Hot accretion flow with radiative cooling: state transitions in black hole X-ray binaries

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ABSTRACT

We investigate state transitions in black hole X-ray binaries through different parameters by using two-dimensional axisymmetric hydrodynamical simulation method. For radiative cooling in hot accretion flow, we take into account the bremsstrahlung, synchrotron and synchrotron-self Comptonization self-consistently in the dynamics. Our main result is that the state transitions occur when the accretion rate reaches a critical value \( \dot{M} \sim 3\alpha \dot{M}_{\text{Edd}} \), above which cold and dense clumpy/filamentary structures are formed, embedded within the hot gas. We argued this mode likely corresponds to the proposed two-phase accretion model, which may be responsible for the intermediate state of black hole X-ray binaries. When the accretion rate becomes sufficiently high, the clumpy/filamentary structures gradually merge and settle down onto the mid-plane. Eventually the accretion geometry transforms to a disc-coroana configuration. In summary our results are consistent with the truncated accretion scenario for the state transition.

Key words: accretion, accretion discs – black hole physics – hydrodynamics: HD – ISM: jets and outflow

1 INTRODUCTION

The accretion of matter onto black holes, where a huge amount of energy (both radiative and kinematic) is liberated out, is the key process in black hole X-ray binaries (BHBs) and active galactic nuclei (AGNs). Interestingly, it is argued that both BHBs and AGNs can be divided into two different broad states/types. This is most evident in BHBs, where two states with distinctive spectral and timing properties are identified (see e.g. Zdziarski & Gierlinski 2004; Remillard & McClintock 2006; Belloni 2010 for recent reviews on the state classifications and their observational properties). As the source entering its outburst, generally it will first go through a hard state, which is fainter but with a hard power-law spectrum. The thermal component is highly suppressed in this state. As it brightens, it will enter into a soft state, in which the spectrum is characterized by a thermal component, supplemented with a weak power-law tail. Between these two standard states, there also exist a hybrid one, where the thermal and power-law components are comparable. We call it intermediate state. In AGNs, although less clear, more pieces of evidences are gathered recently to propose that the bright AGNs are analogy to BHBs in their soft states, while the low-luminosity AGNs (LLAGNs) are analogy to BHBs in their hard states (e.g. Ho 2008; Antonucci 2012; Done 2014).

As widely accepted, the basic theoretical picture of these two states is the truncated accretion – jet model (Esin et al. 1997; Yuan et al. 2005; see Yuan & Narayan 2014 for the latest review). In this model, the geometrically thin, optically thick, cold disc (Shakura & Sunyaev 1973; hereafter SSD) is believed to be truncated at a certain radius \( R_* \), inside which it is replaced by a hot accretion flow such as the advection-dominated accretion flows (ADAFs; Narayan & Yi 1994, 1995; Abramowicz et al. 1995) or the luminous hot accretion flow (LHAF; Yuan 2001, 2003; Yuan et al. 2007; Xie & Yuan 2012). Additionally there is also a jet component, which is believed to be connected to the hot accretion flow (Fender et al. 2004; Wu et al. 2013). With this configuration, observationally the system will be in the hard state. As the accretion rate increases, \( R_* \) becomes smaller. Eventually when \( R_* \) is smaller than the innermost stable circular orbit (ISCO), i.e. the whole accretion is now a cold disc and the system will be in the soft state. The readers are referred to Yuan & Narayan (2014) for the most recent
review of the theory of hot accretion flows, including its
dynamics, radiation, and the applications to observations.

Although non-radiative numerical simulations on hot
accretion flows have been investigated for years (hydrody-
namic [HD]; e.g. Stone, Pringle & Begelman 1999; Igumen-
shchev & Abramowicz 1999, 2000; Yuan, Wu & Bu 2012;
magnetohydrodynamic [MHD]; e.g. Stone & Pringle 2001;
Hawley, Balbus & Stone 2001; Igumenshchev, Narayan &
Abramowicz 2003; De Villiers et al. 2003; Machida et al.
2004; Narayan et al. 2012; Yuan, Bu & Wu 2012), only
until very recently the dynamical importance of radiative
cooling on the hot accretion flow has been examined through
numerical simulations (e.g. Machida et al. 2006; Fragile &
Meier 2009; Ohsuga et al. 2009; Yuan & Bu 2010; Ohsuga
& Mineshige 2011; Dibi et al. 2012; Li, Ostriker & Sunyaev
2013; Sadowski & Narayan 2015). Moreover, the theoretical
interpretation of the state transition still lacks direct support
from numerical simulations. The first step, to our knowledge,
is taken by Das & Sharma (2013), where they studied the
effects of radiative cooling on hot accretion flow. They found
that, there exist a certain critical density (depend on the
viscous parameter $a$), above which the whole simulation
domain will be in a globally stable two-zone configuration,
i.e. an outer cold disc and an inner hot accretion flow. Such
geometric configuration has been expected for almost two
decades (e.g. Esin et al. 1997).

Apart from its success, several shortages of their work
could be noted. Firstly, only bremsstrahlung emission is
considered, while in reality the innermost regions ($R <
20–30 \ R_s$, where $R_s = GM_{BH}/c^2$ is the Schwarzschild radius
of the black hole; cf. Fig. 2 and Section 3.2.2) of hot accretion
flow are cooled by synchrotron and its Comptonization
(Narayan & Yi 1995; hereafter NY95). Secondly, they adopt
an optically-thin assumption and the radiative transfer is
not taken into account, which should not be the case
when state transition occurs. Thirdly, an one-temperature
/ion temperature $T_e$ equals electron temperature $T_e$
/ at every location) structure is assumed, which is known to be
/incorrect (e.g. NY95; Moscibrodzka et al. 2009) for most
cases in hot accretion flows.

In this paper, we investigate the dynamical impacts of
radiative cooling in hot accretion flows by using two di-
/ensional radiative HD simulations, aiming at understanding
/the state transitions in BHBs (and the evolution of AGN
/activity). The main improvement of our work is that, besides
/the bremsstrahlung, both the synchrotron radiation and its
/Comptonization are considered, following the treatment of
/NY95. Moreover, the impacts of gas pressure to magnetic
/radial pressure ratio $\beta$, the ratio of electron temperature to
temperature $T_e/T_0$, on the evolutionary stages of accretion
/flows, are also investigated.

Another motivation of this work is to verify, from state
transition simulations, the existence of LHAF (Yuan 2001,
2003; Yuan et al. 2007; Yuan & Bu 2010; Xie & Yuan
2012), a hot accretion flow with accretion rate and radiative
efficiency higher than those of typical ADAFs. This new
solution is promising for the understanding of bright hard
state (Yuan et al. 2007; Ma 2012). Besides, with even
higher accretion rate, analytical investigations of the height-
integrated one-dimensional solutions indicate that thermal
instability could be triggered. Consequently the accretion
flow will be possibly in two-phase (a type II LHAF branch,
see Yuan 2003; Xie & Yuan 2012, 2016), i.e. numerous cold
clumps will be formed, embedded in the hot gas. This two-
phase accretion flow may relate to the intermediate state
in BHBs (Yang et al. 2015). Eventually as the accretion
rate increases further, the cold clumps may grow, merge
and settle down to the mid-plane, resulting the usual SSD.
Together with the low-density hot medium above the cold
disc, the disc-corona configuration is recovered, and the
accretion system enters into a soft state.

The structure of this paper is as follows. In Section 2, we
describe our numerical method and the treatment on
various radiative cooling mechanisms in hot accretion flow.
The main results are described in Section 3. Finally Section
4 is devoted to a brief summary and discussions.

2 METHOD

2.1 Basic hydrodynamic equations

In our numerical simulations, we use the ZEUS-2D code
(Stone et al. 1992) to solve the basic hydrodynamic equations
in spherical coordinates $(R, \theta, \phi)$:

$$\frac{d\rho}{dt} + \rho \nabla \cdot \mathbf{v} = 0, \tag{1}$$

$$\rho \frac{d\mathbf{v}}{dt} = -\nabla P_{\text{gas}} - \rho \nabla \phi + \mathbf{v} \cdot \mathbf{T}, \tag{2}$$

$$\rho \frac{d(e/\rho)}{dt} = -P_{\text{gas}} \mathbf{v} \cdot \mathbf{v} + \mathbf{T}^2/\mu - Q_{\text{rad}}. \tag{3}$$

Here $\rho$ is the mass density, $\mathbf{v}$ is the gas velocity, $P_{\text{rad}} \equiv P_{\text{gas}} + P_{\text{mag}}$ is the total (gas+magnetic) pressure, $\mathbf{T}$ is the anomalous
stress tensor. A pseudo-Newtonian gravitational potential
(Paczynski & Wiita 1980) is adopted, $\phi = -GM_{BH}/(R -
R_s)$. Besides, for an ideal gas, the internal energy $e$ can be
expressed as, $e = P_{\text{gas}}/(\gamma - 1)$, where $\gamma = 5/3$.

The stress tensor $\mathbf{T}$ is now widely believed to be asso-
ciated with MHD turbulence driven by the magneto-
rotational instability (MRI, see Balbus & Hawley 1998 for a
review). However, since our simulation is hydrodynamic,
we follow the conventional strategy, to use the $\alpha$-viscosity
description (Shakura & Sunyaev 1973). Following Stone,
Pringle & Begelman (1999), we include the non-zero az-
imuthal components of $\mathbf{T}$,

$$T_{\theta \theta} = \mu R \frac{\partial}{\partial R} \left( \frac{v_\theta}{R} \right), \tag{4}$$

$$T_{\phi \phi} = \frac{\mu \sin^2 \theta}{R} \frac{\partial}{\partial \theta} \left( \frac{v_\phi}{\sin \theta} \right). \tag{5}$$

Here, the viscous coefficient $\mu = \nu \rho$ and the kinematic
coefficient $\nu \propto R^{1/2}$, which is similar to the standard $\alpha$
description (Shakura & Sunyaev 1973).

Although global large-scale magnetic field may exist in
the hot accretion flow, we here neglect it since our
simulations are hydrodynamic (see Section 5 for discussions

\footnote{Note that, the NY95 approach is also adopted in the general relativistic radiative MHD simulations by Fragile & Meier (2009). However, they assumed $T_e = T_0$, i.e. an one-temperature accretion flow.}
on this point.). The local and randomly-oriented (tangled) magnetic field is determined through the plasma $\beta$ parameter $\beta = P_{\text{mag}}/P_{\text{gas}}$. Consequently the total pressure can be expressed as $P_{\text{tot}} = P_{\text{gas}} + P_{\text{mag}} = (1+\beta) P_{\text{gas}}$. MHD numerical simulations of hot accretion flow indicate that, the value of $\beta$ is not a constant, but varies with time and location $(r, \theta)$ within the simulation domain. Moreover, the average value of $\beta$ likely depends on the magnitude of the net initial magnetic field (e.g. Stone & Pringle 2001; Hawley & Krongold 2001; Hawley et al. 2001; Beckwith et al. 2008. See Yuan & Narayan 2014 for review.). For simplicity we assume $\beta$ is a constant.

The hot accretion flow is two-temperature (NY95); with electron temperature is much lower than ion in the innermost regions of the accretion flow, due to the fact that both the strong radiative cooling of electrons and different adiabatic index between electrons and ions (e.g. Narayan, Mahadaven & Quataert 1998; Yuan & Narayan 2014). We use the following formulae to mimic the results from high-temperature numerical calculations (NY95; Xie et al. 2010),

$$T_e = \frac{T_i}{(R/R_0)^{\frac{1}{k}} + 2.0^k}$$

Fig. 1 shows $T_i$ and $T_e$ as a function of radius for $k = 0.5$ and $k = 1$. The ion temperature is taken in fiducial Run A (see Table 1) at a quasi-steady state before considering radiative cooling.

2.2 Radiative cooling rate $Q_{\text{rad}}$

We now provide the numerical formulae to derive the radiative cooling rate $Q_{\text{rad}}$. Technically there are several treatments on the radiation, e.g. some researchers employ the flux-limited diffusion approximation to solve the radiation energy equation (Ohsuga et al. 2009). Here we take a different and also simplified approach to calculate all the radiative cooling terms. We note that the radiative processes in hot accretion flow include the bremsstrahlung, synchrotron, and the Comptonization (the seed photons are mainly the synchrotron photons, whose energy is much lower compared to that of bremsstrahlung photons.). The radiative cooling rate can then be expressed as,

$$Q_{\text{rad}} = Q_{\text{brem}} + Q_{\text{syn}} + Q_{\text{syn, C}}$$

where $Q_{\text{brem}}$ is the bremsstrahlung emission including both the electron-ion and electron-electron collisions (Eq. 3.4 in NY95).

Optically thin synchrotron radiation is emitted mainly at the local self-absorption peak frequency $\nu_c$, below which the synchrotron emission is self-absorbed. Once $\nu_c$ is determined at each location, we adopt a simplified formula (Eq. 3.18 in NY95) to estimate the synchrotron emission,

$$Q_{\text{syn}} = \frac{2\pi}{3c^2} T_i R^3 R_0^2$$

where $\delta = R/\delta S$, where $S$ is the scale height and $\delta S$ is the surface area, instead of $4/3 \pi R^3$. We then equate the synchrotron emission to the Rayleigh-Jeans blackbody emission from the surface of that volume. The net effect of this modification is to replace the $R$ of eq. 3.14 in NY95 with $3H_i$, i.e. the surface density plays it role in determining $\nu_c$.

One key process is the Compton scattering within the hot accretion flow. In this work we use the Compton enhancement factor $\eta$ (Dermer, Liang & Canfield 1991) which is determined by the electron energy, the seed photons, the geometry of the Compton region and also the effective optical depth $\tau_{\text{eff}}$. In the calculation of $\tau_{\text{eff}}$, we adopt coefficients that correspond to a disc configuration and seed photon energy at 1 eV (Dermer et al. 1991; Eq. 3.19 in NY95). Moreover, since we do not do any integrations in the calculation of either synchrotron or Comptonization, we take $\nu_c$ as the mean seed photon energy. The radiative cooling rate due to the Comptonized synchrotron emission is,

$$Q_{\text{syn, C}} = (\eta(T_i, \nu_c, \tau_{\text{eff}}) - 1) Q_{\text{syn}}$$

Note two should be emphasized. First, all these formulae are based on the local Comptonization approximation which is crude and additional corrections should be made (e.g. Yuan, Xie & Ostriker 2009; Xie et al. 2010; Niedźwiecki et al. 2012). Second, the radiative transfer of the optically thin medium has been approximately taken into account under this approach, although we can not quantitatively estimate how goodness it is. We devote it to future work.
Table 1. Model Parameters and Simulation Properties

| Run | ρ_{max}\,(\rho_0) | β   | k   | t_{awp} (\,\,\text{orb. time at } \rho_{\text{max}}) | M_1/M_{\text{Edd}} | \text{(10)} |
|-----|-----------------|-----|-----|-----------------------------|----------------|----------|
| A   | 1.0             | 10  | 1.0 | 9.0                        | 0.006          |          |
| A1  | 0.5             | 10  | 1.0 | 8.0                        | 0.03           |          |
| A2  | 0.1             | 10  | 1.0 | 5.0                        | 0.006          |          |
| B   | 1.0             | 100 | 1.0 | 5.0                        | 0.06           |          |
| B1  | 1.0             | 1   | 1.0 | 5.0                        | 0.06           |          |
| E   | 0.5             | 10  | 0.5 | 5.0                        | 0.03           |          |
| E1  | 0.5             | 10  |     | 5.0                        | 0.03           |          |

(a) The whole physical simulation time \(t_{\text{awp}}\) in unit of the orbital period at \(R_0 = 100 \, R_\odot\).
(b) The net accretion rate at location \(2 \, R_\odot\), at time 1.4 orbits, when the whole accretion flow is in a quasi-steady state.
(c) In this case, we set \(T_e = T_i\).

2.4 Numerical methods

The public ZEUS-2D code (Stone & Norman 1992) are used in this work to solve the basic hydrodynamic equations. The shear stress and the radiative cooling are considered. The radiative cooling term is treated explicitly. The Courant–Friedrichs–Lewy condition for the radiative cooling is not taken into account. Instead, a sub-cycle technique is adopted for regions with high density (in the cold clumps), i.e. whenever the radiative cooling time-step is smaller than the time-step used for the hydrodynamic equations, we sub-cycle that. The radiative terms are repeatedly at the smaller time-step until one hydrodynamical time-step has elapsed for the purpose of saving computational time.

In our simulations, the inner and outer boundary (in radial direction) of the computational domain are set as, \(R_\text{in} = 1.3 \, R_\odot\) and \(R_\text{out} = 400 \, R_\odot\). We use outflow boundary conditions at these two boundaries so that mass is allowed to flow out of the computational domain but not into it. In the angular direction, the axis symmetric boundary condition is taken.

We followed the treatment of Stone et al. (1999) for the grid setup, i.e. logarithm in both the radial and the angular directions. More specifically, we choose in radial direction \((\Delta \theta)_{\text{rad}}/(\Delta r) = \sqrt{\frac{R}{10}}\) with \(N_r\) grid points in radius. To resolve the cold clumps (and also the possible think disk) near the equator, we adopt non-uniform angular zones with \((\Delta \theta)/(\Delta \theta)_{\text{rad}} = \sqrt{\frac{\pi}{2}}\) for \(0 \leq \theta < \pi/2\), and \((\Delta \theta)/(\Delta \theta)_{\text{rad}} = \sqrt{\frac{\pi}{2}}\) for \(\pi/2 \leq \theta < \pi\). This gives a refinement by a factor of 4 in the angular grid zones between the poles and equator. The standard resolution is \(N_r = 128\) and \(N_\theta = 80\), giving a grid with a total size of 334 \times 160 grid zones. The standard resolution gives the minimum grid of \(\Delta r = 0.015 R_\odot\), \(\Delta \theta = 0.5^\circ\). We have also computed a high resolution model with \(N_r \times N_\theta = 128 \times 160\) (total 334 \times 320). We caution that although the resolution is high enough to resolve the cold clumps formed near the equatorial plane, it is still insufficient to resolve clumps possibly formed at high altitude. Simulation with higher resolution is computationally expensive, which is beyond the scope of current work.

Technically, we first run the simulation (of the equilibrium torus) without radiative cooling. After the accretion flow enters a quasi-steady state, we then start to include the radiative cooling.

3 RESULTS

As summarized in Table 1, we have run seven HD simulations with different model parameters. We take Run A \((\alpha = 0.01, \beta = 10, k = 1.0)\) as the fiducial run. The accretion time is in the units of the orbital period at radius \(R_0 = 100 \, R_\odot\), in all models. We will first present results of this fiducial model to provide a general picture of our simulations, and then will go further to discuss the influences of individual model parameters. Note that the dynamical impact of the viscosity parameter \(\alpha\) on hot accretion flows is well-known and easy to understand (see e.g. Stone, Pringle & Begelman 1999; Igumenshchev & Abramowicz 1999, 2000 for cases without radiation; Yuan & Bu 2010; Das & Sharma 2013 for cases with radiation), and we omit to discuss it here.
Hot accretion flow with radiative cooling

3.1 Fiducial Run A

When the accretion flow reaches a quasi-steady state ($t = 1.4$ orbits), a representative of hard state in BHBs, we start to include the radiative cooling. Very quickly the state transition is triggered. If initially we setup with a very low density torus and the radiative cooling is weak (dynamically un-important), then such state transition will not happen at all. At a later time, $t = 2.0$, the whole accretion flow is now dominated by a cold, geometrically thin accretion disc and surrounded by a hot tenuous media, similar to the disc-corona configuration. This is likely a representative of the soft state. For completeness, we also present the results at $t = 4.0$ and $t = 8.0$, where the mass accretion rate is reduced gradually, because there is no continuous mass supply in our simulations.

To better understand the detailed structure of the accretion flow, Fig. 2 shows the density (upper panels) and ion temperature (bottom panels) for the four representative times as mentioned before and the density panels over-plotted with poloidal velocity field. At time $t = 1.4$, the accretion flow is globally geometrically thick with high temperature. Moreover, it is turbulent inflow near the equatorial plane, and outflow above certain altitude (see Yuan, Wu & Bu 2012 and references therein). At time $t = 2.0$, the cold, dense, geometrically thin accretion flow extends to radii $R > 100 R_s$. After this stage, the dynamical structures of thin disc becomes turbulent as shown in the right two panels of Fig. 2. This is because there is no continuous mass supply in our simulations. Consequently as the system evolves, the gas density reduces, and the radiative cooling becomes less important. In other words, if we run the simulation for sufficiently long time, the accretion flow will eventually return back to a hot accretion flow, i.e. a transition from soft state to hard state. Note that, a soft-to-hard state transition is realized in Run A1, where the initial density and consequently the accretion rate are lower (cf. Fig. 6).

Fig. 3 shows the ratio of disc scale-height to radius $H/R$ at the equatorial plane versus radius for Run A. The black solid, red solid, blue dotted and green dashed curves are respectively, correspond to the time $t = 1.4, 2.0, 4.0, 8.0$.

Figure 2. The logarithmic density (over-plotted with poloidal directions of the velocity as arrows) and temperature, for Run A. Panels from left to right show time at $t = 1.4, 2.0, 4.0, 8.0$, respectively.

Figure 3. The aspect ratio $H/R$ ($H = c_s/Ω_K$) at the equatorial plane versus radius for Run A. The black solid, red solid, blue dotted and green dashed curves are respectively, correspond to the time $t = 1.4, 2.0, 4.0, 8.0$. 

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Figure 4. Radial scaling of time-averaged physical quantities for Run A. From left to right, upper to lower are the density, surface density, velocity and angular momentum, respectively. In each panel, colors denote different evolution time, the black color lines denote $t = 1.4$ and red color line denotes $t = 2.0$. The density, surface density and angular momentum are in units of cgs. Only the velocities are in units of light speed. In the velocity panel, the solid lines denote radial velocity and dotted lines denote sound speed. In the angular momentum panel, dotted line represents the Keplerian angular momentum.

Figure 5. Mass accretion rate $\dot{M}$ (in units of Eddington accretion rate) at $R = 2 R_0$ (black solid line) and $R = 10 R_0$ (red dotted line), as a function of time (in units of the orbital time at $R_0$) for Run A1.

The first and most apparent parameter is the accretion rate, or density of the initial torus, which will control the whole mass supply of the system. To investigate its effect, we run another two simulations with different accretion rate including Run A1 with $\dot{M} = 0.03 M_{\text{Edd}}$ and Run A2 with $\dot{M} = 0.006 M_{\text{Edd}}$. We use a slightly high numerical resolution (total: $334 \times 320$) in Run A1 to discern filament structures. Our simulation results indicate that the critical mass accretion rate is about $0.03 M_{\text{Edd}}$ below this critical accretion rate the hard-to-soft state transition can not happen. For example, the whole accretion flow remains purely hot (not shown here) throughout the evolution time for Run A2 ($\dot{M} = 0.006 M_{\text{Edd}}$). Run A2 is similar to the non-radiative hydrodynamical simulations (e.g. Stone, Pringle & Begelman 1999). In other words, the mass accretion rate is lower than the critical value, thus the radiative cooling of this model are sufficiently low which results in very weak dynamical impact.

Figure 5 shows the evolution of the net mass accretion rate (see Stone, Pringle & Begelman 1999 for definition) at radius of $2 R_0$ (black solid line) and $10 R_0$ (red dotted line) in units of Eddington accretion rate for Run A1. The net accretion rate increases by a factor of $\sim 2$ after considering the radiative cooling. As the gas enters into the central black hole, the mass accretion rate decreases due to no supplementary matter. At time $t \gtrsim 6.5$ orbits, the net accretion rate $\sim 0.02 M_{\text{Edd}}$, which is below the critical rate, so the accretion flow enters into a hard state again. This may
be similar to a soft-to-hard state transition in the decay of BHBs outburst. Moreover, fluctuation in the mass accretion rate at radius of $R = 10 R_s$ is weaker again at time $t \gtrsim 6.5$, mainly because of the disappeared cold clumps.

Fig. 6 shows the density and temperature of the accretion system of Run A1, from left to right panels corresponding to time 1.7, 2.0, 4.0 and 8.0 orbits, respectively. From this plot, obviously there is no conventional cold thin SSD, but instead of filamentary and sometimes clumpy cold and dense structures except for the rightmost panel. These cold and dense structures are concentrated to lower altitude, but clearly they are highly turbulent, without any sign of settling down onto the mid-plane. The truncated radius of this model $R_t \approx 35 R_s$ at time $t = 2.0$ (not show picture here) is the same as Run A which has same accretion rate $\dot{M} \sim 0.03 \dot{M}_{\text{Edd}}$ at time $t = 8.0$. This result is highly close to the two-phase LHAF model proposed by Yuan (2003) (see also Xie & Yuan 2012 and Yang et al. 2015 for more discussions). It is clear that at time $t = 8.0$ the accretion flow becomes hot and dilute again, namely in hard state.

### 3.2.1 Run A1: luminous hot accretion flow?

Fig. 7 shows advection factor of hot accretion flow at time $t = 1.7$ in Run A1. The advection factor is defined as (Narayan & Yi 1994; Yuan & Narayan 2014),

$$f \equiv \frac{Q_{\text{adv}}}{Q_{\text{vis}}} = 1 - \frac{Q_{\text{rad}}}{Q_{\text{vis}}}$$

where $Q_{\text{adv}} = \rho \frac{dT}{dt} + p \nabla \cdot \mathbf{v} = \rho \frac{dS}{dt}$ (S is the entropy) is the so-called advection term. In this plot, the curves represent advection factor $f$ at different angular locations, $\theta = 90^\circ$ (black solid; equatorial plane), $\theta = 70^\circ$ (blue dotted) and $\theta = 50^\circ$ (red dashed), respectively. The advection factor is negative at radii larger than $40 R_s$ although the flow is still in hard state. This is exactly the case of type I luminous hot accretion flow (e.g. Yuan 2001, 2003), which the radiative
cooling is very strong even larger than viscous heating, but the advection term plays a heating role, rather than a cooling role, so the accretion flow remains hot.

3.2.2 radiative cooling in accretion flows

One advantage of Run A1 is that, it shows the purely hot accretion flow (at \( t \geq 6.5 \)), the LHAF (at \( t = 1.7 \)), and the two-phase accretion flow (at \( 1.8 \leq t \leq 6.5 \)). We here investigate in detail the contributions of different radiative cooling processes. We show in Fig. 8 the total radiation cooling rate (multiplied by the volume \( R^2H \), i.e. \( Q_{\text{rad}} \)) and the ratio of the Comptonized synchrotron emission to that of the total radiation (\( (Q_{\text{syn}} + Q_{\text{syn,co}})/Q_{\text{rad}} \)) for both the two-phase accretion flow (right two panels; at \( t = 1.8 \sim 2.0 \)) and the ADAF (right two panels; at \( t = 7.8 \sim 8.0 \)). The first result we derive from this plot is that, the Comptonized synchrotron emission dominates for the regions within \( 40R_\odot \), where most of the radiation comes from. This justifies our argument on the necessity of including synchrotron and the corresponding Comptonization processes. Besides, the radiation of hot accretion flow (both ADAF and the hot-phase medium of the two-phase accretion flow) is spatially extensive, covering a large volume. We note that, if the magnetic field is relatively stronger at high latitude regions where the \( \beta \) value is smaller compared to that of the mid-plane (e.g. in MHD simulations by Hawley & Krolik 2001; Hirose et al. 2009), then the spatial distribution of the radiation will be even more extensive and smooth. Additional emission from the cold clumps make the total emission of the two-phase accretion flow concentrates towards the equatorial plane.

We should point out that, the concentration of emission site of the two-phase accretion flow also depends on the wavebands observed. The hard X-ray emission should still be spatially extensive, as they are produced by the Compton scattering of the hot electrons. The soft X-ray and UV emission, on the other hand, will likely concentrates to the mid-plane, if the clumps are indeed most abundant there. Currently it remains unclear whether there will be more numerous smaller clumps formed, and whether these clumps can be supported to stay at higher altitude by the magnetic field. Further MHD simulations with higher spatial resolution will be required to clarify these questions.

3.3 Effects of plasma parameter \( \beta \)

In order to investigate the effect of magnetic field strength, we also run two additional models with \( \beta = 100 \) (Run B) and \( \beta = 1 \) (Run B1). The other model parameters are the same as those of the fiducial run. Fig. 9 shows the profiles of density, ion temperature and radial velocity at the mid-plane, before (averaged over the period \( t = 1.3 \sim 1.4 \) top panels) and after (averaged over the period \( t = 1.8 \sim 2.0 \); bottom panels) including the radiative cooling.

We first analyze the results without radiative cooling. If magnetic field is weak (Run A and Run B), the accretion flow will be gas pressure dominated. Consequently, there will be no significant differences between these simulations. The ion temperature of Run B is slightly higher, mainly because the gas pressure (roughly \( \rho T_e \)) increases as \( \beta \) increases, as the total pressure remains less affected. If magnetic field becomes strong enough, i.e. \( \beta \sim 1 \), its dynamical impact will be evident, even without considering the radiation processes. Indeed, when comparing Run B1 with Run A, we find that the density becomes lower in low \( \beta \) case. This is because, we additionally put the magnetic pressure during the initial setup of the equilibrium (gravity against gas pressure gradient force) torus. Consequently, more fraction of the accreting material will be driven out as outflow due to the enhanced pressure gradient force at the early stage of evolution. The density will be lower and the accretion rate will be reduced. Besides, the ion temperature of Run B1, whose accretion rate is lower, is lower compared to that of Run A. This is because of weaker compression work (\( pd(1/\rho) \)), which is proportional to density gradient; cf. the density profile in Fig. 9) done onto the gas in Run B1.

We then compare results with and without radiative cooling. The synchrotron emission, which provides the seed photons for the Compton scattering process, is more sensitive to the electron temperature than the magnetic field strength, i.e. we roughly have \( Q_{\text{syn}} \sim T_e^2B^2 \sim T_e^2P_{\text{mag}}^2 \) (This can be derived from Eq. 8, and note that \( v_c \) is roughly proportional to \( BT_e^2 \); NY95; See also Mehadevan 1997 for the dependence on \( T_e \)). In this sense, The impact of including radiative cooling in Run B, whose electron temperature is the highest before considering radiative cooling (cf. top middle panels of Fig. 9), is significant, while the impact of including radiative cooling in Run B1 is much weaker.
Figure 9. Comparison of the properties of simulations with different $\beta$ values, i.e. the solid, dotted and dashed curves correspond to models with $\beta = 1$, $\beta = 10$ and $\beta = 100$, corresponding to Run B1, Run A and Run B, respectively. The upper panels show the averaged quantities (density, temperature and radial velocity) at the mid-plane over the period $1.3 - 1.4$ (without radiative cooling), while bottom panels show averaged quantities over the period $1.8 - 2.0$ (with radiative cooling).

Figure 10. Comparison of simulations with different $T_e/T_i$ values over period of $t=1.8 - 2.0$ orbits. As labeled in the figure, Run A1, Run E and Run E1 have $k = 1$, $k = 0.5$ and $T_e = T_i$ (sole temperature simulation), while they share the same other initial conditions. In the leftmost panel, the solid and dashed curves are the viscous heating rate and the total radiative cooling rate (both are multiplied by $R^2 H$). In the right three panels, show the density, temperature (solid line denotes the ions temperature, the dashed lines denotes the electrons temperature) and radial velocity in units of light speed (the dashed line denotes free fall speed.)

especially on the temperature of the gas. The difference of the total radiative cooling rates of the three runs are highly reduced, i.e. they differ by factors of $\lesssim 2$ for our chosen parameters. We note that the similarities of the three runs after including radiative cooling may be a coincidence, and should not be taken seriously. Radiative MHD simulations are required for a better understanding on the impacts of the magnetic fields, both dynamically (through magneto-rotational instability) and radiatively (through synchrotron emission).

Moreover, we also point out that our approach can only mimic the randomly-oriented turbulent magnetic fields, but not the global large scale one. The non-radiative simulations are different from those non-radiative MHD simulations, e.g. Stone & Pringle (2001). Their results show that the mass accretion rate depends on the initial magnetic field...
strength, i.e. stronger magnetic field will lead to higher mass accretion rate. The reason is that, the Maxwell stress, which determines the capability of transporting angular momentum, will be stronger for cases with stronger vertical fields. In other words, stronger magnetic fields will generally have larger ‘effective’ $\alpha$, while it is fixed to a constant in our HD simulations. Furthermore, in the MHD case, when radiative cooling becomes dominant, ADAF-like gaseous disk is shrunked in the vertical direction with frozen-in magnetic field, and magnetic pressure becomes relatively dominant in the cool disk. Hirose et al. (2009) showed that the plasma parameter $\beta$ in coronal region became lower value than the disk region when radiative cooling becomes important. Therefore, the assumption of constant $\beta$ will under-estimate synchrotron radiation.

### 3.4 The dependence on temperature ratio $T_e/T_i$

Obviously, only when radiative cooling becomes important in determining the dynamical structure of the accretion flow, the simulation results will be sensitive to the value of $T_e/T_i$ (Mościbrodzka et al. 2009; Dibi et al. 2012).

From theoretical point of view, $T_e/T_i$ depends mainly on two quantities. The first is the accretion rate. At sufficiently low accretion rate (e.g. the quiescent state and also the low luminosity hard state), the radiative cooling of electrons is insignificent, thus the $T_e/T_i$ will be moderate (it depends on the energy deposition and also the adiabatic index differences between electrons and ions; cf. Yuan & Narayan 2014). As the accretion rate increases (i.e. normal hard state), the electrons will suffer additional radiative cooling and $T_e$ will be reduced. $T_i$ on the other hand, remains nearly unaffected, mainly because the Coulomb collision is small compared to the viscous heating to ions at such moderate accretion rates. In other words, the value of $T_e/T_i$ will become smaller in this regime. When the accretion rate is very high (e.g. the bright hard state and the intermediate state, i.e. the LHAF regime), the Coulomb coupling between electrons and ions, which scales as $\rho^2 (T_e - T_i)$, is so strong that the temperature differences between electrons and ions are reduced, i.e. $T_e/T_i$ will become moderately large again (see e.g. Yuan 2001; Xie et al. 2010). Note that for cases of SSD and possibly the cold clumps of the two-phase accretion flow, we will have $T_e = T_i$. Observationally we do observe an anti-correlation between electron temperature and the bolometric luminosity for the outburst of black hole X-ray binaries (e.g. Joinet et al. 2008; Miyakawa et al. 2008; Motta et al. 2009; Natalucci et al. 2014). The second quantity is the fraction of viscous heating that goes to electrons directly (see e.g. Xie & Yuan 2012 for a brief summary on constrains of this value). For given density and accretion rate, if more fraction of viscous heating goes to the electrons, obviously it will have a higher value of $T_e/T_i$, and the corresponding bolometric luminosity will also be enhanced.

We run two new simulations to investigate the impact of temperature ratio. With other parameters the same to those of Run A1, we set in Run E $k = 0.5$ and in Run E1 $T_e = T_i$. The densities are high in these runs, and the radiative cooling is dynamically important, representing the bright hard (and also intermediate) state. As radiative cooling is triggered, differences are evident, as shown in Fig. 10 when a quasi-steady state is reached.

Reducing the differences between $T_e$ and $T_i$ (as in Run E1 and Run E, compared to Run A1) is equivalent to enhance the coupling between electrons and ions. The $T_e$ will then be increased, while the $T_i$ will be reduced. The scale height of the hot accretion flow ($\propto (T_e + T_i)^{1/2}$) will also be reduced, which leads to an enhancement in gas density. The radiative cooling rate (synchrotron and its Comptonization, bremsstrahlung) will also be increased. All these results, as shown in Fig. 10, are easy to understand.

In the case of Run E series they start to form clumpy structures at lower accretion rate, i.e. they have more lower critical accretion rate of the purely hot accretion flow compared to those with $k = 1$ models. We also run some other simulations (not shown here) and the results confirm this point, such as the critical accretion rate is reduced to ~ 0.006 $\dot{M}_{\text{Edd}}$ in Run E1. The ratio between electron and ion temperature is the key point for the radiative cooling process, but it is very complex and depends on many factors and microphysical processes. Two-fluid (electrons and ions) radiative (MHD) simulations are needed to solve this problem. Only under this approach, we can investigate impact of the fraction of viscous heating to electrons, and come out a more reasonable $T_e/T_i$ value (i.e. varies at different location and/or time).

### 4 SUMMARY AND DISCUSSIONS

In this paper, we present a number of two-dimensional hydrodynamical simulations of accretion flow on to a black hole, focusing on the impact of radiation cooling and state transitions. Compared to previous works (e.g. Das & Sharma 2013), our technical improvements includes: (1), the hot accretion flow in our simulations is two-temperature (NY95); (2), we follow the treatment of NY95 to consider the synchrotron emission and the corresponding Compton up-scattering process (see also Fragile & Meier 2009), which are crucial to generate the radiation observed in black hole X-ray binaries.

We summarize our major results as follows. Firstly, as the accretion rate increases, the truncated radius moves inward. This result is consistent with theoretical expectation (Esin et al. 1997) and the numerical simulations by Das & Sharma (2013). Secondly, when the accretion rate reaches the critical value $M = 0.03 \dot{M}_{\text{Edd}} \approx 3 \times \dot{M}_{\text{Edd}}$ (Xie & Yuan 2012), we observe the formation of clumpy/filamentary structures within the hot medium, i.e. the whole accretion flow is likely to be in a two-phase accretion mode (see, e.g. Yuan 2003; Xie & Yuan 2012). As the accretion rate increases, these clumps grow and/or merge. Eventually the turbulent motions of the hot gas can not support them any longer, and they will settle down onto the mid-plane and form a thin disc. In other words, the clumpy accretion mode do exist before the eventually formation of the disc-corona configuration. It is speculated that this accretion mode may be responsible for the intermediate state of X-ray binaries (e.g. Yang et al. 2015).

We still lack direct evidences for the existence of cold clumps (or alternatively filaments, clouds), especially in BHBs. In AGNs, some clues can be derived from variability...
studies. For example, in a broad absorption line (BAL; outflowing gas with velocity ~ 0.1 c) quasar monitoring campaign, Capellupo et al. (2012) found that the variability favours the idea that those absorptions varies (on timescale of years) because of the movements of individual outflowing clouds (or at least substructures in the moving flow). Another example comes from NGC 5548 (Kaastra et al. 2014), where clumpy streams of ionized hot gas blocking the emission from the nucleus are observed.

Below we will highlight several notable shortages of our work, and discuss their consequences.

4.1 Oversimplification in the radiative cooling

In this work, we highly simplify the calculation of radiative cooling. The first is that, we use the critical frequency $\nu_c$ of the self-absorbed synchrotron emission to estimate the synchrotron cooling rate. Besides, we also use this frequency as the representative seed photon frequency for the Compton scattering process. Such approximation, which avoids any integration (in frequency $\nu$), highly accelerates the computational speed. However, its accuracy is difficult to constrain.

More importantly, our Comptonization is based on the local Compton scattering approximation. In the optically thin medium as the case of hot accretion flow, seed photons can propagate a long distance from one location to another, and scatter with the hot electrons there. Such “global” Compton scattering process is not taken into account in our simulations. The global Comptonization process is repeatedly examined in one-dimensional calculations, and it will make the electron temperature profile more flattened and the absolute $T_e$ value reduced (by a factor of ~ 1/3, see e.g. Yuan, Xie & Ostriker 2009; Xie et al. 2010; Niedźwiecki et al. 2012). If we do not want to employ the computationally-expensive Monte Carlo technique, then one easy way to consider this effect is to create a correction factor profile from these numerical calculations and then apply it to the simulations here. However, such treatment is also very crude, and we omit it in this work.

4.2 Magnetic fields (large-scale ordered) in the hot accretion flow

In this work we used plasma $\beta$ parameter to consider the effects of magnetic field (and correspondingly the synchrotron emission). Such approach has several limitations, which we will discuss below.

Firstly, this approach can not properly take into account the large-scale ordered magnetic fields. Plasma $\beta$ can only provide a feasible and reasonable approximation for the consideration of the tangled, turbulent magnetic fields, but even in this case, it is not a constant within the accretion flow (Hawley & Krolik 2001). Besides, the dynamical structure of hot accretion flow will be highly revised if such large-scale magnetic field exist (Machida et al. 2006; Fragile & Meier 2009). The “effective” $\alpha$ will be highly enhanced if the accretion flow has large net magnetic flux (cf. Penna et al. 2013; Bai & Stone 2013). Besides, the large scale magnetic field may also be crucial for the formation and acceleration of the Poynting-dominated jet powered by the black hole spin (Blandford & Znajek 1977; Tchekhovskoy, Narayan & McKinney 2011), the disk-jet powered by the disk rotation (Yuan & Narayan 2014; Yuan et al. 2015), or jet powered by the radiation from supercritical accretion flow (Sadowski & Narayan 2015).

Secondly, one evolutionary stage of our simulations (e.g. in Run A1), before the eventual collapse into a cold thin disc, is the formation of the cold and dense clumpy/filamentary components. If these cold clumps are still coupled to the magnetic fields (valid likely only for small clumps), then the magnetic fields may service as natural barrier for the merger of those small clumps, i.e. the clumps will be confined by the large-scale magnetic fields. Besides, the large scale magnetic fields will provide additional vertical support for those cooling gas components, thus maintain them at higher altitude (Machida et al. 2006).

Finally, one interesting property related to the accumulation of the large scale global magnetic fields is that, they may help on the understanding of the the hysteretic cycle of black hole state transitions (Begelman & Armitage 2014 and references therein for alternative scenarios), i.e. the critical luminosity of the hard-to-soft transition is generally brighter than that of the soft-to-hard state. In the scenario proposed by Begelman & Armitage (2014), they expect that, during the hard state, vertical magnetic flux, which is generated near the truncated radius $R_t$, is advected inward and accumulates stochastically within the inner hot flow. Consequently the $\alpha$ will be large, leading to a higher critical luminosity for the hard-to-soft state transition (see also Yu et al. 2015 for the stability analysis of low-$\beta$ hot accretion flows). On the other hand, the magnetic flux is gradually diffused out in the cold SSD during the soft state, thus the critical luminosity of the soft-to-hard state transition will be much lower. The numerical realization of this scenario is beyond the scope of current work.

4.3 Turbulent dissipation and corona heating

Another issue is related to the turbulent/viscous dissipation within the accretion flow. In this work (and also most of the hydrodynamical simulations of accretion flows), we adopt the $\alpha$-viscosity prescription (Shakura & Sunyaev 1973), which assumes that the dissipation heating rate is proportional to the density of the gas. One consequence of this assumption is that, it assumes most of the “viscous” heating happens close to the mid-plane of the accretion flow.

Contradict to this assumption, recent MHD simulations, in which the magnetic reconnection is considered as the “physical” heating mechanism of the accretion disc, indicate that the vertical distribution of energy dissipation actually do not follow the density profile, but can extend to much higher height (e.g. Jiang, Stone & Davis 2014 and references therein). If this is indeed the case, then it will provide a natural heating mechanism for the corona in those disc-corona models (see e.g. Liu et al. 2002; Cao 2009), where they generally assume a significant fraction (of order ~ 50%) of the “viscous” heating energy goes into the corona region. Besides, if such differences are taken into account, the mid-plane regions will be cooler, while the high altitude regions should be hotter.
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