed $r$ is over the jet cap, which we define by the criterion $B_e > 0.05$ (Narayan et al. 2012; Sadowski et al. 2013b). The time-averaged jet power measured by Eq. 14 is roughly constant with radius from around $r = 10 \, r_g$ out to $r = 1000 \, r_g$. We measure the average jet powers at 100 $r_g$ from the averaged data to be $6.6 \times 10^{43}$ erg s$^{-1}$ for model H10 and $1.2 \times 10^{43}$ erg s$^{-1}$ for R17 (Table 1).

While the jet powers obtained from the two simulations agree to within a factor of two, the value obtained for model R17 is more consistent with the measured values for M87 of $\sim 10^{43}$–$10^{44}$ erg s$^{-1}$ (Reynolds et al. 1996; Stawarz et al. 2006). Comparing the jet power to the accretion rate gives a jet efficiency $\eta_J = P_T / \dot{M} c^2$ of 1.4 and 0.5 for R17 and H10, respectively, indicating that spin energy is being extracted from the black hole. This is especially true in model R17, which has greater than 100 per cent efficiency (Tchekhovskoy et al. 2011).

Because much of H10’s energy and momentum is converted to radiation in the jet, it has a correspondingly lower mechanical jet power. Including radiation in the jet power measurement, we define

$$P_{J, \text{rad}} = -\int \left( T'_r + R'_r - \rho u'_r \right) \sqrt{-g} \, d\tau d\phi,$$

This increases the measured jet powers to $P_{J, \text{rad}} = 1.2 \times 10^{43}$ erg s$^{-1}$ for H10 and $P_{J, \text{rad}} = 1.4 \times 10^{43}$ erg s$^{-1}$ for R17, and increases the jet efficiencies in the two models to 0.9 and 1.6, respectively.

Much of the intense radiation in H10 is produced from regions near the $\sigma_i, \text{max} = 100$ ceiling. Because mass is constantly being injected in these regions, energy and momentum is added to the simulation and is then efficiently converted into radiation. As discussed in Section 3.3, we do not trust the radiation produced in this region, and we do not include these regions in our post-processing computation of spectra and images in the following sections. In the (optically thin) GRRMHD simulation itself, because the frequency-averaged radiation field produced in these regions spreads through the simulation volume at nearly the speed of light, it is difficult to extract meaningful radiation quantities that are unaffected by the density floors. For these reasons, when interpreting the jet power and other quantities from model H10, it is important to note that the results may be strongly dependent on the specific choice of density floors. Because the jet electrons stay at lower temperatures in model R17, the jet radiation is less intense and model R17 is less prone to these uncertainties (see Sec. 4.5).

### 4.2 Spectra

Fig. 6 shows the spectral energy distributions (SEDs) for both of our simulations. These were obtained from post-processing with HEROIC over the full 5000 $r_g$ duration of the simulation after rescaling the density to approximately match the compact 230 GHz flux measured by Doeleman et al. (2012) and Akiyama et al. (2015). Synchrotron, free-free, and inverse Compton emission are all included in the HEROIC calculations, although we found that free-free emission does not contribute significantly even in the X-ray. These spectra do not include radiation produced from regions with $\sigma_i > 25$, and consequently the total luminosity from the post-processing spectra is less than that produced in the simulation bolometric $L_{\nu}$. In computing spectra, HEROIC used data from the simulation out to 1000 $r_g$, so diffuse emission from a good fraction of the jet is included, but the outermost regions of the jet are ignored.

Comparing simulated SEDs with observations of M87 requires some care. Specifically, because the total radio flux density along the extended jet of M87 is comparable to that of the significantly brighter but extremely compact region near the black hole (the “core”), a meaningful comparison with observations requires excising jet contributions that are outside our simulated domain. Prieto et al. (2016) have compiled measurements (their Table 1) from radio to X-ray of the M87 spectrum in a quiescent state, using only measurements that achieve at least 0.4″ resolution (to securely exclude emission from the brightest jet knot, HST-1). They also compiled measurements of the total flux density of the most compact component identified by VLBI observations (their Table 4). However, these latter measurements have some notable limitations. For example, at 86 GHz, Prieto et al. (2016) include the value $S_1 = 0.16 \pm 0.07$ Jy reported by Lee et al. (2008), corresponding to the measured flux density on the longest baseline for observations with the Coordinated (global) Millimeter VLBI Array (CMVA/GMVA). This approach is problematic because the core is resolved at 86 GHz (Kim et al. 2018) and because interference among compact components can significantly affect the correlated flux density on a single baseline. At 22 GHz, Prieto et al. (2016) include the value 0.35 Jy, reported by Junior & Biretta (1995). However, during the epochs reported by Junior & Biretta (1995), the total flux density of M87 was only $\sim 1.1$ Jy, which is significantly lower than the values measured more recently with the Very Long Baseline Array (VLBA) and the KVN/VERA Array (KaVA) (Hada et al. 2017). It is also lower than the values measured at 15 GHz and 43 GHz since 2000 (Lister et al. 2018; Walker et al. 2018).

Because we normalize our results to have a total flux density at 230 GHz that matches measurements taken in 2009 and 2012, we provide updated estimates of the total and compact flux density of M87 in Table 2. To estimate the flux density of the compact component, we measure the peak flux density of a beam convolved image from VLBI observations at a given frequency. This procedure gives a direct comparison between our simulated images and reconstructed images from VLBI. In contrast, the total flux density measured with VLBI may still contain significant contributions from outside our simulation domain, especially at centimeter wavelengths, and VLBI measurements may resolve out...
Figure 6. SEDs for the two models calculated with HEROIC for an observer at 17° inclination (Walker et al. 2018) measured up from the simulation south pole. Spectra were computed from 3D simulation snapshots every 10\,t\,g over the 5000\,t\,g period from \( t = 10^4 \)\,t\,g to \( t = 1.5 \times 10^4 \)\,t\,g after rescaling the density to approximately match the 0.98 Jy flux density of compact emission at 230\,GHz (Doeleman et al. 2012; Akiyama et al. 2015). The solid curve shows the median spectrum for each model, and the shaded region shows the nominal 1\,\sigma time-variability if we assume that the variability distribution is Gaussian at each frequency. Data points are taken from Table 1 of Prieto et al. (2016) (black) and Table 2 of this work. Measurements of the total flux density at radio wavelengths from Table 2 are displayed in cyan, while measurements of the compact flux density of the core are displayed in magenta.

Table 2. The total and compact radio spectrum of M87 from recent VLBI observations.

| Frequency (GHz) | Total Flux Density [Jy] | Core Flux Density [Jy] | Data Used |
|----------------|-------------------------|------------------------|-----------|
| 15.4           | 2.2 ± 0.3               | 1.3 ± 0.1              | 19 MOJAVE observations from 2001-2011 (Lister et al. 2018). |
| 22             | 2.1 ± 0.1               | 1.2 ± 0.1              | 10 KaVA & 3 VLBA (24 GHz) observations from 2013-2014 (Hada et al. 2017). |
| 43.1           | 1.6 ± 0.4               | 0.7 ± 0.2              | 50 VLBA observations from 1999-2016 (Table 3 in Walker et al. 2018). |
| 86.3           | 1.1 ± 0.5               | 0.8 ± 0.4              | 5 GMVA observations analyzed by Kim et al. (2018). |
| 230.0          | 2.05 ± 0.15             | 0.98 ± 0.05            | 2 EHT observations: Doeleman et al. (2012) & Akiyama et al. (2015). |

emission on yet larger scales, which can be estimated using connected-element interferometers.

SEDs from both our models are largely consistent with the radio spectrum data up to the synchrotron peak at 230\,GHz; this flat spectrum is also produced by analytic models of relativistic jets (Blandford & Königl 1979; Falcke & Biermann 1995). The model spectra underpredict the total measured flux at frequencies < 10^{10} Hz; at these low frequencies, the jet on scales larger than 1000\,r_s \approx 3600\,µas likely makes substantial contributions to the total emission.

Neither SED matches the measured flux at infrared through ultraviolet frequencies, although the hot jet electrons in the H10 model do extend the thermal synchrotron spectrum to the \( \approx 3 \times 10^{13} \) Hz measurements by Perlman et al. (2001) and Whysong & Antongucci (2004). The excess emission from the near infrared to ultraviolet may be explained by the addition of a high-energy non-thermal electron population in the disc (Broderick & Loeb 2009) or the jet (Dexter et al. 2012; Prieto et al. 2016).

The hotter electrons in the H10 simulation produce more inverse Compton power at X-ray frequencies. Both models produce flat SEDs from Comptonization at X-ray frequencies, which roughly match the slope of the Chandra measurements of the core of M87, but underpredict the total flux density in the X-ray. However, the shape of the spectrum at frequencies > 10^{12} Hz depends strongly on the choice of
extract the variability time-scale (see Dexter et al. 2014 for details). For the SMA observations of M87, they measured a damped random walk time-scale $\tau = 45^{+64}_{-24}$ days = $127^{+173}_{-60} t_g$.

Here, since we have regularly sampled, noise-free lightcurves from our simulations over $\sim$ 5 years, to measure the correlation time-scale in our 230 GHz lightcurves we simply compute the structure function and perform a least squares fit to obtain $S_\infty^2$ and $\tau$. We compute the structure function for time-scales $\Delta t$ ranging from our sampling cadence $10t_g = 3.5$ days up to 500 days. We first remove a linear slope from both lightcurves before computing the structure functions. For R17 this slope is negligible, but as discussed in Sec. 4.1, simulation H10 shows a slight but steady increase of 230 GHz flux density over the interval considered; the slope we subtract before performing the structure function analysis is $\approx 0.5$ Jy/5000 $t_g$.

Our structure functions and the best-fitting damped random walk models (Eq. 17) are presented in Fig. 7. The structure functions of both models show the characteristic form of a damped random walk, but with quite different measured time-scales. For H10, we measure a correlation time-scale $\tau \approx 43$ days, while for R17 we measure $\tau \approx 130$ days. The difference in variability time-scales may point to distinct origins for the emission in the two models; as we show in Section 4.4.3, the 230 GHz flux density in R17 has a larger contribution from the disc. The shorter variability time scale in H10 may also point to faster dynamics in the inner accretion flow from the strong radiative feedback (Section 4.1).

We find that model H10 has a variability time-scale consistent the measurement of Bower et al. (2015), while R17 varies on longer time-scales than observed. However, these results should be viewed with some caution as our lightcurves were produced with the “fast light” approximation, ignoring the evolution of the fluid as photons propagate through the accretion flow and jet. Properly accounting for the fluid evolution during the radiative transfer may significantly change the characteristics of the produced variability.

Furthermore, the SMA observations presented in Bower et al. (2015) do not resolve the core component directly and include an extra $\sim 1$ Jy of flux from the outer jet at 230 GHz which is not captured in our simulations.

### 4.3 Variability

The 230 GHz lightcurves from our high-resolution grtrans images are presented in the bottom panel of Fig. 5. Both simulations show variability on time-scales of days to years, and the normalized root-mean-square variability of both lightcurves is $\approx 15$ per cent relative to the mean.

Bower et al. (2015) modeled the 230 GHz light curve of M87 observed over more than 10 years with the Submillimeter Array (SMA) as a damped random walk process. They define the structure function $S^2(\Delta t)$ of the lightcurve $s(t)$ as

$$S^2(\Delta t) = \frac{1}{N} \sum (s(t) - s(t+\Delta t))^2,$$

where $N$ is the number of data points in the sum. For a damped random walk, this structure function takes the form

$$S^2_{\text{DRW}}(\Delta t) = S^2_\infty \left(1 - e^{-\Delta t/\tau} \right),$$

where $S^2_\infty$ is the power on long time-scales and $\tau$ is the correlation or damping time of the random walk.

For real data that is noisy and irregularly spaced, the structure function defined in Eq. 16 can have spurious artefacts. Therefore, Bower et al. (2015) did not measure the structure function directly but instead used a Bayesian approach to model the irregularly sampled lightcurve and measure a damped random walk time-scale $\tau$. However, since we have regularly sampled, noise-free lightcurves from our simulations over $\sim$ 5 years, we can directly measure the correlation time-scale in our lightcurves.

### 4.4 Images

To compare our models against existing images and data from VLBI observations, we computed high resolution images at 15, 22, 43, 86, 230, and 345 GHz using grtrans. We only include emission from the jet out to 3000 $r_g$. With the small inclination angle of $17^\circ$ (Mertens et al. 2016; Walker et al. 2018), this corresponds to a maximum projected jet length in our images of 850 $r_g$, or $\approx 3$ mas. In contrast, jet emission at 43 GHz extends out to at least 20 mas (Walker et al. 2018). We chose a representative snapshot for each model where the core flux density at 230 GHz was close to the measured value of 0.98 Jy from the EHT (Doeleman et al. 2012; Akiyama et al. 2015).

In Fig. 8, we show log-scale images of both models at 43, 86, and 230 GHz, each with a dynamic range of $10^4$. The jet structure is similar in both simulations, with a wide apparent opening angle that increases to $> 90^\circ$ at the jet base in
Figure 8. Log scale images of simulation snapshots of the two models at 43 GHz (left), 86 GHz (middle) and 230 GHz (right). Snapshots were observed at an inclination angle of 17° up from the simulation south pole and rotated 108° counterclockwise to match the observed jet orientation. The intensity scale is different at each frequency, but for each frequency the scale displays a dynamic range of $10^4$ and is the same for both the image from the H10 simulation (top) and R17 simulation (bottom). The image length scale changes with frequency; dotted boxes on the 43 and 86 GHz images show the fields-of-view of the 86 and 230 GHz images, respectively. The jet structure is qualitatively similar in the two simulations, with wide apparent opening angles which narrow with distance from the SMBH at lower frequencies. Images at all frequencies show a faint counterjet, and the black hole shadow is evident in both models even down to 43 GHz.

4.4.1 Core-Shift

VLBI observations with absolute phase referencing allow estimates of the relative image location at different frequencies. Because the M87 jet is optically thick at wavelengths longer than a few millimeters, the image centroid of the bright, compact core emission moves with frequency, giving rise to the so-called “core-shift” effect (Blandford & Königl 1979).

Hada et al. (2011) conducted measurements of the core-shift of M87 at frequencies from 2.3 to 43 GHz, finding that the millimeter core is coincident with the SMBH and disc that launch the jet. They estimated that the radio core has a right ascension displacement (relative to the 43 GHz core) given by $\Delta \alpha \approx A \lambda^\alpha + B$, where $A = (1.40 \pm 0.16) \text{ mas}$, $B = (−0.041 \pm 0.012) \text{ mas}$, and $\alpha = 0.94 \pm 0.09$. We computed the analogous core-shifts for our simulated images by first convolving the images with the wavelength-dependent observing beam of Hada et al. (2011) and then measuring the location of the peak in the resulting image.

Fig. 9 shows the results of this analysis for our simulations. Both H10 and R17 produce images with core-shifts that are compatible with the results of Hada et al. (2011) at frequencies as low as 15 GHz. Even though VLBI constrains the core-shift at yet lower frequencies, we did not attempt to estimate core-shifts at these frequencies from our simulations because the image sizes and observing beams are comparable to the raytracing domain.

Radio-jet core-shifts can be used to measure the jet magnetic field. Recently, Zamaninasab et al. (2014) and Zdziarski et al. (2015) used core-shifts to measure the jet magnetic fields of several LLAGN sources (including M87) on $\sim \text{pc}$ scales; they found their values of magnetic field and jet powers were consistent with jets launched by MADs. This is consistent with the findings of this work for M87.
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Section 4.4.2 43 and 86 GHz Images

VLA images show that the jet of M87 is wide, with an apparent opening angle that decreases with radius. The apparent opening angle is $\approx 55^\circ$ at 43 GHz (Walker et al. 2018), increasing up to $\approx 127^\circ$ in the innermost regions ($\approx 50 r_g$) of 86 GHz images (Kim et al. 2018). Fig. 10 compares the simulation images of the inner 2 mas at 43 GHz with an image of 2007 VLBA data Walker et al. (2018) reconstructed with a new closure phase and amplitude imaging method implemented in the eht-imaging software library (Fig. 9 of Chael et al. 2018b). The top row shows the snapshot images for both simulations, and the bottom row shows the VLBI image and the snapshot images convolved with a beam that has a size half of the nominal value reported in Walker et al. (2018). Lines indicating the $55^\circ$ apparent opening angle measured by ridge line analysis in Walker et al. (2018) are overlaid on all images. While previous simulations of M87 using weakly magnetized discs produced narrow jets (Mościbrodzka et al. 2016; Ryan et al. 2018), the observed wide apparent opening angle is naturally produced in these MAD simulations. The simulation images blurred to the same resolution as the VLBI image also show limb brightening in the jet, though the contrast between intensity on the jet edges and along the axis is in general less prominent than in the VLBI image. The counterjet is faint but visible at the edge of our dynamic range in both model images, but it is more prominent in the VLBI image at this epoch.

At smaller scales, the magnetic field strength can also be estimated using the measured core sizes from VLBI. Kino et al. (2014) used this method with a model of the optically thick 43 GHz synchrotron emission to estimate a field strength $1 \lesssim |B| \lesssim 15$ G at an angular scale of 100 $\mu$as. Assuming an inclination angle of $17^\circ$, an angular scale of 100 $\mu$as corresponds to $r \approx 100 r_g$. From the time- and azimuth-averaged simulation data, we find for both models $1 \lesssim |B|_{100 r_g} \lesssim 2$ G, within the measured range. Closer to the black hole, Kino et al. (2015) used EHT data to estimate a field strength $58 \lesssim |B| \lesssim 127$ G on scales $\sim 10$ $\mu$as, corresponding to a de-projected distance of $\approx 10 r_g$. From our averaged simulation data, we estimate $|B|_{10 r_g} \approx 20$ $\mu$G at this radius in the jet. However, Kino et al. (2015) obtain their estimate by assuming a spherical 230 GHz emission region that is optically thick to synchrotron self-absorption, which we do not observe at 230 GHz in our simulations.

4.4.3 230 GHz Images

At 1.3 mm, The EHT has already constrained the emission size at the core to be on the order of $\approx 40$ $\mu$as (Doelmann et al. 2012; Akiyama et al. 2015), approximately the size of the lensed black hole shadow for $M \approx 6.2 \times 10^9 M_\odot$ (Gebhardt et al. 1997) and $D = 16.7$ Mpc (Mei et al. 2007). More recent EHT observations, with the full array, are expected to directly image the jet launching region and potentially the black hole shadow itself (Lu et al. 2014; Akiyama et al. 2017).

Fig. 12 compares three linear scale images at 230 GHz showing time evolution of the source over roughly 1 week, the usual length of observation campaigns by the EHT. In each set, the top row shows high-resolution images from grtrans, and the bottom row shows the same images blurred with a Gaussian beam with a 15 $\mu$as full width at half maximum (FWHM). This beam size is approximately the imaging resolution of the EHT (Chael et al. 2018b).

The 230 GHz images for both models show a distinct black hole shadow that is large enough to be imaged by the EHT. The diameter of the shadow is approximately $9.7 r_g \approx 36$ $\mu$as (slightly less than the diameter for a Schwarzschild black hole, $2\sqrt{27} r_g \approx 38$ $\mu$as), given the assumed mass and distance. Because we do not view the accretion flow directly face-on, the bottom half of the ring structure is brighter due to Doppler boosting of the jet and disc emission. Based on the differential limb brightening and velocities in the large-scale jet (Walker et al. 2018), we chose the jet to rotate clockwise in the plane of the sky. Given the sense of rotation determined by Walker et al. (2018) and the direction of the projected jet on the sky, the 230 GHz ring is expected to be brighter on the bottom by simple geometry.

Bright ridges tracing the rotating helical magnetic field are visible in both simulations. These extended structures are fainter than the bright ring, but they are visible even in the images blurred to the EHT’s resolution. The brightest spot in the 230 GHz image moves with the rotation of the magnetic field lines, particularly in the H10 model, which
lacks bright disc emission. When blurred to EHT resolution, this evolution will generally follow the clockwise sense of rotation of the disc and jet. In the specific frames selected from the H10 simulation however, shifts in the relative brightness of two filaments as they rotate produce an apparent evolution in the blurred frames that is slightly counterclockwise.

The precise shadow diameter and shape is sensitive to the inclination and black hole spin (Bardeen et al. 1972; Chandrasekhar 1983). Even for these simulated images, which have a prominent shadow, any one EHT image reconstruction would leave substantial uncertainty in the shadow size due to the limited resolution and contributions to the source structure from the foreground jet. It may be possible, however, to make a more precise measurement of the shadow size with multi-epoch imaging. While the shadow is a persistent feature set only by the mass and spin of the black hole, the foreground jet at 230 GHz rotates quickly, completing a full revolution on a time-scale of weeks to months.

In Fig. 13, we compare visibility amplitudes extracted from the fiducial 230 GHz images in Fig 12 with observations from the EHT with stations in Hawaii, California, and Arizona in 2009 (Doeleman et al. 2012) and 2012 (Akiyama et al. 2015). The compact flux density measured in these two years was \( \approx 0.98 \) Jy. At the chosen values of inclination \( \theta = 17^\circ \) up from the south pole and distance to the black hole \( D = 16.7 \) Mpc, the snapshot images from our simulations are somewhat too large and underpredict the measured visibility amplitudes on the Hawaii-California and Hawaii-Arizona baselines. Compared to the extremely compact images obtained from the less magnetized simulations in Ryan et al. (2018), our wide-opening-angle jets produce extended emission that increases the overall image size, though the bright \( \sim 40\mu\text{as} \) photon ring remains the most prominent feature.

We note that the image size in our simulations is highly sensitive to the assumed viewing inclination. At larger values of inclination angle than the \( \theta = 17^\circ \) taken from Walker et al. (2018), the image size decreases as we see less Doppler-boosted emission from the wide opening angle jet. While the jet inclination angle is constrained to \( \lesssim 20^\circ \) at distances \( \sim 100\text{pc} \) from the black hole from apparent superluminal velocities measured near the HST-1 knot (Giroletti et al. 2012), it is not as definitively constrained on scales closer to the black hole. In their conservative estimate, Mertens et al. (2016) give an upper limit \( \theta = 27^\circ \). At \( \theta = 30^\circ \), the visibilities from R17 nearly match the observations, and the simulated visibilities from R10 are much less discrepant than at \( 17^\circ \) (Fig. 13). Furthermore, at this larger value of the inclination angle, the limb-brightening in the 43 and 86 GHz images from both simulations more closely matches the VLBI maps (left panels of Figs. 10 and 11), and the counterjet is more prominent at both frequencies.
The image size in our simulations at a fixed inclination angle is also sensitive to the $\sigma_i$ cut that we impose in post-processing. Our choice of $\sigma_i = 25$ was determined to avoid including emission from the regions with densities set to the floor value. In nature, these regions will produce emission which is likely to be significant in the regions probed by the 230 GHz image. We find that including radiation from $\sigma_i > 25$ regions makes the image more compact (see Section 4.5). Thus, it remains possible that different prescriptions for including radiation in post-processing from highly magnetized regions may produce images from MAD simulations that are consistent with the EHT size measurements.

In Fig. 14, we consider the same 230 GHz snapshots as in the central column of Fig. 12 and decompose the emission into three component parts: disc, counterjet, and forward jet. As above we define regions with $\text{Be} > 0.05$ to be in the jets, and $\text{Be} \leq 0.05$ to be in the disc. We produce these component images by zeroing out the emissivities outside the selected regions when doing the radiative transfer in \textsc{grtrans}. Throughout, we maintain the global cut on emissivities from regions with $\sigma_i > 25$.

Because the entire M87 accretion flow is optically thin at 230 GHz, the total flux densities of the component images in Fig. 14 nearly add up to the total flux density in the image generated from the entire emissivity distribution. In both simulations, the majority ($\sim 60$ per cent) of the emission comes from the forward jet. The forward jet emission in both cases consists of a spiral structure surrounding the black hole where emission traces magnetic field lines, as well as a persistent component from the photon ring. We note that a substantial fraction of the jet emission comes from between the $\text{Be} = 0.05$ and $\sigma_i = 1$ surfaces, so adopting $\sigma_i > 1$ as the definition of the jet region will result in assigning much more of what we classify as forward jet and counterjet emission to the disc.

In the H10 model, nearly all of the emission not from the forward jet is produced from the counterjet, which adds to the prominent photon ring. Because the electrons in this model are so much hotter at the base of the jet than in the disc, the emission from these regions dominates the total and the disc emission is negligible. In contrast, the reconnection heating model R17 has approximately equal contributions from the counterjet and disc. The disc emission shows a persistent bright spot from the Doppler-boosted accretion flow. Unlike the bright spots produced from rotating magnetic field lines, this spot remains constant in position and does not rotate around the photon ring with time. This points to the possibility of using multi-epoch imaging with the EHT to disentangle the source structure and identify whether or not the accretion disc emission makes a substantial contribution to the image flux.
4.5 The effects of our choice of $\sigma_{\text{cut}}$

In this section, we explore the effects of our choice of choosing $\sigma_{\text{cut}} = 25$ in the results presented in the previous sections. For our chosen simulation snapshots, we re-generate spectra and 230 GHz images testing five different values of $\sigma_{\text{cut}}$: $\sigma_{\text{cut}} = 1, 10, 25$ (our fiducial value), 50, and no ceiling.

From Fig. 15, we see that in both simulations, any choice of $\sigma_{\text{cut}}$ above unity has relatively little effect on the radio spectrum up to 230 GHz. Most emission in this part of the spectrum comes from less-magnetized regions farther from the black hole. As a result, the predictions of both models at frequencies $< 230$ GHz should be relatively insensitive.

Figure 12. Linear scale images of sequential snapshots at 230 GHz from two models showing time evolution over $20 t_g \approx 7$ days, the usual length of an EHT observing campaign. The top two rows show images from simulation H10, and the bottom two rows show images simulation R17. In each set of images, the top row shows high-resolution images from grtrans, and the bottom row shows the same images blurred with a circular Gaussian beam with a 15 $\mu$as FWHM, approximately representing the imaging resolution of the EHT.
The two MAD simulations presented in this paper produce spectra and images that are broadly consistent with observations of the M87 jet at centimetre and millimetre wavelengths. Model R17 produces a jet power that is in line with observations, while model H10 produces a jet power lower by a factor of two, unless radiation (made suspect by its origin in regions with active density and energy floors) is considered. Both models reproduce the wide jet opening angle observed at 43 GHz, and they are both consistent with core-shift observations. However, at a viewing angle of 17°, the 230 GHz images from both models are too large to match EHT observations from 2009 and 2012.

M87 has been simulated before using GRMHD and GRMHD codes, though not as frequently as Sgr A*. The model of Mościbrodzka et al. (2016) is a representative state-of-the-art combination of single-fluid GRMHD and radiative transfer. The authors used discs with relatively weak magnetization from Shiokawa (2013) and added thermal electrons in post-processing that were assumed to be hot in the jet and cold in the disc (Mościbrodzka et al. 2014). They produce a model with a jet power in the correct range, a flat radio spectrum, and an limb-brightened jet image at 43 GHz. The jets in their simulations show substantial variability and apparent superluminal motion from field lines along the fun-
nel wall. However, the jets in their simulation have apparent opening angles at 43 GHz smaller than the observed ~ 55°.

As in the MAD models in this work, 230 GHz images from Mościbrodzka et al. (2016) show spiral structures from helical field lines in the jet; this is a common prediction of both weakly magnetized and MAD models. At 230 GHz, their images are dominated by the counterjet (see also Dexter et al. 2012). Unlike in our MAD models, their 230 GHz images satisfy the constraints on the image size from the EHT at 20° inclination.

The authors have also investigated the polarized emission from their models (Mościbrodzka et al. 2017). They have shown that in their counterjet-dominated models it is possible to produce images with rotation measure and polarization fraction in line with observations through the depolarization of the counterjet emission as it passes through the cooler disc. The emission in both of our MAD models at 230 GHz is dominated by the forward jet, although counterjet emission is still significant. It is possible that, with less opportunity for depolarization through the disc, our forward-jet-dominated images might produce a net polarized flux that exceeds the observed value (~ 1 per cent; Kuo et al. 2014). This is an important direction for future work.

Recently, Ryan et al. (2018) performed the first two-temperature simulation of M87 using the code ebhlight in axisymmetry. Unlike KORAL, which considers only the frequency-integrated radiation field, ebhlight uses a Monte Carlo method where photons with distinct frequencies are emitted and absorbed on the simulation grid. Consequently, they obtain spectra as a natural product of their simulations without having to perform radiative transfer in post-processing.

Ryan et al. (2018) considered discs that were far less magnetized than those explored here. Consequently, to match the observed 230 GHz flux density they required higher accretion rates than in our simulations. In their best-fitting model, the accretion rate is \( \dot{M}/\dot{M}_{\text{Edd}} \sim 10^{-4} \), about 3× what we find in our lower-accretion-rate model H10. In all their models, they found that Coulomb heating of electrons becomes important in the outer disc. As in our simulations, they see that radiation plays a significant role in the inner disc. Notably, they explore simulations with both high (Gebhardt et al. 2011) and low (Walsh et al. 2013) black hole mass and consider two values of the black hole spin (\( a = 0.5 \) and \( a = 0.9375 \)). They find their high-spin, high mass model produces both a spectrum and 230 GHz image consistent with the available data – we have adopted these preferred parameter values in this study. Unlike in the present work, at \( M = 6.2 \times 10^8 M_\odot \) and \( a = 0.9375 \), they obtain a compact, counterjet-dominated 230 GHz image that is consistent with past EHT measurements of the overall image size. However, the jet powers produced in their simulations are several orders of magnitude too low, and their weakly magnetized discs also produce jet opening angles that are too small when compared against VLBI observations. These problems are not present in our MAD models, which match the observed spectral and image characteristics of the M87 jet well at all frequencies between 15 and 86 GHz.

Simulating the jet interior remains a problem in all sim-
Figure 15. Snapshot spectra from the two simulations generated with different values of $\sigma_{\text{cut}}$ in the radiative transfer. Spectra were generated with HEROIC zeroing out emissivities from all regions with fluid frame magnetization $\sigma_i \geq \sigma_{\text{cut}}$. We generate spectra with $\sigma_{\text{cut}} = 1, 10, 25$ (our fiducial value), $50$, and with no ceiling. In both simulations, any choice of $\sigma_{\text{cut}}$ above unity has little effect on the radio spectrum up to $230$ GHz. Most emission in this part of the spectrum comes from less-magnetized regions farther from the black hole. The choice of $\sigma_{\text{cut}}$ has a drastic effect on the spectrum at higher frequencies as direct synchrotron emission and Compton scattering in the most magnetized, high-temperature regions close to the black hole is added, increasing the radiative power. When no $\sigma_i$ ceiling is imposed, model H17 has an extreme total luminosity $> 10^{43}$ erg s$^{-1}$.

Figure 16. Snapshot images from the two simulations generated with different values of $\sigma_{\text{cut}}$ in the radiative transfer. From left to right, we present images generated using $\sigma_{\text{cut}} = 1, 10, 25$ (our fiducial value), $50$, and with no ceiling. In both simulations, the overall image structure is similar at all cuts up to $\sigma_{\text{cut}} = 50$. Because $\sigma_i$ increases rapidly with decreasing polar angle in the jet region (Fig. 2), including regions of higher and higher magnetization does not open up very different regions of the accretion flow to the radiative transfer. In contrast, including the entire interior of the jet (the rightmost images) produces substantial new emission in front of the photon ring.
magnetic reconnection heating model of Rowan et al. (2017); Chael et al. (2018a) used in simulation B17 launches a jet powered by the black hole spin with a mechanical jet power $\sim 10^{43}$ erg s$^{-1}$, in the correct range for M87 (Reynolds et al. 1996; Stawarz et al. 2006). In contrast, the jet heated by the turbulent cascade model of Howes (2010) in simulation B17 produces intense radiation at the jet base; this radiation sets the jet of mechanical energy, resulting in a mechanical jet power a factor of two less than in B17.

Despite the differences in their kinematics, the spectra and images of the two models are quite similar at centimetre, millimetre, and submillimetre wavelengths. The simulations did not run long enough to produce the full extent of the jet observed with VLBI, but the images from the inner milliarcsecond show similar structure to VLBA and GMVA images at the corresponding wavelengths. Notably, both models reproduce the measured $\approx 55^\circ$ opening angle (Walker et al. 2018), and both show emission from the counterjet, though our models produce less counterjet emission than observed in some 43 GHz VLBI images.

At 230 GHz, our simulations produce images that are somewhat larger than the size measured by the EHT in 2009 and 2012 (Doel et al. 2012; Akiyama et al. 2015). However, the image size is strongly dependent on the viewing inclination, and it becomes consistent with EHT observations around $\theta = 30^\circ$, the upper end of the most conservative range established by VLBI observations (Mertens et al. 2016). The image size also shrinks when the extremely magnetized on-axis regions we exclude in our radiative transfer are included. In this work, we do not trust the emission from these most magnetized regions of our models due to the imposed position of density floors; we exclude emission from regions with a magnetization greater than $\sigma_{\text{cut}} = 25$. A different treatment of these regions from a simulation with a higher overall ceiling on the magnetization could potentially bring the size of the 230 GHz images in line with observations.

Our 230 GHz images all show distinct black hole shadows, which are primarily illuminated by emission originating in the forward jet. Future investigations of polarized images from our simulations will investigate whether our forward-jet-dominated models satisfy constraints on the 230 GHz polarization fraction; these constraints are naturally satisfied by Faraday depolarization of counterjet emission (Moscibrodzka et al. 2017). Furthermore, the rotation of the accretion disc and jet makes our 230 GHz images dynamic on time-scales of weeks to months. It may be possible to distinguish between models by identifying jet, counterjet, and disc contributions to the 230 GHz image by tracking stationary and moving features through repeated observations with the EHT.

While our models under-produce the measured flux in the optical, ultraviolet, and X-ray, it seems clear that the spectrum at these high frequencies is dominated by the hottest, most magnetized regions closest to the black hole. A different treatment of the magnetization ceilings imposed in the simulation and radiative transfer may bring the simulation spectra from these regions more in line with observations. At the other end of the spectrum, we are unable to investigate the jet on large scales and at frequencies $<15$ GHz given the relatively short simulation time considered in this work.

Finally, we note that it is likely that nonthermal electrons in the disc (Broderick & Loeb 2009) or jet (Dexter et al. 2012) contribute to the emission at 230 GHz and into the infrared, optical, and ultraviolet. It is relatively straightforward to add non-thermal electron distributions to GRMHD simulations in postprocessing (Dexter et al. 2012; Ball et al. 2016; Davelaar et al. 2018). However, the present work, as well as previous work by Ressler et al. 2017,Ryan et al. (2018), and Chael et al. (2018a), has shown that including self-consistent evolution of electron populations in simulations substantially modifies the resulting images and spectra in ways that cannot be captured by simply assigning electron temperatures and distribution functions in post-processing. In this spirit, the logical next step is to self-consistently evolve non-thermal electron distributions in simulations of M87. In future work, we will use the method developed in Chael et al. (2017) to this end.

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