Role of magnetically dominated disc-outflow symbiosis on bright hard-state black hole sources: ultra-luminous X-ray sources to quasars

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ABSTRACT
We present optically thin solutions for magnetized, advective disc-outflow symbiosis around black holes (BHs). The main objective is to explain the bright hard state observations of accreting systems with stellar mass to supermassive BHs. We include the effects of magnetic field and radiation counterpart in entropy gradient based on first law of thermodynamics to represent energy advection. The cooling process includes bremsstrahlung, as well as synchrotron radiation. One of our main ventures is to explain some long standing issues of ultra-luminous X-ray sources (ULXs). The existing physical scenarios to explain their unusual high luminosity are either the existence of the missing class of intermediate mass BH (IMBH), or super-Eddington accretion around a stellar mass BH. However, most ULXs with steep power-law spectrum can be well explained through super-Eddington accretion, while the existence of IMBH is indeed disputed extensively. Nevertheless, the interpretation of ULXs with a hard power-law dominated state remains mysterious. Here we show that our magnetically dominated disc-outflow symbiosis around rapidly spinning stellar-mass BHs can achieve such large luminosity even for sub-Eddington accretion rate. The magnetic field at the outer zone of the advective flow is more than the corresponding Eddington limit. Such field becomes dynamically dominant near the BH through continuous accretion process due to flux freezing, but maintaining its Eddington limit. This unique field configurations enhance the synchrotron cooling process to achieve very large luminosity. Through the same mechanism, our solutions for supermassive BHs can explain the unusual large luminosity of ultra-luminous quasars.

Key words: accretion, accretion discs – black hole physics – MHD – gravitation – X-rays: binaries – quasars: supermassive black holes

1 INTRODUCTION

Galactic black hole candidates (BHCs) are known to pass through different X-ray spectral states based on their X-ray luminosity and spectral shape. Typically, the most familiar states are high/soft (HS), low/hard (LH), and very high/steep power-law (VH/SPL) states (Fender et al. 2004; Remillard & McClintock 2006). In HS state, the X-ray spectrum is dominated by a hard power-law component with an exponential high-energy cutoff at $\sim 200$ keV. The luminosity $L$ is quite below the Eddington luminosity $L_{\text{Edd}}$ ($L \sim 0.1 L_{\text{Edd}}$), which is defined as

$$L_{\text{Edd}} = \frac{4\pi GM_{\text{BH}}}{c}\kappa_{\text{es}} \approx 1.4 \times 10^{38} \left(\frac{M_{\text{BH}}}{M_\odot}\right)\text{erg s}^{-1},$$

where $M_{\text{BH}}$ is the mass of the black hole (BH), $G$ the Newton's gravitation constant, $c$ the speed of light, and $\kappa_{\text{es}}$ the electron scattering opacity. The X-ray spectrum in LH state is dominated by a standard geometrically thin, optically thick, radiatively efficient (Shakura & Sunyaev 1973) Keplerian
disc is formed. The HS state is governed by this accretion paradigm. For larger accretion rate \((\dot{m} \gtrsim 1)\), a slim disc (Abramowicz et al. 1988) is formed and the corresponding spectral state is VH/SPL. In this regime, photon trapping may take place due to very large optical thickness and, hence, some disc luminosity could be advected into the BH. On the other side, when accretion rate is very small \((\dot{m} \lesssim 0.01)\), the flow becomes optically thin. The Coulomb coupling between ions and electrons becomes very weak. The generated heat in dissipation process partly can be stored as entropy rather than being radiated out. The different classes of such radiatively inefficient accretion flows could be described by many models: an ion-supported torus (Rees et al. 1982) or an advection-dominated accretion flow (ADAF; Narayan & Yi 1995) or a two component accretion flow (TCAF; Chakrabarti & Titarchuk 1995) or an adiabatic inflow-outflow solutions (ADIOS; Blandford & Begelman 1999) or a convection-dominated accretion flow (CDAF; Narayan et al. 2000) or a general advective accretion flow (GAAF; Rajesh & Mukhopadhyay 2010). All these models can explain the hard spectral states.

Apart from these most familiar canonical states, there are large number of BH sources which are very bright but their spectra are hard power-law dominated. The RXTE observations of the BHC GX 339-4 show the luminosity up to \(\sim 0.3 E_{\text{edd}}\) in its hard spectral state (Miyakawa et al. 2008). Also a significant fraction of ultra-luminous X-ray sources (ULXs) in their hard spectral states are observed with luminosity in the range of \(3 \times 10^{39} - 3 \times 10^{40}\) erg s\(^{-1}\). The true nature of such sources (Antennae X-11, X-16, X-42, X-44, Feng & Kaaret 2009; NGC 1365, Soria et al. 2009; M99 X1, Soria & Wong 2006; M82 X42.3+59, Feng et al. 2010; Holmberg IX, Kaaret & Feng 2009) remains mysterious over decades. Not only stellar mass BHs, a large fraction of supermassive BHs also appears very bright in their hard power-law dominated states. Some of these such sources are ultra-luminous quasars (e.g. PKS 0743-67, Punsly & Tingay 2005; HS 1946+7658, Hagen et al. 1992), luminous BL Lac objects (e.g. PKS 0301-243, 1ES 0502+675, Ackermann et al. 2011).

Exact in this line our proposal comes in. Our group already initiated to explain the importance of strong magnetic fields in BH accretion sources in advective paradigm (Mukhopadhyay & Chatterjee 2015; Mondal & Mukhopadhyay 2018, 2019a,b). In Mondal & Mukhopadhyay (2019a), for the first time in literature, we suggest a plausible mechanism to explain the unusually large luminosity for certain ULXs. We showed that hard-state ULXs are magnetically powered sub-Eddington, advective accretors around stellar mass BH. However, we did not include the cooling processes explicitly therein. In this paper, we address a magnetized disc-outflow symbiosis with explicit cooling to explain the unusually large luminosity for certain BH sources. Finally we end with conclusion in Section 7.

The origin as well as strength of large-scale strong magnetic fields in BH accretion flows is still not well understood. However, a strong correlation between hard spectral states, powerful jets, and, dynamically dominant magnetic fields has been found. Recent observations confirm the signatures of dynamically dominant magnetic fields in the vicinity of BHs (Eatough et al. 2013; Zamaninasab et al. 2014). The magnetically dominated accretion flows have been studied in different versions of models, namely magnetically arrested disc (MAD; Narayan et al. 2003), or magnetically choked accretion flow (McKinney et al. 2012). Such models are based on the idea originally proposed by Bisnovatyi-Kogan & Ruzmaikin (1974), where strong vertical poloidal fields are dragged towards the central BH by continuous accretion process. This has been also verified numerically via pseudo-Newtonian magnetohydrodynamics simulations (Igumenshchev et al. 2003) and via general relativistic magnetohydrodynamics (GRMHD) simulations (Tchekhovskoy et al. 2011; McKinney et al. 2012). Apart from MAD, alternate magnetically supported accretion flow models have been discussed, in which the magnetic field geometry is dominated by radial and toroidal fields both and it operates only when the accretion rate is relatively high (Machida et al. 2006; Oda et al. 2010).

In this light of discussion, we propose a general advective, sub-Eddington disc-outflow symbiotic model in the presence of large-scale strong magnetic fields. Unlike MAD, the advection of both toroidal and poloidal magnetic fields is happening here. This model can explain the bright hard spectral state of BH sources of mass ranging from stellar mass to supermassive scales. Most importantly, we address the hidden nature of hard-state ULXs without incorporating the existence of intermediate mass BHs.

The paper is organized as follow. In the next section, we model the coupled general advective disc-outflow symbiosis including thermodynamic properties in the presence of magnetic field and radiation. The solution procedure along with appropriate boundary conditions is mentioned in Section 3. In Section 4, we discuss our results which cover the disc flow behaviours, magnetic field properties and the energetics of this magnetized accretion induced outflows. In Section 6, we discuss the possible origin and the strength of magnetic fields in this advective paradigm. We also discuss the implication of our model, particularly for very luminous hard-state BH sources. Finally we end with conclusion in Section 7.
2 MODELING THE COUPLED MAGNETIZED DISC-OUTFLOW SYSTEM

2.1 General equations of magnetized advective accretion flow

We formulate a general magnetized disc-outflow symbiotic model around BHs in geometrically thick, optically thin, advective accretion framework. Unlike previous exploration (e.g. Kuncic & Bicknell 2004), we consider all possible viscous stresses as well as large-scale magnetic stress. Indeed, it was initiated earlier by us, without considering cooling effect explicitly (Mondal & Mukhopadhyay 2019a). As a consequence, we could not comment out luminosity of the system explicitly. In this disc-outflow symbiosis, we adopt the cylindrical coordinate system assuming a steady and axisymmetric flow such that \( \partial / \partial t \equiv \partial / \partial \phi \equiv 0 \). This 2.5-dimensional quasi-spherical advective flow describes the inner part of the accretion where gravitational force dominates over the centrifugal force of the flow, unlike standard Keplerian disc. All the dynamical flow parameters, namely, radial velocity (\( v_r \)), specific angular momentum (\( J \)), outflow or vertical velocity (\( v_z \)), adiabatic sound speed (\( c_s \)), fluid pressure (\( P \)), mass density (\( \rho \)), radial (\( B_r \)), azimuthal (\( B_\phi \)), and vertical (\( B_z \)) components of magnetic field, are functions of both radial and vertical coordinates. In this formalism, we use pseudo-Newtonian potential, given by Mukhopadhyay (2002), which mimics certain features of general relativity quite accurately. The other key features in this model are as follow. First, the vertical flow is included here explicitly, and unlike Bhattacharya et al. (2010) it is coupled to the other flow parameters through fundamental equations of motion. The outflows are more likely to emanate from the hotefficiency up region of the accretion flow. Second, we include the effect of viscosity by taking care of all possible components of viscous shearing stress. Third, the effect of large-scale magnetic field geometries is included explicitly, unlike Mandal & Chakrabarti (2005) or Rajesh & Mukhopadhyay (2010). Fourth, we include the effects of magnetic field and radiation counterpart in the computation of entropy gradient and adiabatic exponents based on first law of thermodynamics.

Throughout in our calculations, we express all the flow variables in dimensionless units. The radial and vertical co-ordinates are expressed in units of the gravitational radius \( r_g = GM_{BH}/c^2 \). Any flow velocities are expressed in units of \( c \), the specific angular momentum in \( GM_{BH}/c \), the fluid pressure, mass density and magnetic fields, accordingly, to make all the variables dimensionless. Hence, the continuity equation and the components for momentum balance equation are respectively,

\[
\frac{1}{r} \frac{\partial}{\partial r} (r \rho v_r) + \frac{\partial}{\partial z} (\rho v_z) = 0, \tag{2}
\]

\[
v_r \frac{\partial v_r}{\partial r} + v_z \frac{\partial v_z}{\partial z} - \frac{L^2}{r^3} + \frac{1}{\rho} \frac{\partial P}{\partial r} + F = \frac{\partial W_{r\phi}}{\partial r} \frac{\partial B_r}{\partial \phi} + B_r \frac{\partial B_r}{\partial z} - \frac{\partial B_z}{\partial \phi}, \tag{3}
\]

\[
v_r \frac{\partial v_z}{\partial r} + v_z \frac{\partial v_z}{\partial z} + \frac{\partial P}{\partial z} + F \frac{\partial B_z}{\partial \phi} \frac{\partial B_r}{\partial \phi} + B_r \frac{\partial B_r}{\partial z} - \frac{\partial B_z}{\partial \phi} \left\{ \frac{B_r}{r} \frac{\partial B_r}{r} + B_z \frac{\partial B_z}{\partial \phi} \right\}. \tag{4}
\]

Here \( F \) is the magnitude of the force corresponding to the gravitational potential for a rotating BH in the pseudo-Newtonian framework (Mukhopadhyay 2002), as given by

\[
F = \frac{(r^2 - 2a\sqrt{r} + a^2)^2}{r^3 \left[ \sqrt{r} (r - 2) + a^2 \right]^2}, \tag{5}
\]

where \( a \) is the Kerr-parameter. \( W_{r\phi} \) are the components for viscous shearing stress tensor. Following standard practice, the \( W_{r\phi} \) component is written using standard-disc (Shakura & Sunyaev 1973) prescription with proper modification due to advection (Chakrabarti 1996) as \( W_{r\phi} = a(p + \nu v_r^2) \), where \( a \) prescribes the turbulent viscosity. The other components can be simplified in terms of \( W_{r\phi} \) as \( W_{r\phi} = \frac{5}{3} \pi W_{r\phi} \) (Ghosh & Mukhopadhyay 2009).

The magnetohydrodynamics (MHD) flow provides two fundamental equations for magnetism, namely, the equation for no magnetic monopole and the induction equation. These are respectively,

\[
\frac{1}{r} \frac{\partial}{\partial r} (r B_r) + \frac{\partial B_z}{\partial z} = 0, \tag{6}
\]

\[
\frac{\partial}{\partial z} \left[ r (v_z B_r - v_r B_z) \right] = 0. \tag{7}
\]

\[
\frac{\partial}{\partial r} \left( \frac{v_r B_\phi - \frac{1}{2} B_r^2}{r} - \frac{\partial B_r}{\partial z} - \frac{\partial B_z}{\partial \phi} \right). \tag{8}
\]

\[
\frac{\partial}{\partial r} \left[ r (v_r B_\phi - v_z B_r) \right] = 0. \tag{9}
\]

Here, the induction equation is written in the limit of very large magnetic Reynolds number, which is the case for an accretion disc.

2.2 Thermodynamics of the gas: in the presence of radiation and magnetic field

The first law of thermodynamics allows us to calculate the entropy gradient in terms of temperature and density gradients. This entropy gradient changes in different types of accretion processes depending on the detailed balance of heating, cooling, and advection. Different accretion rate, and hence matter density, provides the information of cooling mechanisms. Also magnetic field plays an important role in heating, as well as, cooling processes. Hence, the equation of state for a mixture of perfect gas and radiation in the presence of magnetic field is

\[
\rho_T = \rho_p + \rho_m = \rho_p + \rho_r + \rho_m = \frac{k_b \mu T}{\mu m_p} - \frac{1}{3} \frac{a T^4}{8 \pi} + \frac{B^2}{8 \pi}. \tag{10}
\]
where \( p_i \) is the total pressure, \( k_B \) the Boltzmann constant, \( \mu \) the mean molecular weight, \( m_p \) the proton mass and \( a \) the Stefan constant. The three different pressures: gas (\( p_g \)), radiation (\( p_r \)) and magnetic (\( p_m \)) can be written in terms of two fundamental parameters \( \beta \) and plasma-\( \beta \) (\( \beta_m \)) given by

\[
p_g = \beta p \text{ , } p_r = (1 - \beta) p \text{ , } p_m = (\beta/\beta_m)p \, .
\]  

(12)

Throughout we take the parameter \( \beta = (p_g/p) \) to be independent of \( r \), unlike the parameter \( \beta_m = (p_g/p_m) \). Note that we keep the same definition of \( \beta \) as discussed in the context of gas-radiation mixture (e.g., Clayton 1983).

The internal energy per unit mass of the system, in the presence of magnetic field, is

\[
U = \frac{3}{2} \frac{pk_BT}{\mu m_p} V + dT^4 V + \frac{B^2}{4\pi} V,
\]

(13)

where \( V \) is the volume of unit mass of gas, and hence \( \rho V = 1 \).

Using first law of thermodynamics and flux-freezing assumption, the entropy gradient is then given by

\[
TdS = \frac{p}{\rho} \left( \frac{12 - 21\beta}{3\beta} \right) \left( \frac{dT}{T} \right) - 4 - 3\beta + \frac{1}{\beta_m} \left( \frac{dp}{p} \right).
\]

(14)

Following Chandrasekhar (1967), we can define the adiabatic exponents \( \Gamma_1 \) and \( \Gamma_3 \) by the equations

\[
\frac{dp}{p} + \frac{dV}{V} = 0 \text{ and } \frac{dT}{T} + \left( \frac{\Gamma_3 - 1}{\Gamma_3} \right) \frac{dV}{V} = 0.
\]

(15)

For such gas-radiation mixture in the presence of magnetic fields, these exponents are

\[
\Gamma_1 = \frac{32}{24 - 2\beta}, \quad \Gamma_3 = \frac{32}{24 - 2\beta} - 2\beta/\beta_m.
\]

(16)

In the absence of magnetic fields (\( \beta_m \rightarrow \infty \)), the entropy gradient and \( \Gamma_1 \) and \( \Gamma_3 \) will reduce to their respective hydrodynamical forms as given by Abramowicz et al. (1988) and Clayton (1983) respectively.

### 2.3 Radiation mechanisms for two-temperature plasma

For simplicity, we assume the gas consists of ions and electrons. From charge neutrality, the number density of ions and electrons are equal. However, the plasma in this advective paradigm behaves like a two-temperature system due to their large mass difference. Here we allow the electron temperature \( T_e \) and ion temperature \( T_i \) to be different and hence the gas pressure of the accreting gas can be written as

\[
p_g = \beta p = \frac{\rho k_B T_e}{\mu m_p} + \frac{\rho k_B T_i}{\mu e m_p}.
\]

(17)

where \( \mu_i \) and \( \mu_e \) are the effective molecular weights for ions and electrons respectively. Since ions are much heavier than electrons, we normally expect all the generated heats due to viscous and/or magnetic dissipations primarily act on ions. Some part of this heat may transfer from ions to electrons via Coulomb coupling. Finally electrons take part in radiating heat. The separate energy equations for ions and electrons are considered by taking detailed balance of heating, cooling and advection. The energy equation for ions is given by

\[
\frac{24 - 21\beta}{2(4 - 3\beta)} \left[ v_r \left( \frac{\partial p}{\partial r} - \Gamma_1 \frac{p}{\rho} \frac{\partial p}{\partial r} \right) + v_z \left( \frac{\partial p}{\partial z} - \Gamma_1 \frac{p}{\rho} \frac{\partial p}{\partial z} \right) \right]
\]

\[
= Q^+ - Q_i. \quad (18)
\]

Here, the radial and vertical advection of the ion flows are represented by the first and second terms of the left-hand side of the equation (18). The right-hand side of equation (18) represents the difference between the rate of energy gained (generated via dissipation process, \( Q^+ \)) and loosed (transfer from ions to electrons, \( Q_i \)) by ions per unit volume. \( Q^+ \) consists both the viscous and magnetic dissipation parts as \( Q^+ = Q^+_{vis} + Q^+_{mag} \). The rate of generated heat per unit volume due to viscous dissipation is given by

\[
Q^+_{vis} = a \left( p + \rho v_r^2 \right) \left( \frac{1}{r} \frac{\partial v_r}{\partial r} - \frac{1}{r} \frac{\partial v_z}{\partial r} + \frac{z}{r} \frac{\partial v_z}{\partial r} + \frac{v_r}{\rho} \frac{\partial v_z}{\partial r} \right). \quad (19)
\]

On the other hand, the magnetic heating is basically due to the abundant supply of magnetic energy and/or due to the annihilation of the magnetic fields. The rate of generated heat per unit volume due to magnetic dissipation is given by (Mondal & Mukhopadhyay 2019a)

\[
Q^+_{mag} = \frac{1}{4\pi} \left( B_r B_z \left( \frac{\partial v_r}{\partial r} + \frac{\partial v_z}{\partial r} \right) + B_{\phi} B_r \left( \frac{1}{r} \frac{\partial \phi}{\partial r} + \frac{21}{r^2} \right) \right).
\]

(20)

The energy transfer from ions to electrons happens through Coulomb coupling. This term behaves like heating term in the energy balance equation for electrons. The volume transfer rate of energy from ions to electrons is given by (Bisnovatyi-Kogan & Lovelace 2000)

\[
q_{ie} = \frac{8\sqrt{2\pi} e^4 n_i n_e}{m_i m_e} \left( \frac{T_e}{m_e} + \frac{T_i}{m_i} \right)^{-3/2} \ln(\Lambda)(T_i - T_e)
\]

\[ \text{ergs cm}^{-3} \text{ s}^{-1}, \quad (21) \]

where \( \ln(\Lambda) \approx 20 \) is the Coulomb logarithm. Here the dimensionless form \( (Q^+)^{\ast} \) is linked through \( q_{ie}^{\ast} = Q^+^{\ast}c_1^{11}/(G^4 M_B^2 H) \).

The energy equation for electrons is then given by

\[
\frac{24 - 21\beta}{2(4 - 3\beta)} \left[ v_r \left( \frac{\partial p_e}{\partial r} - \Gamma_1 \frac{p_e}{\rho} \frac{\partial p_e}{\partial r} \right) + v_z \left( \frac{\partial p_e}{\partial z} - \Gamma_1 \frac{p_e}{\rho} \frac{\partial p_e}{\partial z} \right) \right]
\]

\[
= Q_e - Q^-. \quad (22)
\]

where \( Q^- \) represents the radiative cooling rate through electrons via different cooling processes including bremsstrahlung \( q_{br} \), synchrotron \( q_{syn} \). The radiative cooling rate per unit volume is

\[
q^- = Q^- c_1^{11}/(G^4 M_B^2 H) = q_{br} + q_{syn}.
\]

Various cooling formalisms are adopted from Narayan & Yi (1995) and Rajesh & Mukhopadhyay (2010) and can be written as

\[
q_{br} = 1.4 \times 10^{-27} n_i n_e T_e^{1/2} (1 + 4.4 \times 10^{-10} T_e), \quad \text{ergs cm}^{-3} \text{ s}^{-1}, \quad (23)
\]
\[ \dot{q}_{\text{syn}} = \frac{2\pi}{3} \frac{kT_e}{c^3} \frac{v_c^3}{R} \text{ ergs cm}^{-3} \text{ s}^{-1}, \]  
where \( v_c \) is the synchrotron self-absorption cut-off frequency. This can be determined from the relation

\[ v_c = \frac{3}{2} \frac{\nu_0 \theta_0^2}{\lambda_{\text{syn}}}, \quad \nu_0 = \frac{eB}{2\pi mc^2}, \quad \theta_0 = \frac{k_BT_e}{m_ec^2}. \]  

The parameter \( \lambda_{\text{syn}} \) is computed numerically at every radius \( R = rGM_{\text{BH}}/c^2 \) following Narayan & Yi (1995).

### 3 SOLUTION PROCEDURE

In this disc-outflow coupled region, we make a reasonable assumption that within the inflow regime, the vertical variation of any dynamical flow parameters (say, \( A \)) is much less than that with radial variation, which allows us to choose \( \delta A/\delta z \approx sA/2z \), where the constant \( s \) is scale parameter for that corresponding variable. The magnitude of \( s \) is very small compared to unity \((|s| \sim 0.01)\). Note that the BH accretion is transonic. The sub-sonic matter far away from the BH passes through sonic/critical point and becomes super-sonic.

Following standard practice (e.g. Rajesh & Mukhopadhyay 2010), we combine all the above equations to express \( dv_r/dr \) in terms of all dynamical variables and independent variable \( r \), as

\[ \frac{dv_r}{dr} = \frac{N}{D}. \]  

At `critical radius` \( r = r_c, D \) becomes zero. To capture smooth solutions around such point, \( N \) must be vanished therein. Any variables with subscript `c` refers to the values of that respective variables at the critical radius. Here, we prescribe

\[ B_{ic} = \sqrt{4\pi\rho_c} \frac{c_{sc}}{f_i}, \quad \text{and} \quad v_{zc} = \frac{c_{sc}}{f_i}. \]  

where we assume \( c_{sc} = \sqrt{\frac{\mu_0\rho_c}{\rho}}, \) and the constants \( f_i \) and \( f_j \) determine the magnetic field strength and outflow velocity at that critical radius respectively.

The matter density and the outflow velocity at the critical point are prescribed through the following information. Integrating the continuity equation, we obtain the total mass accretion rate \( M \), as given by

\[ \int_0^{r_c} \int_0^{2\pi} \int_{-h}^h \left[ \frac{\partial}{\partial r} (\rho v_r) + \frac{\partial}{\partial z} (\rho v_z) \right] \, dr \, dz \, d\theta = M. \]  

The first term on the left-hand side of the equation (28) signifies the rate at which the radial mass flux changes, whereas the second term indicates the vertical mass flux rate. Hence this total mass accretion rate can decouple into inflow rate \( M_0 \) and outflow rate \( M_j \) to prescribe the information of vertical velocity and mass density at sonic location in this disc-outflow symbiotic model, as \( M = M_0 + M_j \). Therefore, the inflow and outflow rate can read as

\[ M_0 = \int_0^h 4\pi \rho v_r \, dz \quad \text{and} \quad M_j = \int 4\pi \rho v_z \, dr + c_j, \]  

where the constant \( c_j \) is determined from appropriate boundary condition. Throughout in our computation, we express the mass accretion rate in units of the Eddington rate: \( \dot{M}_{\text{Edd}} = L_{\text{Edd}}/(\eta c^2) = 1.39 \times 10^{38} (M/M_\odot) \, \text{g s}^{-1} \), where \( \eta = 0.1 \) here is the radiative efficiency factor.

Now at \( r = r_c, \, dv_r/dr = 0/0 \). Applying the Hospital rule and some algebraic simplification, we find that the velocity gradient at critical point has two values. These two values define the nature of the critical point (Chakrabarti 1990; Mondal & Mukhopadhyay 2018). When both are complex, the critical point is `spiral'-type. When both are real and of same sign, the critical point is `nodal'-type. When both are real and of opposite sign, the critical point is `saddle'-type. The negative slope at `saddle'-type critical point indicates the accretion solution, whereas the positive one refers wind solution. Due to two-temperature prescription in this advective paradigm, we need to supply the electron temperature \( T_{ec} \) and specific angular momentum \( J_c \) at \( r_c \). We need to adjust these parameters along with the relative dependence of magnetic field geometries, to capture the full dynamical solutions connecting the outer boundary to the BH event horizon through \( r_c \). The simultaneous solution for \( N = 0, \, D = 0 \), and, the equation for vertical momentum balance (equation (5)) provide self-consistently the numerical values for \( z_c, c_{sc} \) and \( v_{zc} \) at \( r = r_c \). The outer boundary corresponds to the radius \( r = r_{out} \), where \( J = J_K, J_K \) is the Keplerian angular momentum per unit mass of the flow where the centrifugal force balances the gravity. The inner boundary corresponds to the event horizon of the BH, in which radius the velocity becomes the order of speed of light (in our units unity). We also have to supply \( M_{BH}, \, M, \, \alpha, \) and \( \beta \).

### 4 COUPLED DISC-OUTFLOW PHYSICS

#### 4.1 Disc dynamics: stellar mass black holes

We address here how the large-scale strong magnetic fields can influence the transport of angular momentum, as well as, outflow dynamics in the advective paradigm. The disc is here quasi-spherical and hot puffed-up. In Fig. 1, we show how the fundamental flow parameters vary throughout this disc-outflow coupled system. The location of the critical point and specific angular momentum value at such critical location are respectively \( r_c = 6.0, \, \lambda_c = 3.184 \) for non-rotating BH \((a = 0)\), and \( r_c = 4.2, \, \lambda_c = 1.385 \) for fast-rotating BH \((a = 0.998)\). The typical values for vertical scale parameter `\( s \)' for the corresponding flow variables \( v_r, \lambda, v_z, B_r, B_\phi, B_z, B, \rho \), and \( \rho \) are respectively \( s_1 = -0.014, \, s_2 = -0.01, \, s_3 = 0.04, \, s_4 = -0.03842, \, s_5 = -0.03, \, s_6 = 0.01, \, s_7 = -0.035, \) and \( s_8 = -0.03 \). The magnetic field configuration at the critical location is \( B_{0c} = 1.8B_{Ke}, \quad B_{zc} = -B_{rc}/5 \) when \( a = 0, \) and \( B_{0c} = 0.6B_{ke}, \quad B_{zc} = -B_{rc}/5 \) when \( a = 0.998 \).

Fig. 1(a) describes the Mach number (\( M \)) profile, which is defined as the ratio of the radial velocity to the sound speed. It indicates that very far away from the BH, matter is sub-sonic and is independent of BH’s spin. As it advances towards the central BH, it becomes super-sonic near the event horizon of a BH. The sonic locations for non-spinning and fast-spinning BHs are respectively \( r = 6.6269 \) and \( r = 5.049 \). The inner boundary of our solution is defined as the location where effective matter velocity \( v = \sqrt{v_r^2 + (A/r)^2 + v_z^2} \) reaches the speed of light. This location does not overlap exactly with the corresponding location of the event horizon of a BH, because of the use of pseudo-Newtonian potential.
Figure 1. Variations of (a) Mach number, (b) specific angular momentum, (c) vertical/outflow velocity, and (d) sound speed, as functions of radial coordinate. The model parameters are $M = 20M_\odot$, $m = 0.05$, and $\alpha = 0.01$.

instead of full GR computation. Fig. 1(b) describes the outward transport of the angular momentum in this flow. Very far away from the BH, the transition radius between the Keplerian and sub-Keplerian flows is the outer boundary for our solutions. The Keplerian angular momentum basically signifies the profile in which the centrifugal force balances the gravitation force of the BH. As matter falls towards the central BH, it loses angular momentum to form the disc. However, it is very difficult to model self-consistently the transition region where $\lambda/\lambda_K = 1$. This is because, the set of equations used to model such hot advective flow is not strictly continued to be valid to explain cold, optically thick, Keplerian disc model. Note that, in this paper we do not intend to address this transition zone, rather we concentrate on the hot advective part. In Fig. 1(c), we show the profile for the vertical/outflow velocity. At outer zone, it is almost negligible as usual. The outflow is basically emitted from the hot puffed-up region of this advective disc-outflow surface. Our model is valid vertically up to the upper surface of this region, above which flow will decouple and accelerate to form plausible jets. The vertical velocity near vicinity of BH increases from 0.1053 when $a = 0$ to 0.2512 when $a = 0.998$. The corresponding sound speed profile is shown in Fig. 1(d). The medium sound speed at the horizon increases from 0.2352 when $a = 0$ to 0.4637 when $a = 0.998$. The increase of $c_s$ with spin indicates that the temperature of the system is higher for fast-spinning BHs. All such changes suggest the possibility that the energetics of such advective flows may strongly depend on BH spin.

The details of the magnetic field geometry and different stress components are shown in Fig. 2. Figs. 2(a) and (b) indicate the variations of different field components for
non-rotating and rapidly rotating BHs respectively. Spinning BHs can sustain more magnetic fields in comparison with non-rotating BHs. The magnetic field strength near the event horizon is of the order of $|B| \sim 8.74 \times 10^6$ G when $a = 0$, and $|B| \sim 2.67 \times 10^7$ G when $a = 0.998$ for a $20 M_\odot$ BH. Figs. 2(c) and (d) show the relative magnitude of different magnetic field components with respect to the total field for $a = 0$ and $a = 0.998$ respectively. Initially, very far away from the BH, the disc is mainly vertical poloidal field dominated. As matter drags inward, the advection of both poloidal and toroidal fields happens. Within the plunging region, near to the central BH, the enhancement of the
magnetic field becomes more. This is because of the presence of strong gravity and the absence of differential rotation in this zone. Figs. 2(e) and (f) show the ratio of the different components of magnetic to viscous stresses for \( a = 0 \) and \( a = 0.998 \) respectively. Two main components, \( r\phi \) and \( \phi_z \), reveal the radial and vertical transports respectively. These profiles signify the dominant nature of the magnetic stress over viscous one and this should be the case in any kind of magnetically dominated accretion flows. The plasma-\( \beta \) parameter for both non-spinning and rapidly spinning BHs is shown in Fig. 3. The rapidly spinning BHs can sustain more magnetic fields compared to the non-spinning BHs. It makes the value of the plasma-\( \beta \) parameter smaller for spinning BHs. Also such low \( \beta_m \) value (near or below 10) infers strongly magnetized flows for BH accretion. The nature of the magnetic field vectors is visualized in 2D-plane \((x - y)\) for a rapidly-spinning BH, as shown in Fig. 4. The arrow size is normalized here to scale the field value accordingly, with Cartesian-coordinate value and the colorbar indicates the actual magnetic field strength near such rotating stellar-mass BHs.

### 4.2 Disc dynamics: supermassive black holes

The solutions for magnetized disc-outflow symbiosis around supermassive BHs with mass \( 10^{8}M_\odot \) keeping other model parameters exactly same as in the case of stellar-mass BHs are shown in Fig. 5. The flow parameters, like, Mach number, specific angular momentum, outflow velocity, and medium sound speed, reveal same behaviour with that of stellar-mass BHs as shown in Fig. 1, hence are not depicted again. This is because, the magnetized disc-outflow symbiotic model is effectively scale-free when the physical variables are written in terms of relevant fiducial parameters (Schwarzschild radius, Eddington accretion rate, light velocity etc.). It makes some quantities, like, radial and vertical velocities, specific angular momentum, sound speed are essentially independent of BH mass for such advective flows. However, not all the features are similar. Some physical quantities, like, density, magnetic field strength, pressure etc., differ quite largely with the mass of the BH, though their profiles appear similar. The variations of different magnetic field components for non-rotating, as well as, fast-rotating supermassive BHs are shown in Fig. 5(a) and (b) respectively. Since the magnetic field strength \( (B) \) varies as \((M/M_\odot)^{-1/2}\), the \( B \) value drops here in comparison with stellar mass BHs as shown in Fig. 2. Near the event horizon of the BH, the magnetic field strength is of the order of \( |B| \sim 3.91 \times 10^3 \) G when \( a = 0 \), and \( |B| \sim 1.16 \times 10^4 \) G when \( a = 0.998 \), for a \( 10^8 \) \( M_\odot \) BH. The relative magnitude of different magnetic field components with respect to the total field for \( a = 0 \) and \( a = 0.998 \) are same in comparison with the respective stellar-mass BHs. This is because of the scale-free nature with respect to the mass of the BH. The stress ratios of different components of magnetic to viscous ones also show similar behaviour as in case of stellar-mass BHs. This is expected, because density and pressure vary as \((M/M_\odot)^{-1}\), and hence stress ratios remain independent of BH mass.

### 4.3 Energetics of the magnetized accretion process

We compute the energetics of the magnetized accretion induced outflow for non-rotating, as well as, rapidly rotating BHs spanning from stellar-mass to supermassive scales. The disc luminosity can be computed from the cooling mechanisms and can be defined as

\[
L = \frac{h}{\int_{0}^{a} Q^{-4\pi r} \, dz} \, dr.
\]  

The variation of disc luminosity for stellar mass BH is
shown in Fig. 6. The luminosity at any arbitrary $r$ is computed by integrating (as given in equation 30) from outer disc radius to that corresponding $r$. The total luminosity is basically integration over whole disc from outer radius to inner one and, hence, the total luminosity is the value corresponding to that at inner radius. Figs. 6(a) and (c) show the variations of luminosity for $a = 0$ and $a = 0.998$ respectively. For stellar-mass BH of mass $M_{BH} = 20 M_\odot$ with total mass accretion rate $\dot{m} = 0.05$, the maximum attainable disc luminosity integrated over whole disc is $L \sim 1.1389 \times 10^{39}$ erg s$^{-1}$ when $a = 0$ and $L \sim 6.4968 \times 10^{39}$ erg s$^{-1}$ when $a = 0.998$. One question may automatically arise: how can such large luminosity possible? The answer is the presence of externally generated large scale strong magnetic field and its configurations. Figs. 6(b) and (d) show the variations of magnetic field strength for $a = 0$ and $a = 0.998$ respectively. The red dotted lines show the variation of the Eddington magnetic field limit ($B_{Edd}$) for $20 M_\odot$ BH. The estimate of such $B_{Edd}$ is based on the simple assumption that the luminosity associated with the magnetic energy density is comparable to the Eddington luminosity, given by equation (1). This can be expressed as

$$B_{Edd} = \frac{6.15 \times 10^6}{r} \left( \frac{M}{M_\odot} \right)^{-1/2} \text{ G}.$$  \hspace{1cm} (31)

The magnetic field near the outer disc region is large enough and even more than the corresponding Eddington limit $B_{Edd}$ at that zone. Near the BH, the magnetic field maintains its Eddington limit. Such field configurations enhance the synchrotron cooling process and, hence, help to achieve very large luminosity. If magnetic field strength drops due to less supply of magnetic field, and goes down to the corresponding $B_{Edd}$ limit throughout the flow, then synchrotron cooling also reduces accordingly. That accretion environment certainly can not achieve very large luminosity. This is shown in Fig. 7 for the case of non-rotating stellar-mass BHs with mass $20 M_\odot$, keeping other parameters same as shown in Fig. 6. The maximum attainable disc luminosity integrated over the whole disc is $L \sim 4.3689 \times 10^{38}$ erg s$^{-1}$. Fig. 7(b) shows the variation of the magnetic field strength along with its corresponding Eddington limit. Near the event horizon of the BH, the magnetic field strength is of the order of $|B| \sim 6.2288 \times 10^5$ G, and it is one order less than the case as shown in Fig. 6. The field strength here is quite below than the corresponding $B_{Edd}$ limit throughout the flow. Hence the magnetic field strength is playing very important role to achieve very large luminosity.

The accretion disc luminosity and the corresponding magnetic field strength for a supermassive BH of mass $M_{BH} = 10^9 M_\odot$ with $\dot{m} = 0.05$ are shown in Fig. 8. The luminosity is strongly BH mass dependent. For instance, the synchrotron emission is self-absorbed below the critical frequency $\nu_c$ and it is radiated at $\nu_c$ which is proportional to the magnetic field strength $B$ (equation 25) and hence this critical frequency itself varies as $(M/M_\odot)^{-1/2}$. The maximum attainable disc luminosity integrated over whole disc is $L \sim 3.0421 \times 10^{35}$ erg s$^{-1}$ when $a = 0$ and $L \sim 1.2345 \times 10^{36}$ erg s$^{-1}$ when $a = 0.998$. This is as usually because of the fact that the rapidly spinning BHs can sustain more magnetic field compared to non-rotating BHs and therefore enhance cooling process through synchrotron radiation.

5 OBSERVATIONAL IMPLEMENTATIONS

Galactic BH X-ray binaries show a large varieties in their X-ray spectral states, namely, quiescent, low/hard, intermediate, high/soft, and very high (Remillard & McClintock 2006; Fender et al. 2004). It is believed that these different spectral states correspond to different accretion geometries, as mentioned before. In this paper we focus on the hard-state BH sources with BH mass ranging from stellar mass to su-

![Figure 5. Variations of magnetic field components for (a) $a = 0$, and (b) $a = 0.998$, as functions of radial coordinate. The model parameters are $M = 10^3 M_\odot$, $\dot{m} = 0.05$, and, $a = 0.01$.](image-url)
permassive scales. The so-called hard states are generally observed at low X-ray luminosities ($L < 0.01L_{\text{Edd}}$) and their X-ray spectra are well explained by a power-law component with photon index $\Gamma \sim 1.4$ to 1.8. However, the observations of hard state BH sources often reach higher luminosities. For example, the BHC GX 339-4 achieves luminosity up to 30%$L_{\text{Edd}}$ in its hard state (Zdziarski et al. 2004). Apart from such BH binaries, a large fraction of ULXs are even more luminous in their hard power-law dominated states. Some of these ULXs are listed in Table 1. The true nature of such observations remains mysterious over the decades. Here we address that our highly magnetized, advective, optically thin disc-outflow symbiotic model can achieve these large luminosities. Fig. 6 indicates that maximum attainable luminosity integrated over the disc is $L \sim 6.4968 \times 10^{39}$ erg s$^{-1}$ for rapidly spinning stellar mass BH of mass $M_{\text{BH}} = 20 M_\odot$ with mass accretion rate $\dot{m} = 0.05$. Hence such magnetically dominated advective accretion process can easily explain these long standing issues.

On the other hand, the observations of supermassive BHs also show unusual large luminosities in their hard spectral states. For example, some high synchrotron peak (HSP) BL Lac objects appear very luminous in their hard power-law dominated states. Some of these sources are listed in Table 2 based on the observations of Fermi second catalog of active galactic nuclei (Ackermann et al. 2011). As shown in Fig. 8 that our magnetically dominated disc-outflow symbiotic model can reach luminosity $L \sim 1.2345 \times 10^{46}$ erg s$^{-1}$ for rapidly spinning supermassive BH of mass $M_{\text{BH}} = 10^8 M_\odot$ with mass accretion rate $\dot{m} = 0.05$. 

Figure 6. Variations of (a) disc luminosity, and (b) magnetic field, for non-rotating stellar mass BH. The black solid line represents the magnetic field profile of the disc, whereas red dotted line represents the corresponding Eddington magnetic field limit. (c), and (d) depict the same as (a), and (b) respectively, except for $a = 0.998$. The model parameters are same as in Fig 1.
Variations of Figure 7. Table 1. Some hard-state ULX sources.

| Source      | $\Gamma$ | $L_{0.3-10 \text{ keV}}$ ($10^{38}$ erg s$^{-1}$) | Ref. |
|-------------|----------|---------------------------------|------|
| NGC 1365 X1 | 1.74$^{+0.12}_{-0.11}$ | 2.8 | 1 |
| NGC 1365 X2 | 1.23$^{+0.07}_{-0.06}$  | 3.9 | 2 |
| Holmberg IX X-1 | 2.0$^{+0.1}_{-0.02}$ | 1.0 | 2 |
| NGC 5775 X1 | 1.81$^{+0.02}_{-0.02}$ | 7.5 | 3 |
| NGC 3628 X1 | 1.81$^{+0.02}_{-0.02}$ | 7.5 | 3 |
| M99 X1      | 1.75$^{+0.01}_{-0.01}$ | 1.9 | 5 |
| M82 X42.3+59| 1.44$^{+0.09}_{-0.09}$ | 1.1 | 6 |
| Antennae X-11 | 1.76$^{+0.05}_{-0.05}$ | 2.1 | 6 |
| Antennae X-16 | 1.35$^{+0.03}_{-0.04}$ | 1.8 | 7 |
| Antennae X-42 | 1.73$^{+0.10}_{-0.11}$ | 0.96 | 7 |
| Antennae X-44 | 1.74$^{+0.04}_{-0.04}$ | 1.28 | 7 |

Table 2. BL Lac objects in a hard power-law dominated state based on the observations of Fermi second catalog of active galactic nuclei (Ackermann et al. 2011).

| Source      | $\Gamma$ | $L_{1-100 \text{ GeV}}$ ($10^{38}$ erg s$^{-1}$) | Ref. |
|-------------|----------|---------------------------------|------|
| PKS 0301-243 | 1.975 | 1.938 | 1.10 |
| PKS 0447-439 | 1.975 | 1.855 | 1.21 |
| PKS 0011+393 | 1.975 | 1.489 | 1.83 |
| PKS 0301-243 | 1.975 | 1.23 | 1.05 |
| B3 1307+433 | 1.975 | 1.23 | 1.05 |
| B3 1307+433 | 1.975 | 1.23 | 1.05 |

6 DISCUSSION

In this 2.5–dimensional magnetized, viscous, advective disc-outflow symbiotic model, we address the role of large-scale strong magnetic fields in the formation of strong outflows and in the enhancement of synchrotron cooling. Below we discuss some important aspects of this framework.

First, what could be the possible origin of large-scale strong magnetic fields in an accretion disc? The accretion disc problems have been studied widely in the presence of small-scale fields, as well as, large-scale fields. Small-scale field may be generated locally. Some seed magnetic fields can generate from zero initial field condition via Biermann battery mechanism (Biermann 1950). To operate this mechanism, a non-parallel gradient of density and temperature profiles is required, which is very common in accretion environment. Also, the MHD dynamo process may generate small-scale magnetic field locally (Brandenburg et al. 1995). However, the origin of large-scale magnetic fields is not clear yet. One possibility may be the externally generated fields. The interstellar medium or the companion star may supply...
such fields, which further push towards the central BH by the continuous accretion pressure and may enhance due to flux freezing (Bisnovatyi-Kogan & Ruzmaikin 1974).

Second, is there any upper bound to the amount of magnetic flux to thread the disc and/or BH? Any magnetized accretion flow causes certain amount of magnetic flux to thread the disc and/or BH. Recent observations of radio-loud active galaxies confirm a dynamically dominated magnetic fields in the jet launching region based on the correlation of jet magnetic field and accretion disc luminosity (Zamaninasab et al. 2014). Also, the unusual large Faraday rotation near the centre of our Galaxy infers a signature of strong magnetic field near the BH (Eatough et al. 2013). GRMHD simulations for relativistic jets generally assume highly magnetized plasma at the jet foot-print (McKinney & Gammie 2004; Tchekhovskoy et al. 2011). Theoretical models over the decades try to correlate the observable quantities with the fundamental properties of the disc and BH. Based on the combined effects of Blandford-Payne (Blandford & Payne 1982) and Blandford-Znajek (Blandford & Znajek 1977) mechanisms, the jet kinetic power had been computed in terms of the mass and spin of the BH and the magnetic field strength in the vicinity of BH (Garofalo et al. 2010). From an estimate of the kinetic power of the relativistic jet (Russell et al. 2013), the measurement of the synchrotron cooling time (Baczko et al. 2016) and the observed characteristic frequencies of quasi-periodic oscillations of radiation (Piotrovich et al. 2011), it was suggested that the typical values of the magnetic field near the event horizon is $B \sim 10^8$ G for stellar mass BHs and $B \sim 10^4$ G for supermassive BHs.

Figure 8. Variations of (a) disc luminosity, and (b) magnetic field profile for non-rotating supermassive BHs. The black solid line represents the magnetic field profile of the disc, whereas red dotted line represents the corresponding Eddington magnetic field limit. (c), and (d) depict the same as (a), and (b) respectively, except for $a=0.998$. The model parameters are same as in Fig 1, except $M = 10^8 M_\odot$. 

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In this context, the Eddington magnetic field limit near the
event horizon of a BH, $B_{\text{Edd}} = 10^4 \ G \left( \frac{M}{10 M_\odot} \right)^{1/2}$, might
be the approximate upper bound to the amount of magnetic
field strength what any disc around a BH can sustain (Mondal & Mukhopadhyay 2019a), which limit is perfectly viable
in our computation. More importantly, this upper limit in
our model sets up automatically through the existence of
inner `saddle'-type critical point as discussed before.

Apart from these, our main venture is to explain some
unusual observational key features based on this magnetized
accretion process. One such object is ULXs, which are very
bright, off-nuclear X-ray sources with luminosity exceeding
the standard Eddington limit for a stellar mass black hole.
The theoretical explanation for their large apparent luminos-
ities is either the existence of the missing class of inter-
mediate mass black hole (IMBH; Miller et al. 2004), or the
super-Eddington accretion around a stellar-mass BH (Ebisu-
awa et al. 2003), or beamed emission (King et al. 2001).
Indeed, it is still an open question regarding the existence
of the missing class of IMBH and also the supporting evidence
for IMBH scenario has been disputed extensively (Gonça-
elves & Soria 2006). Most arguments support the idea that a ma-
jor fraction of ULXs is stellar mass BH (Begelman 2002; Motch et al. 2014). Super-Eddington scenario may explain
the ULXs with a steep power law. However, the hidden na-
ture of a large fraction of ULXs with hard-power law state
remains mysterious. Exactly in this point, our model comes.
The magnetically dominated advective disc-outflow symbio-
sis around a stellar mass BH can reach such large luminos-
ity even for sub-Eddington mass accretion rate. The most
important criterion is the magnetic field strength, as well as,
its geometry. Near the transition region from Keplerian
to sub-Keplerian flow, the magnetic field strength is quite
larger than the corresponding Eddington limit. This field
becomes dynamically dominant near the BH due to contin-
uous advection of the magnetic flux through flux freezing.
Near the BH, magnetic field well maintains its correspond-
ing Eddington limit. This is a very efficient way to enhance
the synchrotron cooling to reach such a large luminosity.
Also, the large toroidal fields exert a huge outward pressure
to produce strong outflows in this model. Unlike other mag-
netically dominated accretion model, here the advection
of both poloidal and toroidal field components takes place.

7 CONCLUSIONS

We have obtained the semi-analytical solutions for a mag-
netized, advective, optically thin disc-outflow symbiosis in-
corporating explicit cooling formalisms. The detailed bal-
ance between heating, cooling and advection is taken care
here. We prescribe the generalized viscous shearing stress
in terms of standard $\alpha$-viscosity parameter. However, the
components for magnetic shear are at least one order larger
than that of viscous shear in this formalism. The large scale
strong magnetic field removes the angular momentum from
the in-falling matter and also helps in the formation of strong
outflows. We assume that the energy transfer from ions to
electrons occurs through Coulomb coupling and the radia-
tive cooling processes are bremsstrahlung emission and syn-
chrotron emission. The presence of strong magnetic fields
enhances the synchrotron cooling via hot electrons. We im-
plement our model to explain the observed unusually large
luminosity of BH sources of mass ranging from stellar mass
to supermassive scale in their hard power-law dominated
states. Particularly we focus on ultra-luminous X-ray sources
(listed in Table 1), BH binaries with luminosities up to
$\sim 10\%$ to $30\% L_{\text{Edd}}$ in their hard state (e.g. GX 339-4), ultra-
luminous quasars (e.g. PKS 0743-67), highly luminous BL
Lac objects (listed in Table 2). We propose here for the
first time in the literature that such bright HSP-BL Lacs
with high spectral signature are magnetically powered sub-
Eddington, advective accretors around supermassive BHs.
In a similar framework for stellar mass scale, we suggest that
ULXs are magnetically dominated advective, sub-Eddington
accretors around rapidly-rotating stellar mass BHs neither
incorporating the existence of intermediate mass BHs con-
cept, nor with the super-Eddington accretion phenomenon.

Apart from such important aspects of this strongly mag-
netized advective accretion process, we would like to ex-
amine some poorly understood, long standing questions in
this field in future. First, throughout in our computation,
we assume that the energy transfer from ions to electrons
is happening due to Coulomb coupling. How do the other
non-thermal processes influence the results in such magneti-
cally dominated advective accretion phenomena? Second, in
our disc-outflow model, we use pseudo-Newtonian potential
which might not be suitable to extract the rotational energy
of a BH. We plan to explore such magnetically dominated
advective accretion flows in the framework of GRMHD for-
malism in future. Third, how does electron and ion ener-
gization occur in such optically thin, strongly magnetized,
advective accretion process? Very recently, Zh Dankin et al.
(2019) and Schekochihin et al. (2019) initiated to under-
stand the electron and ion heating, coupling and also the
non-thermal particle acceleration mechanism in the case of
radiatively inefficient accretion flows around BHs. Fourth,
what are the plausible mechanisms to generate large-scale
magnetic fields locally in this advective accretion flows
and what is the strength of such generated magnetic fields?

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REFERENCES

Abramowicz M. A., Czerny B., Lasota J. P., Szuszkiewicz E.,
1988, ApJ, 332, 646
Ackermann M., et al., 2011, ApJ, 743, 171
Ago J., Krollk J. H., 2000, ApJ, 528, 161
Balasz A.-K., et al., 2016, A&A, 593, A47
Balbus S. A., Hawley J. F., 1991, ApJ, 376, 214
Begelman M. C., 2002, ApJ, 568, L97
Bhattacharya D., Ghosh S., Mukhopadhyay B., 2010, ApJ, 713, 105
Biermann L., 1950, Z. Naturforsch., 5, 65
Bisnovatyi-Kogan G. S., Lovelace R. V. E., 2000, ApJ, 529, 978
Bisnovatyi-Kogan G. S., Ruzmakin A. A., 1974, ApSS, 28, 45
