TWO TYPES OF ERGOSPHERIC JETS FROM ACCRETING BLACK HOLES: THE DICHOTOMY OF FANAROFF-RILEY GALAXIES

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\section{1. INTRODUCTION}

Remarkable advances have been made in explaining the observed spectrum of black hole X-ray binaries (BHXBs) and active galactic nuclei (AGNs) with the help of theoretical studies on black hole (BH) accretion disks \citep{Abramowicz99, Kato08}. Current stable BH accretion disk solutions include the advection-dominated accretion flow, or ADAF for short \citep{Narayan94, 1995ApJ...441..733N, 1997ApJ...476..375N}, the thin disk \citep{Shakura73, Novikov73}, and the slim disk \citep{1988ApJ...329..477A, 2010ApJ...714..969A, Sadowski2009, Sadowski2011}. The accretion disk type is a function of the the dimensionless accretion rate, $\dot{m} \equiv \dot{M}/M_{\text{Edd}} = L/L_{\text{Edd}}$, where $\dot{M}$ is the accretion rate, $M_{\text{Edd}}$, $L_{\text{Edd}}$, and $L$ are the Eddington accretion rate, the Eddington luminosity, $L_{\text{Edd}}$, and the Eddington luminosity, $L$, is the efficiency ($\xi = 1$ is adopted for later computation) and $c$ is the speed of light. In Table 1 we summarize the properties of the above disk solutions.

The accretion disk solution is also a function of radius \citep{Abramowicz93, Chen95}. A combined disk which consists an outer thin disk and an inner ADAF has been also proposed to explain observations of BHXBs and AGNs \citep{Narayan96, 1997MNRAS.287..705E, 2004ApJ...605L..81F, 2011ApJ...737..101T}. Such disk is expected to be formed before the entire disk transits from an ADAF to a thin disk, because the hot ADAF would cool down from outside. In other words, the transition radius from outer thin disk to inner ADAF solution decreases as the accretion rate increases until the entire disk become thin disk \citep{Honma96, Kato98}. For a comprehensive consideration of how accretion disk type vary with $\dot{m}$, we also include the the combined disk in Table 1.

As a pure relativistic effect, the specific angular momentum of a Keplerian rotating object near the BH, $\ell_k$, has a local minimum at the radius of the innermost stable circular orbit, $R_{\text{ISCO}}$. The flow geometry near the BH horizon is determined by how the angular momentum distribution of the accretion disk, $\ell(R)$, deviates from the Keplerian values \citep{Abramowicz98}. If $\ell(R) > \ell_{\text{ms}}$, where $\ell_{\text{ms}}$ is the Keplerian angular momentum at $R_{\text{ISCO}}$, $\ell(R)$ must equal the Keplerian value twice (assuming $\ell(R) = \ell_k$ at $R_1$ and $R_2$, and $R_2 > R_{\text{ISCO}} > R_1$). In this case, a pressure maximum is formed inside the disk (at $R_2$) and plasmas are "pushed" towards the BH, forming a "pressure-driven" flow and resulting in a "disk-like" structure on the equatorial plane near the horizon. On the contrary, if $\ell(R) < \ell_{\text{ms}}$, the sub-Keplerian plasmas fall into the BH because of the nature of their sub-Keplerian angular momentum ("viscous-driven"). The plasmas are able to fall closer to the axis (i.e., higher latitudes of the horizon), forming a "quasi-spherical" geometry near the horizon. These two kinds of accretion flow geometry near the horizon, disk-like and quasi-spherical, was first recognized by \citep{Abramowicz81}.

Figure 1 schematically depicts the geometric and dynamic properties of different accretion disks and the extraction of the BH rotational energy when different accretion flow geometries near the horizon is realized. If the magnetosphere near the horizon is filled with plasmas, those plasma can load onto the large-scale field lines and reduce the electromagnetic extraction of the BH energy \citep{Takahashi90}. In this letter, by assuming that a relativistic jet are launched when the BH energy is extracted outward, we investigate the formation of BH relativistic jets when the BH is surrounded by different type of disks listed in Table 1. This work is an extension of \cite{Pu12}, hereafter PHC, in which the formation of relativistic jets at low accretion rate (e.g. when $\dot{m} \lesssim 0.1$) is considered. It is found that, ergospheric jets can preferentially take place not only when the BH is sur-
rounded by a combined disk, as suggested in PHC, but also when the BH is surrounded by a slim disk. We refer the two types of jets as ‘EJC’ and ‘EJS’, respectively (see section 3). We conclude that the characteristic accretion rate of a EJC and a EJS are \( \dot{m} < 0.01 \) and \( \dot{m} > 0.01 \), respectively, and that the EJS is in general more powerful and faster than the EJC for the same central BH mass and BH spin. The difference between these two types of jets may result in the observed dichotomy of of FR I and FR II galaxies.

2. IDEAL MHD FLOW

In an axis-symmetry and stationary BH magnetosphere, the total energy of the ideal MHD flow, \( E_r \), is conserved along a specific field line (see, e.g., PHC for details). We can separate \( E \) into the plasma part (the first term) and electromanetic part (the second term),

\[
E = -\mu u_t + \frac{\Omega_F}{4\pi \eta} B_\phi ,
\]

where \( \mu \) is the relativistic specific enthalpy, \( u_t \) is constant time component of the four—velocity, \( \Omega_F \) is the angular velocity of the field, \( \eta \propto nu'^2/B_r \) is the particle flux per unit flux tube, \( n \) is the proper number density, \( u' \) is the radial component of the four—velocity and \( B_\phi \) and \( B_r \) is the toroidal and radial field observed by distant observer. Similarly, the total outward energy flux \( E' \) has the plasma part, \( E_{'\text{em}} \), and the electromagnetic part, \( E_{'\text{ ELF}} \),

\[
E' = nE'u' = E_{'\text{em}} + E_{'\text{ ELF}} .
\]

3. TWO TYPES OF RELATIVISTIC JET

In general, the strength of a large-scale magnetic field strength\(^4\) can be parameterized by the ratio, \( \varepsilon_2 \), of the gravitational binding energy of the disk at radius \( R \) to the large-scale magnetic field energy inside \( R \) (see also

\[\text{(2)}\]

A plasma is accelerated along the large-scale magnetic field lines inwards (or outwards) if the gravitational force is larger (or less) than the magnetocentricrifugal forces. The conservation of \( E' \) near the separation region, where the magnetocentricrifugal forces balance

with the gravitational force, requires the solution of a powerful outflow (i.e., relativistic jet) to that of an inflow with a positive energy flux \( (E_r > 0) \). As discussed in §3.6 of PHC, for the inflow, \( E_{'\text{em}} > 0 \) and \( E_{'\text{ ELF}} < 0 \) are generally expected; that is, the magnetic field lines contribute to the extraction of the rotational energy of a BH, while the plasma plays an opposite role and reduces the extraction. Therefore, an MHD flow can extract the BH energy (accordingly a ergospheric jet is launched), only when the magnetosphere is magnetically dominated, \( |E_{'\text{em}}| > |E_{'\text{ ELF}}| \). The extraction of the rotational energy by the MHD flow is described by the ‘MHD Penrose process’ \((\text{Takahashi et al. 1990})\). When \( E_{'\text{ ELF}} \) is negligibly small compared to \( E_{'\text{em}} \), (i.e., a nearly vacuum environment), the force-free limit becomes a good approximation. The electromagnetic extraction of the rotational energy of a BH (by large-scale, hole-threading field lines) is described by the ‘Blandford-Znajek process’ \((\text{Blandford \\& Znajek 1977})\).

\[\text{(3)}\]

\[\text{(4)}\]

\[\text{TABLE 1}\]

| Disk Type          | ADAF   | Combined Disk | Thin Disk | Slim Disk |
|-------------------|--------|---------------|-----------|-----------|
| Accretion Rate    | \( \dot{m} \ll 0.01 \) | \( \dot{m} \lesssim 0.01 \) | \( 0.01 \lesssim \dot{m} \lesssim 0.3 \) | \( 0.3 \lesssim \dot{m} \) |
| Cooling Process   | advection | radiation | advection |
| Optical Depth     | thin   | thick         | thick     | thick     |
| Disk Geometry     | thick \( (H \sim R)^a \) | outer thin disk + inner ADAF | thin \( (H \ll R) \) | thick \( (H \ll R) \) |
| Viscosity         | \( \alpha \sim 0.1 \) | \( \alpha \sim 0.01 \) | \( \alpha \ll 1 \) |
| Infalling Process | viscous-driven \( (t_{in} < t_{ms})^b \) | viscous-driven \( (t_{in} \sim t_{ms}) \) | pressure-driven \( (t_{in} > t_{ms}) \) |
| Flow Geometry near the Horizon\(^d\) | quasi-spherical | quasi-spherical | disk-like |

\( a \) \( H \) is the disk height and \( R \) is the radial distance to the central BH.

\( b \) \( \alpha \) is the parameter in \( \alpha \)-prescription \((\text{Shakura \\& Sunyaev 1973})\).

\( c \) \( t_{in} \) is the angular momentum of the flow when it falls onto the BH.

\( d \) Note that the geometry of the accretion flow near the horizon is determined by the ‘infalling process’ (see text), therefore it can be different from the ‘disk geometry’, which describes the geometry of an accretion disk far from the horizon.

\( E_r \) is the radial component of the four—velocity and \( r \) is the radial distance to the central BH. Seem also Figure 1.
The mathematical similarity of the idea MHD assumption and the force-free assumption guarantees that $P_{em}$ is identical to the output power of Blandford–Znajek process:

$$P_{em} = \frac{1}{32} \omega_{F}^{2} B_{H}^{2} R_{H}^{2} \beta_{c} c ,$$

where $j \equiv J/J_{max} = J/(GM^2/c)$, $J$ the angular momentum of the BH, $B_{H}$ the large-scale hole-threading field strength, and $\beta_{c} \equiv \Omega_{F}/(\Omega_{H} - \Omega_{F})/\Omega_{H}$, $\Omega_{H}$ the rotational velocity of the BH. Assuming $\omega_{F} = 1/2$, which maximize $P_{em}$, we have,

$$P_{em} = \frac{1}{32} \frac{1}{4} B_{H}^{2} \left[ 1 + \sqrt{(1 - j^{2})} \right]^{2} \left( \frac{GM_{BH}}{c^{2}} \right)^{2} j^{2} c .$$
In slim disk case, the accreting plasma, which contributes to $P_{\text{plasma}}$, stay on the equatorial plane. As a result, there is only tiny amount of plasma loading onto the large-scale magnetic field lines that thread the horizon at different latitudes. The jet power of a EJS can be therefore estimated by Equation (7), namely that $P_{\text{EJS}} \approx P_{\text{em}}$. On the other hand, because the plasma with quasi-spherical geometry near the horizon can be loaded onto the field lines and can further reduced the power by the amount of $P_{\text{plasma}} < 0$, Equation (7) represents the maximum power of a EJC.

By Equation (3), $B_H$ can be computed if $\varepsilon$ and $\Sigma$ are known. Although the determination of $\varepsilon$ is unclear, it should in general has a value $1 > \varepsilon > 0$. Therefore, the changing of $B_H$ at different accretion rate should mainly be determined by the variation of the surface density of the disk, $\Sigma(m)$, because it will be higher variable than $\varepsilon(m)$. Here we assume that the large-scale, disk-threading fields can be dragged inward by accretion flow for all types of disk (Narayan et al. 2003; Spruit & Uzdensky 2003; Rothstein & Lovelace 2008) and hence $\varepsilon > 0$ is always satisfied. Note that, in a slim disk, it is the pressure–driven nature (instead of the viscosity) that determines the advection motion of the plasma near the BH. Therefore, although the the magnetic Prandtl number is expected to be small due to a relative tiny value of $\alpha$, compare to other types of disk (Table 1), the field strength near the horizon can be further enhanced when the disk transits form a thin disk to a slim disk, by further pushing a large-scale, disk-threading fields can be dragged inward.

Figure 2 presents typical BH accretion disc solution profiles on the $\Sigma - m$ plane, following the model in Abramowicz et al. (1995) and Chen et al. (1995). The solid curves represent the solutions of a ADAF, a thin disk and a slim disk with their typical viscosity (see figure caption). With the relation of $\Sigma - m$ and Equation (3), relative values of $B_H$ as a function of $m$ are qualitatively presented in Figure 3. The field strength gradually increases with increasing $m$ when the disk is either an ADAF or a thin disk because of the increase of $\Sigma$, where the radiation-pressure dominated thin disk solution is not of interest here because the disk become unstable in that case. For a combined disk, $B_H$ is essentially described by the surface density of the disk at the transition radius, $R_{tr}$, because the outer thin disk has a much larger density than that of the inner ADAF. The decrease of $R_{tr}$ with increasing $m$ results in a rapid growth of $B_H$ with $m$ (see also the Figures 1a and 1e of PHC). For a slim disk, although $B_H$ can be further enhanced due to the pushed-in large-scale, disk-threading field lines, the increase of $B_H$ is expected to saturate eventually. This is because the fields diffuse outward outside the pressure–driven region (that is, outside the super-Keplerian part of the disk).
The transition radius of a combined disk continuously decreases with increasing \( m \) before the entire disk become a thin disk. Previous studies (Honma [1996] Mammoto et al. [2000]) show that \( R_{tr} \) can reach down to \( \sim 10 R_g \). For illustration purpose, assuming a modest BH spin, \( j = 0.6 \), we compute the maximum value of \( B_H \) for a combined disk by using the maximum \( \Sigma \) of the (outer) thin disk solution at \( R = 10 R_g \) (the filled circle in Figure 2) and calculate the resulting maximum power of a EJC by equation (3). The jet power for three different \( \varepsilon \)'s are plotted on the \( M_{BH} - P \) plane (Figure 4). Noting that that the plasma contribution, \( P_{\text{plasma}} < 0 \), as well as a smaller BH spin, will reduce the power of EJC, we find these jet powers represent the upper limits. In other words, the power of a EJC will appear below these lines. On the contrary, for the same parameters (BH mass, BH spin, \( \varepsilon \)), sources with EJSs will appear above the lines because of a larger \( B_H \) (Figure 3) and a negligible \( P_{\text{plasma}} \) in the force-free limit.

3.2. Jet Speed

By denoting the quantities of the inflow \((u^i < 0)\) and outflow \((u^o > 0)\) with the subscript “in” and “out”, respectively, the ratio of the jet speeds of a EJC and a EJS can be estimated as follows: At large distances, the jet Lorentz factor, \( \Gamma \), can be defined by

\[
\Gamma = \frac{E_{out}}{E_{in}}.
\]

In addition, near the separation region of the inflow and the outflow, the conservation of \( E^r \) gives,

\[
E_{out} = \frac{n_{in} u_{in}^r}{n_{out} u_{out}^r} E_{in}.
\]

Thus, for fixed BH mass and spin, the ratio of the power of the EJC and the EJS, \( \chi \), can be written as,

\[
\chi \equiv \frac{\Gamma_{EJC}^\varepsilon}{\Gamma_{EJS}^\varepsilon} \approx \frac{(n_{out} u_{out}^r)^{EJS}}{(n_{out} u_{out}^r)^{EJC}} \frac{(B_r B_\phi)^{EJC}}{(B_r B_\phi)^{EJS}},
\]

by the help of Equation (1), where the superscript “EJC” (“EJS”) denotes the parameter of EJC (EJS). Because both \((n_{out} u_{out}^r)^{EJS}/(n_{out} u_{out}^r)^{EJC} < 1\) and \((B_r B_\phi)^{EJC}/(B_r B_\phi)^{EJS} < 1\) hold, we obtain \( \chi \ll 1 \).

4. DISCUSSION

The characters of EJC and EJS offer the key to an understanding of the FR I/ FR II dichotomy. The FR I/ FR II division line found by Ledlow & Owen (1996), which is shown by the thick solid line in Figure 4, is roughly consistent with the maximum power that a EJC can reach. It is, therefore, reasonable to suppose that the division reflects the two types of relativistic jets. That is, most FR I (or FR II) galaxies are associated with a EJC (or a EJS). Noting that a EJC has a slower speed than a EJS, why FR I and FR II galaxies have a edge-darken and an edge-brighten morphology, respectively, can be consistently understood as a result of different jet velocities (Bicknell [1985]).

Ghisellini & Celotti (2001) found that the accretion force of FR I and FR II galaxies can be separated at \( m \approx 0.01 \). Similar division of the accretion rates is also indicated for BL Lacs and radio quasars (which are believed to be ‘face-on’ FR I and FR II galaxies, respectively) (Xu et al. 2009, Ghisellini et al. 2011), and for radio galaxies of ‘high-excitation’ and ‘low-excitation’ (Best & Heckman 2012). The corresponding \( \varepsilon \) of EJC \((\varepsilon < 0.01)\) and EJS \((\varepsilon > 0.01)\) are consistent with the above findings.

The major drawback in our calculation is that, the Pseudo-Newtonian potential (Paczynsky & Wiita 1980) is adopted when we use used the disk model of Abramowicz et al. (1995) and Chen et al. (1995). We should notice a full-relativistic model (e.g. Sadowski (2009)), could significantly change the range of \( m \) of the slim disk solution (see captions of Figure 2). Nevertheless, the surface density of the disk solution will not be changed dramatically (Sadowski 2009, Abramowicz et al. 2010, Sadowski et al. 2011). We therefore expect that the results will be qualitatively unchanged when general–relativistic corrections are taken into account. The effect of the BH spin on the jet power, which is beyond the present model, can be investigated in the future by incorporating general relativistic corrections.

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