Formation of Episodic Jets and Associated Flares from Black Hole Accretion Systems

Miljenko Ćemeljic1,2, Hai Yang1,4,5, Feng Yuan1,4, and Hsien Shang3

1 Shanghai Astronomical Observatory, Chinese Academy of Sciences, 80 Nandan Road, Shanghai 200030, People’s Republic of China; fyuan@shao.ac.cn
2 Nicolaus Copernicus Astronomical Center, Polish Academy of Sciences, Bartycka 18, 00-716 Warsaw, Poland
3 Academia Sinica, Institute of Astronomy and Astrophysics, P.O. Box 23-141, Taipei 106, Taiwan
4 University of Chinese Academy of Sciences, 19A Yuquan Road, Beijing 100049, People’s Republic of China

Received 2021 December 21; revised 2022 April 23; accepted 2022 May 16; published 2022 July 4

Abstract

Episodic ejections of blobs (episodic jets) are widely observed in black hole sources and usually associated with flares. In this paper, by performing and analyzing three-dimensional general relativity magnetohydrodynamical numerical simulations of accretion flows, we investigate their physical mechanisms. We find that magnetic reconnection occurs in the accretion flow, likely due to the turbulent motion and differential rotation of the accretion flow, resulting in flares and formation of flux ropes. Flux ropes formed inside of 10–15 gravitational radii are found to mainly stay within the accretion flow, while flux ropes formed beyond this radius are ejected outward by magnetic forces and form the episodic jets. These results confirm the basic scenario proposed in Yuan et al. Moreover, our simulations find that the predicted velocity of the ejected blobs is in good consistency with observations of Sgr A*, M81, and M87. All of the processes were found to occur quasiperiodically, with the period being the orbital time at the radius where the flux rope is formed. The predicted period of the flares and ejections is consistent with those found from the light curves or image of Sgr A*, M87, and PKS 1510–089. The possible applications to protostellar accretion systems are discussed.

Unified Astronomy Thesaurus concepts: High energy astrophysics (739)

1. Introduction

Jets are ubiquitous in astrophysical accreting systems. Two types of jets have been observed, namely, continuous jets and episodic ones. In addition to the temporal difference, the radiation spectrum, polarization, and velocities of these two types of jet are also different, as summarized in Fender & Belloni (2004). Episodic jets and radiation flares are often found to be physically associated with each other. One such example is Sgr A*, the supermassive black hole in the Galactic center. Flares are observed in multiple wave bands such as radio, submillimeter, infrared, and X-ray (Baganoff et al. 2001; Genzel et al. 2003; Eckart et al. 2006b; Yusef-Zadeh et al. 2006; Tripp et al. 2007; Marrone et al. 2008). The flares are strongest in the IR and X-ray (Baganoff et al. 2001; Genzel et al. 2003) and occur simultaneously (Genzel et al. 2003), followed by flares at submillimeter and radio wave bands (Eckart et al. 2006b; Marrone et al. 2008). The peak of the light curve at 43 GHz is found to lead that at 22 GHz by 20–40 minutes (Yusef-Zadeh et al. 2006). This time lag is regarded as the evidence for the ejection and adiabatic expansion of plasmoid from the accretion flow (Yusef-Zadeh et al. 2006; Marrone et al. 2008). This interpretation is confirmed by the 7 mm Very Long Baseline Array (VLBA) observations that found a possible image of the ejected blob 4.5 hr after the IR flare, and the speed of the blob is estimated to be 0.4 speed of light (Rauch et al. 2016).

The phenomena of multiple wave band flares with a frequency-dependent time lag and associated ejection are common in black hole sources. For example, large radio flares were detected in three active galactic nuclei, namely, M81 (King et al. 2016), blazar PKS 1510–089 (Park et al. 2019), and M87 (Hada et al. 2014), preceded by an X-ray flare, γ-ray flare, and γ-ray flare in these three sources, respectively. Discrete blobs associated with flares were clearly detected by the VLBA or very long baseline interferometry (VLBI), with the speed of the blob measured to be 0.5 and 0.2 speed of light in the case of M81 (King et al. 2016) and M87 (Hada et al. 2014). A similar phenomenon occurs in a black hole X-ray binary, Cygnus X-1 (Wilms et al. 2007).

Flares and episodic ejection (as jet knots) have also been observed in protostellar accretion systems (Feigelson & Montmerle 1999; Wolk et al. 2005; Lee 2020). Although consensus has not been reached on the mechanisms of the ejection and flares, both of them are believed to be powered by magnetic activities in the accretion flow (Feigelson & Montmerle 1999), so it is natural to conjecture that they are physically associated with each other. More interestingly, even in the case of the Sun, large solar flares and coronal mass ejections (CMEs) also occur together (Lin & Forbes 2000; Zhang & Low 2005).

Many works have been published to explain the origin of black hole flares, usually by invoking synchrotron radiation or Compton scattering by relativistic electrons. These electrons are accelerated either by shocks in jets (Markoff et al. 2001) or by magnetic reconnection in accretion flows (Yuan et al. 2004; Dodds-Eden et al. 2009; Yuan et al. 2009; Dodds-Eden et al. 2010; Dexter et al. 2020; Nathanail et al. 2020; Ripperda et al. 2020; Zhao et al. 2020; Nathanail et al. 2021; Ripperda et al. 2021; Scepi et al. 2021). These works usually focus only on the flares, with little discussion of episodic ejections and their associations with flares.

An exception is Yuan et al. (2009). In this work, by analogy with the CMEs in the Sun, they proposed that the turbulent
motion and differential rotation of the black hole accretion flow will deform the magnetic field lines; consequently, magnetic reconnection occurs, and flux ropes are formed in the coronal region of the accretion flow, which will then be ejected by the magnetic compression force.

These processes must be highly dynamic, but assumptions and simplifications had to be adopted in the analytical treatment of Yuan et al. (2009). In the present work, we perform three-dimensional magnetohydrodynamical numerical simulations to examine the scenario proposed in Yuan et al. (2009), investigating the formation of episodic jets and their association with flares. The paper is structured as follows. In Section 2, we briefly describe our model of general relativity magnetohydrodynamical simulations of hot accretion flows. In Section 3, we present the main results obtained from our simulations, including the formation of the flux rope, its ejection, and the periodicity of the whole process. In Section 4, we discuss the applications of our model to observations, including the periodic flares detected in Sgr A* and other sources; the hot spot in Sgr A* detected by the GRAVITY observations; the velocity and possible periodicity of the ejected blobs in Sgr A*, M81, and M87; and finally, the periodic ejection and flares in protostellar accretion systems. A comparison between the present work and previous works is presented in Section 5. Finally, we summarize our work in Section 6.

2. Model

All black hole sources mentioned in the Introduction are powered by hot accretion flows (Ho 2008; Yuan & Narayan 2014). We therefore perform a set of three-dimensional general relativity magnetohydrodynamical numerical simulations of a hot accretion flow around a black hole using the ATHENA++ code. Spherical coordinates are adopted. We have simulated two kinds of accretion flows, namely, standard and normal evolution (SANE) and magnetically arrested disk (MAD), with the magnetization of the accretion flow in MAD being much stronger than that in SANE (Igumenshchev et al. 2003; Narayan et al. 2003; Tchekhovskoy et al. 2011). Two values of black hole spin are considered, $a = 0$ and 0.98. The details of the simulations are presented in the Appendix. We find that the value of $a$ is not important for our problem, so in the following, we simply refer to the models as SANE and MAD. Snapshots of the distributions of magnetic field lines and plasma $\beta$ (defined as the gas-to-magnetic pressure ratio) of the models are shown in Figure 1.

3. Results

3.1. Formation of Flux Ropes

Solar flares and CMEs have been intensively studied, and the physical mechanism is known to be magnetic reconnection (Zhang & Low 2005; Gou et al. 2019). After reconnection, some plasma will be enclosed within helical magnetic field lines and form a “flux rope.” The flux rope will then be ejected by a strong magnetic pressure force and form the CMEs. It was proposed in Yuan et al. (2009) that similar processes should also occur in the black hole accretion flows and result in the formation of episodic jets.

Stimulated by the above scenario, we have analyzed the simulation data of accretion flows, looking for flux ropes. We do find that many such helical structure magnetic field lines (i.e., the flux ropes) exist both within and at the coronal region of the accretion flow, from radii as small as $4r_g$ up to $\sim 30r_g$, where $r_g$ stands for gravitational radius. The flux ropes are found to keep forming throughout the evolution of the accretion flows in all four models once they reach the steady state. In the present paper, we have tried to find as many flux ropes as possible. We then measure the properties of individual flux ropes and present the typical results.

One example of a flux rope identified in the case of MAD at time $T = 20,000r_g/c$ is shown in Figure 2. The poloidal magnetic field structure is presented in the $x-z$ plane at nine various azimuthal angles. As we can see from the figure, in the two-dimensional plane, these flux ropes will be represented by circular magnetic field line patterns—“magnetic islands.” Note that if a constant-$\varphi$ slice through the flux rope were not perpendicular to the axis of the rope, we would see a whirl in the poloidal field lines instead of a circular magnetic island.

To understand the physical mechanism of the formation of flux ropes, after we find a flux rope, we have traced back to the time and location of the flux rope formation. The motion of the flux ropes is traced by examining their movement in different $r-\theta$ planes with various angles $\varphi$. The center of the flux rope is identified by the minimum of the poloidal magnetic field and

---

6. The time-averaged magnetic field configuration at a larger scale will show an “hourglass” shape, as shown by Figure 5 in Yang et al. (2021).
the circular pattern of the magnetic field lines, as we will explain at the end of Section 3.1. Figures 3 and 4 show the tracing results. Figure 3 shows the magnetic field structure in the $r-\theta$ plane at three snapshots. In each panel, we can see two reconnection layers, located above and below the equatorial plane of the accretion flow. These reconnection layers will continue to evolve, resulting in the occurrence of magnetic reconnection and formation of magnetic islands after a time of $\sim 100 r_g/c$. As an example, the evolution process of the reconnection layer located above the equatorial plane in the rightmost (i.e., $t = 19,400$) snapshot of Figure 3 is illustrated in Figure 4. The six magnetic islands that are finally formed are shown by the first column of Figure 11. We emphasize that similar steps of evolution from the reconnection layer to the magnetic island can be traced in all cases of island formation captured in our simulations. It is interesting to note that, similar to our case, magnetic reconnection is also invoked to explain the formation of solar prominence (Chen et al. 2020).

We speculate that the physical mechanism underlying the above processes is likely as follows. In the inner region of the accretion flow, the field lines are tangled due to the MHD turbulence driven by magnetorotational instability (Balbus & Hawley 1991). Thus, the field lines with opposite polarity can come close enough, resulting in the occurrence of reconnection. In addition, the Parker instability and wind launched from the accretion flow will make the magnetic field lines emerge out of the accretion flow into the corona. The differential rotation and turbulent motion of the accretion flow where the field lines are

---

We note that the appearance of a reconnection layer and a subsequent occurrence of magnetic reconnection are found to occur randomly, not necessarily simultaneously, above and below the equatorial plane.

---

Figure 2. Poloidal magnetic field structure of a flux rope identified in MAD and at a simulation time of $T = 20,000 r_g/c$ at various azimuthal angles $\varphi$ (denoted by the grid cell number $i$ at the top of each plot) in the $x-z$ plane. The magnetic field is denoted by black solid lines with arrows indicating the direction of the field, while the colors indicate the logarithmic value of the ratio of poloidal and toroidal magnetic field strength. The flux rope is located at $(x, z) = (5-10, 5-10)$, depending on the value of $\varphi$. In this example, the flux rope extends in the $\varphi$-direction for roughly $(63-41) \times 5.6 = 123^\circ$. 

---

We note that the appearance of a reconnection layer and a subsequent occurrence of magnetic reconnection are found to occur randomly, not necessarily simultaneously, above and below the equatorial plane.
rooted twist the field lines in the coronal region and make the reconnection occur there as well. This scenario of the formation of flux ropes is fully consistent with that proposed by Yuan et al. (2009). It has also been confirmed by other recent MHD numerical simulations (Nathanail et al. 2020, 2021; Ripperda et al. 2020, 2021), and it is confirmed again by the present analysis. A detailed comparison with these works will be presented in Section 5.

We find that the formation of a flux rope occurs not only in the coronal region but also within the accretion flow, as also pointed out in Nathanail et al. (2020). The reconnection rate is proportional to the Alfvén speed. Since $\beta \gg 1$ in the accretion flow, the reconnection rate is much smaller there compared to the coronal region. Combined with the fact that the energy density of the magnetic field in the corona is also stronger than that in the accretion flow (e.g., Yang et al. 2021), we expect that the energy release during reconnection occurs slowly and weakly and thus may not correspond to the observed strong flares. For these reasons, in the present paper, we only focus on the flux ropes formed in the coronal region of the accretion flow.

Because of the reconnection, the poloidal magnetic field at the center of the flux rope should become weaker compared to its surrounding medium. This is confirmed by the top panel of Figure 5. In fact, the movement of the flux ropes is followed by tracing the trajectories of those minima and the circular magnetic field line pattern. The released magnetic energy should be converted into the thermal and kinetic energy of plasma and be used to accelerate electrons whose radiation will be responsible for flares. In addition, the density within the flux rope is larger than its surrounding medium, as shown in the bottom panel of Figure 5. We speculate the reason as follows. The weakening of the magnetic field results in the decrease of magnetic pressure within the flux rope; consequently, the pressure from the surrounding medium compresses the flux rope and increases its density and pressure until a new pressure

---

**Figure 3.** Six reconnection layers from our SANE simulation (marked with blue circles). They will evolve and result in the occurrence of magnetic reconnection and formation of six magnetic islands in $100r_g/c$, shown by the first column of Figure 11. The density is shown by color, and the magnetic field is shown by solid lines with arrows indicating the direction of the field.

**Figure 4.** Evolution of the magnetic field structure around a reconnection layer shown in the rightmost panel of Figure 3, i.e., at $t = 19,400r_g/c$. Formation of the magnetic island and its further motion can be found in the bottom panels of Figure 11.
balance within and outside of the flux rope is established. Combining the top and bottom panels of the figure, we can find that the plasma $\beta$ within the flux rope should be higher, while the magnetization parameter $\sigma \equiv B^2/(\rho c^2)$ should be lower than its surrounding medium. These results are consistent with Ripperda et al. (2020).

To present quantitative results, the change of various components of the magnetic field and density as a function of $\varphi$ at a given cylindrical radius are given in Figure 6. Note that since the cylindrical radius of the flux rope is different for different $\varphi$, in this figure, the flux rope is present only at two places, i.e., $\varphi \sim 100^\circ$ and $350^\circ$. At these two places, the poloidal field reaches its minima, while the density reaches its maxima.

3.2. Ejection of Flux Ropes

Although flux ropes can be formed at both small and large radii, we find that there is an important difference between them. At small radii, within roughly $r \sim 10–15r_g$, although the flux ropes can be formed, few of them are ejected out. Many flux ropes just stay within the accretion flow and finally fall onto the black hole. This is likely because of the strong gravitational force of the black hole, similar to the existence of the “stagnation radius,” within which the motion of the matter is always toward the black hole. This is indicated by the direction of the velocity vectors close to the black hole shown in Yang et al. (2021; refer to Figure 3 therein). Such a result is also consistent with the absence of wind in the innermost region of the accretion flow (Yang et al. 2021). For other flux ropes, although they are found to be able to propagate outward for some distance, they quickly disappear because of the strong differential rotation of the accretion flow at small radii, which tears the flux ropes.

Beyond $10–15r_g$, however, we find that the flux ropes can be ejected out. This is because both the gravity and the differential rotation become weaker. The flux rope is found to become entangled soon after its ejection because its two end points are anchored in two different radii of the accretion disk that have different angular velocities. Such an entanglement would lead to the ejection of the flux rope due to the kink instability within an orbital timescale, similar to the solar and space physics cases (Wang et al. 2016; Sklodowski et al. 2021). The short lifetime is consistent with previous analytical arguments (Broderick & Loeb 2006) and perhaps also related to the short lifetime of the three flares detected by high-resolution GRAVITY observations (Gravity Collaboration et al. 2018).

We trace the motion of the broken flux rope (i.e., plasmoid) by examining its movement. The result is shown in Figure 7. We can see from the figure that the plasmoid is spiraling outward, reminding us of the helical structure of the magnetic field in a continuous jet. It is interesting to note that the corresponding $\theta$ angle of the plasmoid decreases with the increasing distance to the black hole, indicating a collimation effect, which is again similar to the behavior of continuous jets.
To give a quantitative result, we show in Figure 8 the radial location of a plasmoid in MAD as a function of time. The left panel of Figure 9 shows various components of the velocity as a function of time for this plasmoid. It reaches as high as 0.3 speed of light at $T = 23,000r_g/c$ from 0.05 speed of light at $T = 20,000r_g/c$. As a comparison, we show in the right panel of Figure 9 the evolution of various components of the velocity in the case of SANE. Compared to MAD, the typical values of both the magnitude of velocity and acceleration are much smaller.

To understand the physical reason for the acceleration of the plasmoid, we have analyzed the forces acting on the flux rope shown in Figure 8 (and the left panel of Figure 9). They include gravitational force, centrifugal force, Lorentz force, and the gradient of gas pressure. The result is shown in Figure 10. We can see from the figure that the Lorentz force is the dominant force that accelerates the plasmoid. The dominant component of the Lorentz force is the gradient of the magnetic pressure, which is enhanced by the occurrence of reconnection below the plasmoid. From this result, we can understand why the acceleration in the SANE case is much weaker than that in MAD; it is because the plasma $\beta$ in SANE is more than 10 times larger than in MAD (refer to Figures 1 and 4 in Yang et al. 2021). Such a scenario of the magnetic acceleration of the ejected plasmoid confirms the prediction in Yuan et al. (2009) and is very similar to the case of solar CMEs (Zhang & Low 2005).

3.3. Periodicity of the Formation and Ejection of Flux Ropes

We find from the simulation data that, in both SANE and MAD cases, every time magnetic islands are formed, their corresponding azimuthal angles $\phi$ are usually different; but their corresponding radii are almost always within the range of $20r_g \leq r \leq 30r_g$. More interestingly, the formation of the flux rope seems to occur quasiperiodically, once the simulations have reached the steady state. Such a phenomenon does not depend on model parameters or numerical resolution. As an example, we show in Figure 11 three successive cases of magnetic island formation and their subsequent evolution. The formation time of the three islands is $T = 17,500$, 18,500, and 19,500$r_g/c^{-1}$, respectively. More cases before $T = 17,500$ and after $T = 19,500r_g/c^{-1}$ exist but are not shown in the figure due to the spatial limitation. This implies a period of $\sim 1000r_g/c$. Such a period matches the orbiting timescale of both SANE and MAD at 20–30$r_g$ well, which is slightly larger than the Keplerian timescale, since the angular velocities of both SANE and MAD are sub-Keplerian (Yuan & Narayanan 2014). Since the formation of a flux rope at relatively large radii is accompanied by plasmoid ejection, such a periodicity implies that, in addition to flares, the ejection of plasmoids also occurs periodically. At smaller radii, the reconnection (and thus flare) is also found to occur periodically, although not accompanied by ejections.

The presence of periodicity combined with its orbital timescale strongly implies that, compared to turbulence, the differential rotation of the accretion flow seems to play a more important role in twisting the magnetic field lines and resulting in the occurrence of magnetic reconnection. This scenario is enhanced by the fact that there seems to be no difference in terms of the periodicity of the formation of the flux ropes in the SANE and MAD cases, although MAD is much less turbulent compared to SANE. In contrast, the occurrence of CMEs in the Sun mainly depends on solar activity. Although various periodicities exist for CMEs and flares (e.g., Lou et al. 2003), the periods are not related to the rotation period of the Sun. This is because, unlike the differentially rotating accretion flow, the geometry of the Sun is spherical, and the motion of the footpoints of the field loops is driven by the solar dynamo through local convection rather than by the rotation of the Sun.

In addition to the periodicity, another interesting finding is that in some cases, when a plasmoid moves outward, it may catch up and merge with a previously produced plasmoid having a slower speed. The merger of two magnetized islands must be accompanied by magnetic reconnection. This will result in strong radiation, so strong flares are expected to occur. Yuan & Zhang (2012) proposed this scenario as the mechanism of gamma-ray bursts.
The Astrophysical Journal, 933:55 (13pp), 2022 July 1

Čemeljić et al.

4. Applications to Observations

4.1. Periodic Flares and the Hot Spot in Sgr A* and Other Black Holes

The formation of the flux rope occurs because of magnetic reconnection, which must be accompanied by electron acceleration. The synchrotron radiation of the accelerated electrons will be responsible for the observed flares in various wave bands ranging from IR to X-ray and γ-ray flares in black hole sources mentioned above.

Near-infrared flares in Sgr A* were recently resolved by high-resolution GRAVITY observations, with a compact “hot spot” detected orbiting around the black hole with an orbiting radius of 6–10r_g (Gravity Collaboration et al. 2018). This hot spot is produced by the magnetic reconnection found in our simulations at small radii of the accretion flow. The flares are found to be polarized, with the polarization signature being consistent with the orbital motion of a hot spot in a strong poloidal magnetic field (Gravity Collaboration et al. 2018). In our simulation, we find that at the place of reconnection (i.e., the place of the hot spot), which is also roughly where the flux rope is located, the magnetic field could be dominated by either poloidal or toroidal components, depending on time. This is shown in Figure 12. The top panel shows a case where the field is dominated by the poloidal component, which is consistent with the observation. But at another time, shown in the bottom panel of this figure, the field is dominated by a toroidal component. This prediction could be verified by future observations. The detailed phenomenological modeling of the IR and X-ray light curves of Sgr A* indicates that to explain the simultaneity and symmetry of the X-ray and near-infrared light curves, the magnetic field strength is required to significantly decrease during the flare (Dodds-Eden et al. 2010). This is consistent with the weakening of the magnetic field during reconnection that we have described above and explicitly shown in Figure 5.

From the over 2 hr long infrared light curve of Sgr A*, several periods of variability are identified by observations, with periods ranging from 17 to 40 minutes (Genzel et al. 2003; Eckart et al. 2006a; Trippe et al. 2007; Genzel et al. 2010). The observed periods are consistent with the orbital period at 3–5r_g for the 3.6 × 10^6 M_☉ black hole in Sgr A*, assuming the accretion flow in Sgr A* is described by an MAD. Given that the typical lifetime of a hot spot is shorter than one orbit timescale, as we have explained before, the detected periodicity is best explained by the periodic occurrence of reconnection events found in our simulation at small radii. The observed period is not a fixed value but has a range, consistent with our finding that reconnection occurs in a range of radii.

4.2. Ejected Blobs in Sgr A*, M81, and M87: Origin, Velocity, and the Possible Periodicity

In our simulations, the flux ropes (i.e., plasmoids) formed beyond 10–15r_g will be ejected outward from the surface of the accretion flow. This explains the origin of the radio-emitting blobs observed in Sgr A*, M81, M87, and other black hole sources. The observationally measured ejection velocity of 0.4, 0.5, and 0.2 of the speed of light for the ejected plasmoids in Sgr A*, M81, and M87 is also quantitatively similar to what we have obtained in the case of MAD, as shown in Figure 9, which is about 0.3 of the speed of light. The “MAD” nature of the accretion flow in M81 and M87 is consistent with most recent studies (Event Horizon Telescope Collaboration et al. 2021; Shi et al. 2021; Yuan et al. 2022).

Examining the periodicity of the ejection of plasmoids is difficult, since it requires very high spatial resolution and
Figure 11. Three successive cases of magnetic island formation at $T = 17,500, 18,500,$ and $19,500 r_g/c^2$, just after the reconnection events shown in Figure 3, and their subsequent evolution. The data are based on the SANE simulation. Colors denote density, while solid lines denote the magnetic field. The dashed black and red lines are spherical coordinate lines passing through the centers of two magnetic islands located at the crossings of the circular and straight dashed lines of the same color. The value of $\varphi$ is indicated at the top of the upper right corner of each panel.
sensitivity of the telescope and the suitable observational targets. The high-resolution Global Millimeter VLBI Array observation of the M87 jet reveals some discrete blobs at a distance of about 100$\ r_g$ from the central black hole (EHT MWL Science Working Group et al. 2021). By measuring the spacing between two adjacent blobs from the image and adopting a projected moving speed of 0.2 of the speed of light (Hada et al. 2014), we can estimate that the time interval of two adjacent ejections is about 1 yr. This is in good consistency with our predicted period of 1000$\ r_g/c \sim 1$ yr for the $6 \times 10^9$ $M_\odot$ supermassive black hole in M87 (Gebhardt et al. 2011).

Most recently, two quasiperiodic oscillations (QPOs) were detected in the TeV blazar PKS 1510–089 by the Fermi Large Area Telescope (Roy et al. 2021). One is a 3.6 day QPO that lasted for five cycles with a moderate significance of $\sim 3.5 \sigma$, and another is a 92 day QPO that lasted for seven cycles with a significance of $\sim 7.0 \sigma$. As discussed in Roy et al. (2021), models like binary black holes and a procession jet are unlikely; two attractive models are a rotating hot spot in the inner region of the accretion flow and “equipaced magnetic islands” inside the jet. The remaining problem for the first model is that, since the accretion flow in a blazar should have an almost face-on orientation, it may be hard for the rotating hot spot to produce enough variability.

Obviously, these two scenarios are very similar to what we have proposed in the present work. The only difference is that, for the first hot-spot model, we think the variability is not caused by a single hot spot orbiting around the black hole but rather several hot spots periodically appearing in the inner region of the accretion flow; since the typical lifetime of a hot spot is shorter than one orbital time, it is difficult to explain the several observed cycles. Moreover, our “modified” hot-spot model has the advantage of overcoming the “not enough variability” difficulty.

4.3. Possible Application to Protostellar Accretion Systems

Although our simulations are done for hot accretion flows around black holes, we speculate that the basic physics relevant to flares and ejections should also work for accretion disks in protostars. This is why very similar phenomena have been observed in these systems. Specifically, in this case, the period should correspond to the orbital period at the inner edge of the accretion flow, where it is truncated by the magnetosphere of the protostar, which is a few days. Periodic ejections are observed from the spacings of the jet knots (Lee 2020) in those systems, and the period can be estimated from the spatial intervals and jet velocities. Unfortunately, the current instrumental limit of spatial resolution of $\sim 0\arcsec 1$ sets a limit on the period detection no shorter than about 1 yr (Lee 2020). On the other hand, the Chandra light curve does show obvious periodicity with a period of $2 \times 10^5$ s or a few days, according to Wolk et al. (2005) and Flaccomio et al. (2012; reference to Figure 3 in Wolk et al. 2005 and Figure 4 in Flaccomio et al. 2012), consistent with our prediction.

5. Comparison with Previous Works

The most relevant previous work is Yuan et al. (2009). In that work, by analogy with the CMEs in the Sun, an analytical model was proposed to explain the formation of episodic jets and their association with flares. The basic scenario proposed in that work, i.e., flux ropes are formed due to magnetic reconnection and then ejected out by magnetic forces, is fully confirmed by the present numerical simulations. Due to the limitation of analytical calculations, although all of the physical processes involved are highly dynamic, many assumptions and simplifications have to be adopted in Yuan et al. (2009). For example, the formation of a flux rope is purely an assumption there. The quantitative calculation of the ejected velocity of the flux rope is also based on assumptions on the distributions of the magnetic field and density of the coronal gas. Compared to Yuan et al. (2009), our current work is much more realistic and comprehensive. Moreover, important new findings are obtained, e.g., the periodic formation of the flux rope and its...
ejection and the quantitative prediction of the velocity of the ejected blobs.

Recently, several MHD numerical simulation works focusing on the analysis of reconnection in accretion flows have been published. It is thus useful to compare our simulations with these works. In the two-dimensional general relativity ideal magnetohydrodynamic numerical simulations by Nathaniel et al. (2020), the formation and evolution of current sheets are investigated. Special attention is paid to the effect of different initial magnetic field configurations. Similar to our work, they find that plasmoids are formed both within and at the surface of the accretion flow due to magnetic reconnection, which is driven by the turbulent motion of the accretion flow. Moreover, they find that plasmoids are formed during the reconnection, which is again consistent with our work.

In works by another group, two- and three-dimensional MHD simulations including weak ohmic resistivity are performed, focusing on the modeling of magnetic reconnection and the formation of plasmoids (Ripperda et al. 2020, 2021). Since an explicit resistivity is included that can approximately mimic the kinetic effects, their simulations are more suitable than ours to describe magnetic reconnection. Both SANE and MAD are studied as in our work. They find that within about $20r_g$ of the accretion flow current sheets and plasmoids are ubiquitous features that form regardless of the initial magnetic field, the magnetization of the accretion flows, and the spin of the black hole.

Different from our work and Nathaniel et al. (2020), another mechanism for the formation of plasmoids is identified in Ripperda et al. (2020), namely, the tearing instability of the current sheet. When the reconnection layer becomes thin enough, the reconnection layer can break up and produce chains of plasmoids. These small plasmoids can merge and grow to macroscopic scales of the order of a few Schwarzschild radii and be advected along the jet sheath or into the disk. We do not find such an instability in our simulations. For this instability to occur, the Lundquist number $S = v_A L / \eta$ needs to be larger than $10^5$. Here $v_A$ is the Alfvén speed, $L$ is the typical length of the current sheet, and $\eta$ is the resistivity (McKinnery et al. 2012). In our simulations, $S \ll 10^5$ because of the high numerical resistivity caused by the relatively low resolution of our three-dimensional simulations. In reality, both the turbulence of the accretion flow and the tearing instability in the current sheet should be able to drive reconnection and the formation of plasmoids, as seems to be confirmed by a recent work (Rosenberg & Ebrahimi 2021). In this work, resistive MHD equations are solved, and reconnection and plasmoids are found to be formed both via merging current sheets due to turbulent motions of the accretion flow and through reconnection due to current sheet instabilities.

### 6. Summary

Episodic ejections of blobs (episodic jets) have been observed in both black hole and protostellar accretion systems, and they are often associated with radiation flares. Notable sources include Sgr A*, M81, and M87, among others. Yuan et al. (2009) proposed an analytical MHD model for an episodic jet and its association with flares. In the present work, we develop that work by performing a three-dimensional general relativity magnetohydrodynamical (GRMHD) numerical simulation of hot accretion flows. Both SANE and MAD are considered. We have analyzed the simulation data and obtained the following results.

1. We find that once the simulations have reached the steady state, flux ropes keep forming both within and at the coronal region of the accretion flow from radii as small as $4r_g$ up to $\sim 30r_g$. One example of a flux rope in the case of MAD is shown in Figure 2. It extends in the $\varphi$-direction for about $120^\circ$.

2. We find that the flux rope is formed due to the reconnection of magnetic field lines, driven by the differential rotation and turbulent motion of the accretion flow. This scenario is consistent with that proposed in Yuan et al. (2009) and recent MHD numerical simulations. A detailed comparison between the present work and previous works is presented in Section 5.

3. We find that flux ropes formed at small and large radii have different fates. At small radii, inside of $\sim 10$–$15r_g$, few flux ropes can be ejected out, likely because of the strong gravitational force of the black hole, similar to the absence of wind at small radii. Beyond that radius, flux ropes can be ejected out. The ejected flux ropes are found to break within one orbital timescale due to the kink instability.

4. The broken flux ropes (i.e., plasmoids) spiral outward with the radial velocities shown in Figure 9. A clear acceleration in both cases is found, and the velocity can reach $0.3c$ and $0.08c$ for MAD and SANE, respectively. We have analyzed the forces and found that the acceleration is mainly due to the gradient of the magnetic pressure, as shown in Figure 10, again consistent with Yuan et al. (2009).

5. All of the processes mentioned above, i.e., the formation of the flux ropes and their subsequent ejection, occur periodically, with the period being the orbiting time of the radius where the flux ropes are formed. This implies that differential rotation of the accretion flow plays a more important role than turbulence in twisting the magnetic field lines and results in the formation of reconnection.

6. Magnetic reconnection should accelerate some electrons; their synchrotron emission explains the observed flares and the origin of the hot spot in Sgr A* detected by GRAVITY. The periodicity detected from the IR light curves of Sgr A* and X-ray light curves of blazar PKS 1510–089 is explained by the periodic formation of flux ropes.

7. The observed ejection of blobs is explained by the ejection of flux ropes formed at relatively large radii. In Sgr A*, M81, and M87, the measured velocity of the ejected blobs agrees with the velocity predicted in our simulations of MAD. The predicted periodicity of the ejection seems to be supported by some observations of M87 and PKS 1510–089.

8. Finally, we speculate that our results may also be able to interpret the flares and episodic ejection observed in protostellar systems. The measured period in the X-ray light curve is roughly consistent with the orbital period at the inner edge of the accretion flow, where it is truncated by the magnetosphere of the protostar.

We are thankful for useful discussions with Drs. P. F. Chen, J. Lin, and Y. M. Wang on the solar flares and coronal mass ejections and C. White for his help with using the ATHENA++.
code. The anonymous referee is acknowledged for constructive suggestions and comments. This work has made use of the High Performance Computing Resource in the Core Facility for Advanced Research Computing at Shanghai Astronomical Observatory. We also thank ASIAA, Taipei, for the use of their visualization servers. This work was supported by CAS President’s International Fellowship for Visiting Scientists grant 2020VMC0002 (M.C.); Polish NCN grant 2019/33/B/ST9/01564 (M.C.); Natural Science Foundation of China grants 12133008, 12192220, and 12192223 (F.Y., H.Y.); and Ministry of Science and Technology of Taiwan grants 108-2112-M-001-009 and 109-2112-M-001-028 (H.S.).

Appendix
Three-dimensional GRMHD Numerical Simulations of a Hot Accretion Flow

We have performed numerical simulations in three dimensions by solving the equations of ideal MHD describing the evolution of the accretion flow around a black hole in the Kerr metrics using the GRMHD code Athena++ (White et al. 2016; Stone et al. 2020). Readers are referred to Yang et al. (2021) for the details of the simulations; here we only present a brief overview. All of our simulations are performed in Kerr–Schild (horizon-penetrating) coordinates \((t, r, \theta, \varphi)\). The comoving rest-mass density is denoted by \(\rho\), and the component of the coordinate frame four-velocity is denoted by \(u^\mu\). The equation of the state of the gas \(u = \rho p_{\text{gas}} / (\Gamma - 1)\), where \(p_{\text{gas}}\) is the gas pressure of the comoving mass, \(u\) is the internal energy of the gas, and \(\Gamma\) is the adiabatic index, which we set to \(\Gamma = 4/3\). The units we adopt are Heaviside–Lorentz with both the light speed and gravity constant set to unity and the sign convection of metric \((- , + , + , + )\). The metric is stationary, and the self-gravity of the accretion flow is ignored in our simulation.

We have simulated four models: SANE00, SANE98, MAD00, and MAD98. They denote the SANE accretion flow around a black hole of spins \(a = 0\) and \(0.98\) and the MAD accretion flow around a black hole of spins \(a = 0\) and \(0.98\), respectively. The initial condition of all models is a rotating torus (Fishbone & Moncrief 1976) around a black hole, with the inner edge of the torus at \(r = 40.5\) \(r_g\) and the radius of the pressure maximum at \(r = 80\) \(r_g\). We have also added a poloidal magnetic field threading the torus in the way described by Penna et al. (2013). For MAD cases, we set one poloidal loop threading the whole torus. For SANE cases, we initially set up a seed field consisting of multiple poloidal loops of the magnetic field with changing polarity. We use the gas-to-magnetic pressure ratio \(\beta \equiv p_{\text{gas}} / p_{\text{mag}}\) to normalize the

![Figure 13](https://example.com/figure13.png)

Figure 13. Values of \(Q_\theta\) (upper) and \(Q_\varphi\) (lower) in the inner region of MAD98 (left) and SANE98 (right), calculated based on our simulation data.
magnetic field. For MAD, we set $\beta_{\text{min}} = 0.1$. For SANE, we set $\beta_{\text{min}} = 0.05833\ (a = 0)$ and $\beta_{\text{min}} = 0.1\ (a = 0.98)$.

The grid we use is a static mesh refinement grid, which allows the use of higher resolution where it is needed in the simulation. We use a logarithmically spaced radial grid with an inner edge of $1.1r_g$ and outer edge of $120r_g$. The grid cells in the polar and azimuthal directions are uniform. The final grid resolution in our simulation is $288 \times 128 \times 64$ cells in the $r$, $\theta$, and $\phi$ directions for SANE00 and MAD00, while for SANE98 and MAD98, it is $358 \times 128 \times 64$. We use the outflow boundary conditions at the inner and outer boundaries. For the $\theta$ and $\phi$ directions, we use the polar axis and periodic boundary conditions, respectively.

We have run simulations up to $t_f = 40,000, 80,000, 40,000,$ and $40,000r_g^{-1}$ for SANE00, SANE98, MAD00, and MAD98, respectively. They correspond to 8.9, 17.8, 8.9, and 8.9 orbital periods of the accretion flow at the pressure maximum. The “inflow equilibrium” has been reached at roughly 26, 40, 80, and $80r_g$ for these four models, respectively. Figure 1 shows the snapshot distribution in the $x$-$z$ plane of the gas-to-magnetic pressure ratio $\beta$ (color) and the magnetic field lines for SANE and MAD.

The main physical implication of the inflow equilibrium radius is that, within this radius, the inflow has reached the steady state; thus, the radial profiles of density and velocity will not change with time and so are reliable. For our problem, our main concern is not the density of the accretion flow but the configuration of the magnetic field, which is determined by the turbulence. Beyond the inflow equilibrium radius, although the density and velocity are not reliable, the level of turbulence is still reliable up to a much larger “turbulence radius,” the limiting radius of the turbulence steady state. This radius can be estimated as follows. Turbulence in the accretion flow is because of magnetorotational instability (MRI). The fastest growth rate of MRI at a radius $r$ is $\sim \Omega(r)$. More precisely, it takes three to four orbits for MRI to develop (Hawley et al. 1995). At the end of our simulations to the four models, taking a timescale of three orbits, we can obtain that the turbulence radii are $\sim 130, 200, 120,$ and $120r_g$, respectively. Many analyses in the present paper are performed at $r = 20,000r_g/c$. At that time, the turbulence radii are $\sim 100, 100, 80,$ and $80r_g$, respectively.

To check the resolution of our simulations, following Hawley et al. (2011), we have calculated $Q_\theta$ and $Q_\phi$ based on our simulation data,

$$Q_\theta = \frac{\Omega}{\Omega_{\text{MRI}}} = \frac{2\pi |v_{\theta,A}|}{\Omega dx^2},$$

$$Q_\phi = \frac{\Omega}{\Omega_{\text{MRI}}} = \frac{2\pi |v_{\phi,A}|}{\Omega dx^2}.$$

Here $v_{\theta,A}$ is the $\theta$- and $\phi$-directed Alfvén speed, and $\Omega$ is the angular velocity of the fluid. The physical meaning of these two parameters is the number of grids in the fastest-growing wavelength of MRI in the $\theta$- and $\phi$-directions, respectively. According to Hawley et al. (2011), if $Q_\theta > 10$ and $Q_\phi > 20$, the resolution of the simulation will be high enough to give quantitatively converged results for the accretion flow. The calculated values of $Q_\theta$ and $Q_\phi$ in the inner region of MAD98 and SANE98 are given in Figure 13. We can see from the figure that this criterion is well satisfied for MAD98 and marginally satisfied for SANE98. The values of $Q_\theta$ and $Q_\phi$ are larger in MAD than in SANE. This is because in the case of MAD, the Alfvén speed is larger, and the angular velocity is smaller; thus, the fastest-growing wavelength of MRI is larger.

The above approach is used to estimate the convergence of nonlinear MHD turbulence in accretion flows resulting from the magnetorotational instability. In addition to this approach, we have also adopted another more straightforward approach to verify the convergence of our results. That is, we have simulated the SANE98 and MAD98 models with a higher resolution of $704 \times 256 \times 128$. Each simulation costs us about 400K CPU hours. We have used these simulation data to repeat the analysis performed in the present paper, including the occurrence of magnetic reconnection, the formation of flux ropes, and the ejection speed of the plasmoids. We find that all of the results remain quantitatively unchanged compared to the case of low-resolution simulations, indicating the convergence of our results.
Rosenberg, J., & Ebrahimi, F. 2021, ApJL, 920, L29
Roy, A., Sarkar, A., Chatterjee, A., et al. 2022, MNRAS, 510, 3641
Scepi, N., Dexter, J., & Begelman, M. C. 2022, MNRAS, 511, 3536
Shi, F., Li, Z., Yuan, F., & Zhu, B. 2021, NatAs, 5, 928
Skodowski, K. D., Tripathi, S., & Carter, T. 2021, JPlPh, 87, 905870616
Stone, J. M., Tomida, K., White, C. J., & Felker, K. G. 2020, ApJS, 249, 4
Tchekhovskoy, A., Narayan, R., & McKinney, J. C. 2011, MNRAS, 418, L79
Trippe, S., Paumard, T., Ott, T., et al. 2007, MNRAS, 375, 764
Wang, Y., Zhuang, B., Hu, Q., et al. 2016, JGRA, 121, 9316
White, C. J., Stone, J. M., & Gammie, C. F. 2016, ApJS, 225, 22
Wilms, J., Pottschmidt, K., Pooley, G. G., et al. 2007, ApJL, 663, L97

Wolk, S. J., Harnden, F. R. J., Flaccomio, E., et al. 2005, ApJS, 160, 423
Yang, H., Yuan, F., Yuan, Y.-F., et al. 2022, ApJ, 914, 131
Yuan, F., Lin, J., Wu, K., & Ho, L. C. 2009, MNRAS, 395, 2183
Yuan, F., & Narayan, R. 2014, ARA&A, 52, 529
Yuan, F., Quataert, E., & Narayan, R. 2004, ApJ, 606, 894
Yuan, F., Wang, H., & Yang, H. 2022, ApJ, 924, 124
Yuan, F., & Zhang, B. 2012, ApJ, 757, 56
Yusef-Zadeh, F., Roberts, D., Wardle, M., Heinke, C. O., & Bower, G. C. 2006, ApJ, 650, 189
Zhang, M., & Low, B. C. 2005, ARA&A, 43, 103
Zhao, T.-L., Yuan, Y.-F., & Kumar, R. 2020, MNRAS, 499, 1561