Luminous hot accretion disks

Feng Yuan\textsuperscript{1,2}\textasteriskcentered
\textsuperscript{1} Department of Astronomy, Nanjing University, Nanjing 210093, China; Email: fyuan\text@nju\text.edu.cn
\textsuperscript{2} CAS-PKU Joint Beijing Astrophysical Centre, Beijing 100871, China

ABSTRACT

We find a new two-temperature hot branch of equilibrium solutions for stationary accretion disks around black holes. In units of Eddington accretion rate defined as $10L_{\text{Edd}}/c^2$, the accretion rates to which these solutions correspond are within the range $\dot{m}_1 \lesssim \dot{m} \lesssim 1$, here $\dot{m}_1$ is the critical rate of advection-dominated accretion flow (ADAF). In these solutions, the energy loss rate of the ions by Coulomb energy transfer between the ions and electrons is larger than the viscous heating rate and it is the advective heating together with the viscous dissipation that balances the Coulomb cooling of ions. When $\dot{m}_1 \lesssim \dot{m} \lesssim \dot{m}_2$, where $\dot{m}_2 \sim 5\dot{m}_1 < 1$, the accretion flow remains hot throughout the disk. When $\dot{m}_2 \lesssim \dot{m} \lesssim 1$, Coulomb interaction will cool the inner region of the disk within a certain radius ($r_t \sim$ several – tens of Schwarzschild radii or larger depending on the accretion rate and the outer boundary condition) and the disk will collapse onto the equatorial plane and form an optically thick cold annulus. Compared to ADAF, these hot solutions are much more luminous because of the high accretion rate and efficiency, therefore, we name them luminous hot accretion disks.

Key words: accretion, accretion disks – black hole physics – galaxies: active – galaxies: nuclei hydrodynamics – radiation mechanisms: thermal

1 INTRODUCTION

There are currently four black hole accretion disk models. They are the standard thin disk (Shakura & Sunyaev 1973), slim disk (Abramowicz et al. 1988), cooling-dominated hot disk (Shapiro, Lightman, & Eardley 1976, hereafter SLE), and advection-dominated accretion flow (ADAF) (see reviews by Narayan, Mahadevan, & Quataert 1998 and Kato, Fukue & Mineshige 1998). Among them, the former two models are cold, while SLE is thermally unstable. Therefore, ADAF is the only viable hot accretion disk model up to date. The success of this model is in virtue of the two-temperature plasma concept first put forward by SLE and the realization to the importance of the advection term in the energy equation of ions (Ichimura 1977; Rees et al. 1982; Narayan & Yi 1994, 1995; Abramowicz et al. 1995). In a typical two-temperature ADAF, the accretion rate is very low and the Coulomb energy transfer from ions to electrons is very inefficient therefore almost all of the viscously dissipated energy is stored in the ions and advected into the centre black hole rather than transferred to the electrons and radiated away. In other words, in the energy equation of ions,

$$\rho v T_i \frac{ds}{dr} = \rho v \frac{de_i}{dr} - q^c \equiv q_{\text{adv}} = q^+ - q_{\text{ic}},$$

where $s$ and $e_i$ are the entropy and internal energy of the ions per unit mass of plasma, $q^c$ and $q_{\text{ic}}$ are the compressive heating and Coulomb energy transfer from ions to electrons, and $q_{\text{adv}}$ denotes the energy advection, we have for a typical ADAF,

$$q_{\text{adv}} \approx q^+ \gg q_{\text{ic}}.$$

\textsuperscript{*} present address: Max-Planck-Institut für Radioastronomie, Auf dem Hügel 69, D-53121 Bonn, Germany. Email: fyuan@mpef-bonn.mpg.de

© 2000 RAS
It is well known that the optically thin two-temperature advection-dominated accretion flow (ADAF) solution exists only for the mass accretion rate less than a critical value $\dot{m}_1$ (Ichimura 1977; Rees et al. 1982; Narayan & Yi 1995; Abramowicz et al. 1995; Esin et al. 1996, 1997). This is because $q^+ \propto \dot{m}$ while $q_e \propto \dot{m}^2$, i.e., $q_e$ increases faster than $q^+$ with the increasing accretion rate. When the accretion rate reaches a certain critical value, the Coulomb coupling between the ions and electrons becomes so efficient that a large fraction of the viscously dissipated energy is transferred to the electrons and radiated away therefore the accretion flow ceases to be an ADAF. The critical accretion rate of ADAF $\dot{m}_1$ is determined by the balance between the viscous heating and the Coulomb energy transfer from the ions to the electrons

$$q^+ \approx q_e.$$  \hfill (3)

(Narayan, Mahadevan, & Quataert 1998). Due to its low accretion rate and efficiency, an ADAF is unable to emit significant radiation, i.e., it is a dim hot accretion disk.

An interesting question is, what will happen when the mass accretion rate increases above $\dot{m}_1$. It is in general assumed that in this case, hot solutions don’t exist and the standard thin disk is the only viable solution. However, this is not true. From eq. (1), we know that in addition to the viscous dissipation $q^+$, the compression work $q^\circ$ is also a term which can heat the ions. The plasma can remain hot if the ion cooling rate, $q_e$, is lower than the sum of $q^+$ and $q^\circ$. Since $q^\circ \propto \dot{m}$, we expect that there must exist another critical accretion rate below which the accretion flow can be hot. This critical accretion rate, $\dot{m}_2$, is determined by

$$q_e \approx q^\circ + q^+.$$  \hfill (4)

The critical rate $\dot{m}_2$ could be obviously larger than $\dot{m}_1$ if $q^\circ$ is in the same order to, or larger than, $q^+$, as we will show by simple analytical estimate and exact numerical calculation in the following. And more importantly, there must exist another type of hot accretion solutions in addition to ADAF, which corresponds to accretion rate between $\dot{m}_1$ and $\dot{m}_2$. In these solutions, the ions cooling rate is lower than the sum of $q^+$ and $q^\circ$ but higher than the viscous dissipation rate $q^+$. Obviously, compared to ADAF, this new hot solution will be much more luminous.

## 2 LUMINOUS HOT ACCRETION DISK

Before exploring deeply this new solution, we first give some simple analytical estimate to the range of accretion rate within which this solution exists. The viscous dissipation, compression work, and the Coulomb energy transfer in the ions energy equation (1) have the following forms,

$$q^+ = \rho \nu r^2 \left( d\Omega^i \over dr \right)^2 = \alpha \rho c_s H r^2 \left( d\Omega^i \over dr \right)^2,$$

$$q^\circ = -\rho v p d \left( 1 \over \rho \right) / dr,$$  \hfill (5)

and

$$q_e = \frac{3}{2} \frac{m_e}{m_i} n_e n_i \sigma_T c \ln \Lambda \left( k T_i - k T_e \right) \left( \frac{2}{3} \right)^{1/2} \left( \theta_e + \theta_i \right)^{1/2} \frac{\left( \theta_e + \theta_i \right)^{3/2}}{\left( \theta_e + \theta_i \right)^{3/2}}.$$  \hfill (6)

(see Dermer, Liang & Canfield 1991 for $q_e$). Here the subscribes “i” and “e” denote quantities for ions and electrons, respectively. The Coulomb logarithm $\ln \Lambda \approx 15$ for stellar-mass black holes sources and $\approx 20$ for AGNs, $\theta_i = k T_i / m_i c^2$ and $\theta_e = k T_e / m_e c^2$. Following Narayan, Mahadevan, & Quataert (1998), using the self-similar scaling law obtained by Narayan & Yi (1995) (see also Mahadevan 1997), we obtain,

$$q^+ \approx 3 \times 10^{20} m^{-2} \dot{m} (r/r_g)^{-4},$$  \hfill (7)

$$q^\circ \approx 5 \times 10^8 m^{-2} \dot{m} (r/r_g)^{-3} T_i,$$  \hfill (8)

and

$$q_e \approx 1.8 \times 10^8 \alpha^{-2} m^{-2} \dot{m}^2 (r/r_g)^{-3} (T_i - T_e) \theta_e^{-3/2}.$$  \hfill (9)

All the three quantities above are in cgs units. Here the accretion rate is in units of Eddington accretion rate defined as $10 L_{\text{Edd}} / c^2$, the black hole mass is in units of solar mass, $M = m M_\odot$, $r_g$ denotes the black hole radius, $r_g = 2 G M / c^2$. From eqs. (3), (8) and (10), setting the ions temperature as virial, $T_i \simeq 2 \times 10^{12} \beta (r/r_g)^{-1}$, where $1 - \beta$ is the ratio of the magnetic pressure to the sum of the magnetic pressure and the gas pressure, and $\theta_e \simeq 1/3$, we obtain the critical accretion rate for ADAF,

$$\dot{m}_1 \approx 2 \times 10^{12} \alpha^2 T_i^{-1} \theta_e^{-3/2} \approx 0.4 \alpha^{-2}.$$  \hfill (10)
We set $\beta = 0.5$ throughout the present paper, i.e., we assume exact equilibrium between the magnetic and gas pressure. From eqs. (4), (8), (9), and (10), we obtain $\dot{m}_2$,

$$\dot{m}_2 \approx \frac{3 \times 10^{12} (r/r_g)^{-1}_s + 57}{8 T_i} \alpha^2. \quad (12)$$

If we set $T_i \approx 10^{12} (r/r_g)^{-1}_s$, then $\dot{m}_2 \approx \alpha^2$. In fact, the temperature of ions will be moderately lower than the virial value as $\dot{m} \approx \dot{m}_2$ since according to our definition of $\dot{m}_2$, most of the compression work and viscous dissipation are used to compensate the Coulomb energy transfer loss rather than increase the internal energy of ions. Thus, the value of $\dot{m}_2$ will be much larger than $\alpha^2$. The exact value of $\dot{m}_2$ can only be obtained by self-consistently solving the radiation hydrodynamic accretion equations by numerical calculation.

New hot accretion solutions must exist when the accretion rate is higher than $\dot{m}_1$ but less than $\dot{m}_2$. In these solutions, the ion cooling rate $q^i_e$ is so efficient due to the high accretion rate that the viscous dissipation alone is not large enough to compensate it, i.e., the right-hand side of eq. (1) is negative. It is the compression work of the accretion flow, another heating source of plasma, together with $q^g$ plays this role. In other words, the energy advection now is a heating rather than a cooling term in the ions energy balance equation. The entropy of the flows therefore decrease with the decreasing radii, similar to the Bondi accretion and the cooling flows in galaxy clusters. It is also very similar to the electrons in a typical ADAF with $\dot{m} \ll \dot{m}_1$. As pointed out by Nakamura et al. (1997), in that case, the energy advection by electrons plays a heating rather than cooling role that compensates the radiative cooling of electrons.

Another effect arisen when we numerically solve the accretion equations is that $\dot{m}_2$ is actually a function of radius. We will sign $\dot{m}_2$ in the following as the critical rate independent of radius below which a hot solution exists throughout the disk, from the outer boundary to the horizon. When $\dot{m}$ is greater than $\dot{m}_2$ but less than $\sim 1$, we find by numerical calculation presented in the next section that below a certain radius the Coulomb interaction between the ions and electrons will efficiently cool the inner part of the accretion flow. As a result, the hot accretion flows will collapse onto the equatorial plane, forming an optically thick cold annulus. The exact value of the transition radius depends on the parameters such as $\dot{m}$ and the outer boundary condition of the accretion flow, as shown by our numerical calculation below. Typically it equals several or several tens of $r_g$.

This result, i.e., when the accretion rate reaches a high value a transition from an outer hot disk to an inner cold disk will occur, is first anticipated by Pringle, Rees, & Pachocyk (1973) qualitatively and further developed in more detail by Begelman, Sikora, & Rees (1987) in the context of quasi-spherical accretion onto a black hole where the variation of the transition radius with the accretion rates is obtained although the detailed dynamically self-consistent solutions are lacking. Narayan & Popham (1993) also found this result in one of their examples of the numerical solutions of the boundary layer of the standard thin disk.

Some features of this new hot accretion solution can be expected immediately. The ions temperature should be high, close to or moderately lower than the virial value depending on the mass accretion rate $\dot{m}$. The efficiency should be high because not only the viscous dissipation but also the compression work will be transferred to the electrons and radiated away. Compared to ADAF, the emergent luminosity of this hot accretion disk will be much higher because of the high efficiency and accretion rate.

3 NUMERICAL CALCULATION RESULTS

In this section, we present our exact numerical calculation results by self-consistently solving the radiation hydrodynamic accretion equations. Paczyński & Witta (Paczyński & Witta 1980) potential is adopted to mimic the geometry of a Schwarzschild black hole. Steady axisymmetric and two-temperature assumptions to the accretion flow are adopted. A randomly oriented magnetic field is assumed to exist in the accretion flow and the magnetic pressure is in exact equilibrium to the gas pressure ($\beta = 0.5$). In this case, the radiation is very efficient thus $\dot{\theta}_c \ll 1$ for $\dot{m} > \dot{m}_1$. Since thermal pair production is very sensitive to the electron temperature,

$$\dot{\theta}_c \propto \exp(-2/\dot{m}), \quad (13)$$

we neglect pair effect in the present paper. The radiation pressure $p_{rad}$ is also neglected because we find it is less than 10% of the gas pressure even though when $\dot{m} \sim 1$. However, $p_{rad}$ becomes comparable to the gas pressure in the transition region from the hot disk to the thin annulus when $\dot{m} \gtrsim \dot{m}_2$.

The estimates to $\dot{m}_1$ and $\dot{m}_2$ in section 2 are for a diffusion-type viscous description where the kinetic viscous coefficient $\nu$ have the form $\nu = \alpha \sigma c_s H$. Another type of viscous description widely used in the literature is assuming that the viscous stress tensor is proportional to the total pressure,

$$\tau_{\nu \varphi} = \alpha p = \alpha (p_{gas} + p_{mag}). \quad (14)$$

© 2000 RAS, MNRAS 000, 1-3
Figure 1. The variations with the radii of Mach number, electrons and ions temperatures $T_e$ and $T_i$, the specific angular momentum $l$, the surface density $\Sigma(\equiv \rho H)$, and the ratio of disk height to radius $H/r$ for some hot accretion solutions. Solid (ADAF): $\dot{m} = 0.05 < \dot{m}_1$, $T_i = 6 \times 10^8 K$, $T_e = 2 \times 10^8 K$, $v/c_s = 0.4$; dotted (critical ADAF): $\dot{m} = 0.1 \approx \dot{m}_1$, $T_i = 8 \times 10^8 K$, $T_e = 2 \times 10^8 K$, $v/c_s = 0.4$; dashed (new hot solution): $\dot{m} = 0.3 < \dot{m}_2$, $T_i = 1.4 \times 10^9 K$, $T_e = 2 \times 10^9 K$, $v/c_s = 0.6$; long-dashed (new hot solution): $\dot{m} = 0.5 > \dot{m}_2$, $T_i = 1.8 \times 10^9 K$, $T_e = 10^9 K$, $v/c_s = 0.6$. All are for $\alpha = 0.3$, $M = 10M_\odot$ and $r_{out} = 100r_g$. The units of $\Sigma$ and $T$ are g cm$^{-2}$ and K while $l$ is in c=G=M=1 units.

We will adopt this type of viscous description because in this case the no-torque boundary condition required in the horizon of the black hole to solve the set of accretion differential equations can be automatically satisfied hence significantly simplify our numerical calculation (Abramowicz et al. 1988; Narayan, Kato, & Honma 1997).

The equations of mass conservation and the hydrodynamic balance in the vertical direction of the disk are,

$$-4\pi rH \rho \dot{w} = \dot{M}, \quad \text{and} \quad H = c_s/\Omega_k \equiv \sqrt{p/\rho}/\Omega_k. \quad (15)$$

The radial and axial momentum equations are,

$$\frac{dv}{dr} = -\Omega_k^2 r + \Omega_r^2 r - \frac{1}{\rho} \frac{dp}{d\rho}, \quad (16)$$

$$v(\Omega r^2 - j) = \alpha r \frac{p}{\rho}. \quad (17)$$

The energy equation for electrons is,

$$\rho v \left( \frac{d\varepsilon_e}{dr} + p_e \frac{d}{d\rho} \left( \frac{1}{\rho} \right) \right) = q_e - q^-, \quad (18)$$

where $\varepsilon_e$ denotes the internal energy of the electron per unit mass of the gas.

The radiation mechanisms we consider include bremsstrahlung, synchrotron radiation and Comptonization of soft photons. Assuming the disk is isothermal in the vertical direction, the spectrum of unscattered photons at a given radius is calculated by solving the radiative transfer equation in the vertical direction of the disk based upon the two-stream approximation (Rybicki & Lightman 1979). The result is (Manmoto, Mineshige & Kusunose 1997):

$$F_\nu = \frac{2\pi}{\sqrt{3}} B_\nu [1 - \exp(-2\sqrt{r}_\nu)], \quad (19)$$
where \( \tau^*_\nu \equiv (\pi^{1/2}/2)\kappa_\nu H \) is the optical depth for absorption of the accretion flow in the vertical direction with \( \kappa_\nu = \chi_\nu/(4\pi B_\nu) \) being the absorption coefficient, where \( \chi_\nu = \chi_{\nu,\text{brems}} + \chi_{\nu,\text{synch}} \) is the emissivity, and \( \chi_{\nu,\text{brems}} \) and \( \chi_{\nu,\text{synch}} \) are the bremsstrahlung and synchrotron emissivities, respectively. Then the local radiative cooling rate \( q^- \) reads as follows:

\[
q^- = \frac{1}{2H} \int d\nu \eta(\nu)2F_\nu,
\]

where \( \eta \) is the energy enhancement factor first introduced by Dermer, Liang & Canfield (1991) and modified by Esin et al. (1996).

We numerically solve the above coupled radiation hydrodynamic equations (1), (5) – (7), and (15)–(20). The solutions must satisfy the no-torque condition at the horizon, a sonic point condition at a sonic point, and the outer boundary condition (OBCs) at a certain outer boundary \( r_{\text{out}} \). We here choose OBCs as the temperature of ions and electrons, \( T_{i,e} \), and the ratio of the radial velocity of the flows to the local sound speed, \( v/c_s \) at \( r_{\text{out}} \). The numerical approach is presented in Yuan et al. (2000) (see also Nakamura et al. 1997; Matsumoto et al. 1997). We would like to emphasize here that the OBCs are again found to play an important role in determining the dynamics of the solution, as pointed out by Yuan (1999) in the general context of optically thin accretion flows. Both the exact value of \( \dot{m}_{1,2} \) and the transition radius between the outer hot disk and the inner cold thin annulus when \( \dot{m} \gtrsim \dot{m}_{1,2} \) depend on OBCs. The effects of OBCs are expected to play a more important role in calculating the emergent spectrum (Yuan et al. 2000). Since we concern here only the most general dynamics, we don’t investigate in detail the effect of OBCs in the present study.

Figure 1 shows the dynamic features of four hot solutions with different accretion rates. The parameters are \( \alpha = 0.3, m = 10 \). The outer boundary is set at \( 100r_g \), a relatively small radius, since we first concentrate on the outer cold disk plus inner hot disk configuration. This two-components configuration is believed to be the most promising geometry for the hard state of the Galactic black hole candidates and Seyfert galaxies (see Zdziarski 2000 for a review). The solid, dotted, dashed, and long-dashed lines are for \( \dot{m} = 0.05, 0.1, 0.3 \) and 0.5, corresponding to a typical ADAF with \( \dot{m} < \dot{m}_1 \), a critical ADAF with \( \dot{m} \approx \dot{m}_1 \), a new hot accretion solution with \( \dot{m} < \dot{m}_2 \), and a new hot solution with \( \dot{m} > \dot{m}_2 \) (a transition to an inner cold thin annulus occurs in this case), respectively. We find from the figure that the ions temperature decreases with the increasing \( \dot{m} \). This is because the two heating terms \( q^+ \) and \( q^- \) in the ions energy equation are both proportional to \( \dot{m} \) while the cooling term \( q_{\text{cool}} \propto \dot{m}^2 \), therefore, a larger and larger fraction of energy is transferred to the electrons and radiated away rather than stored in the plasma to increase the internal energy of ions. This is also the reason why the ions temperature \( T_i \) increases slower inward with increasing \( \dot{m} \). For the long-dashed line which denotes a hot solution with a transition to an inner cold annulus, the temperature of the ions in the hot disk part almost remains constant. This shows that the sum of the compression work and the viscous dissipated energy are nearly equal to the Coulomb energy transfer from ions to electrons.

The two plots in Figure 2 show the corresponding energy balance relationship for ions (left) and electrons (right) for the four solutions presented in Figure 1. The left plot shows the variation of the “advection factor” \( f \equiv q_{\text{adv}}/q^+ \) with radii. We
note that $f$ is in general set as an “average” value over radius $r$ in a lot of applications of ADAF (e.g. Narayan, McClintock, & Yi 1996; Esin, McClintock, & Narayan 1997; Narayan et al. 1998). By self-consistently solving the radiation hydrodynamic coupled equations, we find $f$ is typically a sensitive function of radius (see also Figure 5 below) if $\dot{m}$ is not very small. From Figure 2 we see that $f \approx 1$ throughout the disk for a typical ADAF especially with $\dot{m} \ll \dot{m}_1$, therefore, the “average” advection factor is a good assumption in this case (Narayan, McClintock, & Yi 1996; Narayan et al. 1998). However, for the critical ADAF solution, which we define as the solution with $f \sim 0$ in “most” region of the disk, $f$ differ significantly from its “average” value $\sim 0$ in the innermost region of the disk which is the most important part of an ADAF, increasing from $\sim 0$ to $\sim 1$ when the flow is accreted inward. Thus, in this case the assumption of a constant $f$ may be a dangerous approximation (Esin, McClintock, & Narayan 1997). For the two luminous hot accretion solutions with $\dot{m} > \dot{m}_1$, $f < 0$, as we expected above.

For the energy balance of electrons, from the right plot of Figure 2, we see that for a typical ADAF with $\dot{m} \ll \dot{m}_1$, the radiated energy is larger than the transferred energy from the ions by Coulomb collision, as first pointed out by Nakamura et al. (1997). In this case, the electrons must lose their entropy to balance the radiative loss, like the ions in our new hot solution. However, when $\dot{m}$ approaches or becomes greater than $\dot{m}_1$, we find that $q^\sim \approx q_\alpha$ approximately holds (However, see Figure 5 below).

The long-dashed line shows the solution with $\dot{m} > \dot{m}_2$. In this case, the accretion rate is so large that the hot solution can only exist beyond a certain radius where the Coulomb energy transfer is still weaker than the sum of the viscous dissipation and compression work. Within this transition radius, Coulomb coupling becomes so efficient that the disk can not remain hot. The only viable solution in this case is the cold optically thick disk, therefore the disk will collapse rapidly onto the equatorial plane and form a cold thin annulus. The structure and the location of the transition region are determined by the parameters and the outer boundary conditions of the flows. Since we assume the gas pressure is much larger than the radiation pressure in our equations, our integration can’t continue after the transition from the hot disk to the cold disk to track the cold disk solution because the radiation pressure can not be neglected in that case. In fact, we find that in the transition region the radiative pressure begins to become comparable to the gas pressure.

Significant energy will be released in the transition region, as pointed out by Pringle, Rees, & Pacholczyk (1973). First, the magnetic flux tubes will have to expand out of the plane and strong magnetic reconnection will occur. Second, in the transition process, the accretion flow cools rapidly from a temperature close to virial to a state with a much lower black body temperature, thus the internal energy of the accretion flow will be released at a luminosity...
Luminous hot accretion disks

\[ L \sim \frac{\dot{m} M_{\text{Edd}}}{m_i} kT_i, \]

where \( T_i \) is the ions temperature just outside of the transition radius. Third, the radial velocity of the accretion flow must drop down rapidly in the transition process because of the large radiation pressure gradient force and the centrifugal force. A shock may happen in the transition region if the radial velocity of the accretion flow before the transition is supersonic.

The pre-shock kinetic energy of the accretion flow will be converted into thermal energy and released. This part of energy is at least in the same order with or several times larger than \( L \) in eq. (21), depending on the value of the Mach number of the accretion flow just outside of the transition radius. Investigating where these energy go is obviously a very interesting subject.

The soft black body photons radiated from the inner cold annulus can partially enter the outer hot disk to serve as the seed photons of Comptonization. In our calculation we neglect this cooling effect. We expect in this case the transition radius will move outward slightly but such two-components configuration still holds. One reason is that from Figure 1 we find the scale height of the hot disk just outside of the transition radius is low thus only a small fraction of black body photons can be intercepted by the hot disk. In addition, the significant energy released in the transition region will serve as a heating source of the outer hot disk, which would partially cancel the cooling effects of Comptonization of black body soft photons.

Figure 3 gives some more examples of the luminous hot accretion solutions with lower viscous parameter \( \alpha = 0.1 \) and 0.01. We see from the figure that even when \( \alpha \) is as low as 0.01, the new hot accretion solution still exists for \( \dot{m} = 1 \).

All above results are for a black hole with stellar mass and the outer boundary is small, \( r_{\text{out}} = 100r_g \). The results are qualitatively the same for a black hole with galactic mass and a much larger outer boundary, as shown by Figures 4 and 5 for the dynamics and energy relationships, respectively. The four solutions are for \( M = 10^8 M_\odot \) and \( r_{\text{out}} = 10^4r_g \). The solid, dotted, dashed, and long-dashed lines are for \( \dot{m} = 0.05, 0.1, 0.3 \) and 1. They correspond to an ADAF with \( \dot{m} \sim \dot{m}_1 \), a luminous hot solution with \( \dot{m}_1 < \dot{m} < \dot{m}_2 \), and two luminous solutions with \( \dot{m} > \dot{m}_2 \). Compared to Figures 1 & 2, we find that \( q_{\text{adv}}/q^{+} \)
and \( q^\alpha \) are more sensitive to the radius. Especially, a “peak” arises in each of the four lines in the two plots in Figure 5. This is because the peak in the electron temperature profile leads to a local minimum in \( q_{ic} \). This result indicates that when the outer boundary is far away from the hole, it is hard to make any approximation such as \( f \equiv q_{adv}/q^\alpha \approx \text{const.} \) or \( q^{-} \approx q_{ic} \). Exact self-consistent global solutions are needed to calculate the dynamics and the emergent spectrum.

Throughout the present paper, we only take into account bremsstrahlung, synchrotron, and their Comptonization. For a corona lying above a cold disk, Comptonization of soft photons from the cold disk should also be included. But we can’t conclude that this cooling effect will deduce the maximum accretion rate below which our luminous hot accretion solution exists because some extra heating processes in addition to viscous dissipation, such as magnetic reconnection, would play an opposite role.

The exact values of \( \dot{m}_1 \) and \( \dot{m}_2 \) are hard to determine. One reason is because the advection factor \( f \) is the function of radii. More importantly, they are sensitively dependent on the outer boundary conditions. All the values given below should only be considered as the approximate values. Our numerical calculation results show that when the outer boundary is set at \( 100r_g \), we have \( \dot{m}_1 \approx 0.25\alpha^{-0.7} \) and \( \dot{m}_2 \approx 1.5\alpha^{-0.7} \). The above formula have different dependence on \( \alpha \) compared to those obtained in section 2. One reason is because we set the outer boundary at a relatively small radius, \( 100r_g \). As is well known, the self-similar scaling law \( \rho \propto \alpha^{-1} \) fails and physical quantities such as radial velocity and density are less sensitive to \( \alpha \) in the inner region of the disk. Another reason is due to different viscous descriptions. In our case, the specific angular momentum of the accretion flow decreases rapidly away from the outer boundary, but almost remains constant in the inner region of the disk (see also Abramowicz et al. 1988; Nakamura et al. 1997). Thus the viscosity plays a smaller role compared with the case of the diffusion type viscous description. When the outer boundary is set at \( 10^4r_g \), we find \( \dot{m}_1 \approx 0.05, 10^{-2}, 10^{-4} \) for \( \alpha = 0.1, 0.01, 0.001 \) and \( \dot{m}_2 \approx 5\dot{m}_1 \). The values of \( \dot{m}_1 \) are obviously different with \( \dot{m}_1 \approx \alpha^2 \) in Esin, McClintock, & Narayan (1997) but are similar with Nakamura et al. (1997). The discrepancy is again due to different viscous descriptions.

Up to now, we don’t concern the stability of the solution. From the surface density plots in Figures 1 & 4, we find that the surface density of the disk increases with the increase of the accretion rate, therefore the solutions are viscously stable. The thermal stability analysis is a little difficult. Optically-thin hot accretion flow is normally thermally unstable when the advection in the energy equation is not included in the stability analysis (Piran 1978). Abramowicz et al (1995) and Narayan & Yi (1995) show that the optically-thin hot ADAF is thermally stable because of the inclusion of the energy advection. Our model is along the line of ADAF, with energy advection is explicitly included in the energy equation. So we expect our new hot accretion solution is also thermally stable. More convincing thermal stability analysis is needed and is a subject of a future investigation.

\( \dagger \) In Figure 4 the solid line shows higher surface density compared with the dotted line in some region of the disk although it corresponds to a lower accretion rate. This is because these two solutions correspond to different outer boundary conditions. The specific angular momentum of the flow denoted by the solid line at the outer boundary \( 10^4r_g \) is relatively high thus the accretion is of disk-like. While for the dotted line, the angular momentum is relatively low thus the solution is of Bondi-like characterized by a much larger sonic radius and lower surface density. See Yuan et al. (2000) for details.

© 2000 RAS, MNRAS 000.
In the $\dot{m}$ vs. $\Sigma$ plot, an “S” curve is usually found for an optically thick accretion flow, thus for certain choices of $\dot{m}$, the only solution available is then an unstable one (Abramowicz et al. 1988; Chen & Taam 1993). It is often argued that in this case, the accretion flow will be forced into a limit cycle behavior where the flow oscillates between the two stable solutions. We note that the “unstable range” of $\dot{m}$ is approximately in the range of $0.01 \lesssim \dot{m} \lesssim 1$, similar to the range within which our new solution exists. Therefore, we propose that possibly the accretion flow would enter into the luminous hot solution rather than enter into the limit cycle.

4 SUMMARY AND DISCUSSION

We find a new branch of hot accretion disk solution. This solution corresponds to the accretion rate higher than the critical rate of ADAF $\dot{m}_1$ but less than 1 (all the accretion rates are in units of Eddington accretion rate defined as $10L_{\text{Edd}}/c^2$). In these solutions, the viscous dissipated energy is less than the Coulomb coupling between the ions and electrons and it is the advective heating together with viscous dissipation that balances the Coulomb energy transfer from ions to electrons. When the accretion rate $\dot{m}$ is within the range $\dot{m}_1 \lesssim \dot{m} \lesssim \dot{m}_2$, where $\dot{m}_2 \sim 5\dot{m}_1 < 1$, the hot solution can extend throughout the disk. When $\dot{m}_2 \lesssim \dot{m} \lesssim 1$, within a certain radius of the disk, the Coulomb cooling of the ions is so strong that even the sum of the viscous dissipation and compression work can not balance it, therefore, the hot accretion flow will be cooled by the Coulomb cooling and collapse to form a cold thin annulus.

The new hot solutions are physically quite different with ADAF although the equations describing them are exactly the same. First, they correspond to different ranges of accretion rate. For a standard ADAF, the viscous dissipation rate is significantly greater than Coulomb cooling rate, therefore almost all the viscously dissipated energy is stored in the plasma as entropy. With the increase of the accretion rate, more and more viscously dissipated energy is lost through Coulomb cooling.

When the accretion rate reaches a certain rate, the viscously dissipated energy is balanced by the Coulomb cooling. This accretion rate is determined by the balance between the Coulomb cooling and the sum of viscous dissipation and compression work. In other words, the new critical accretion rate is determined by the balance between the Coulomb cooling and the sum of viscous dissipation and compression work.

The difference between ADAF and my solution can also be understood in the language of entropy. In the standard ADAF, the entropy of the accretion flow increases inward while in my solution the entropy decreases with the decreasing radius. In the standard ADAF, the energy advection serves as a cooling term in the Lagrangian point of view, while in my solution the energy advection plays a heating role. It is the decrease of the entropy of the plasma that partially supply the radiation of the accretion flow. In this case, my solution is dynamically more similar to the cooling flow in elliptical galaxies rather than to the ADAF.

Up to date we have four viable accretion disk models. They are the standard thin disk, slim disk, ADAF, and the luminous hot disk. These four models belong to two series, with the former two models being cold and the latter two being hot. The transition between the two models in each series is due to the variation of the accretion rate. The transition from the standard thin disk to the slim disk corresponds to the transition of the advection factor $f$ from $f \approx 0$ to $f > 0$ when the accretion rate passes across $\sim 1$, while the transition from ADAF to our new hot solution corresponds to the transition from $f > 0$ to $f < 0$ when the accretion rate passes across the critical value $\dot{m}_1$. In addition, the cold series exists for almost any value of accretion rate while the hot series exists only for accretion rate approximately less than the Eddington rate $10L_{\text{Edd}}/c^2$.

Figure 6 shows the corresponding accretion rate and efficiency of these four accretion disk models. For simplicity we assume the efficiency of thin disk is 0.1.

In previous works (Abramowicz et al. 1995; Chen et al. 1995; Narayan & Yi 1995; Esin, McClintock, & Narayan 1997) both the local and global analysis didn’t find the new hot solution. This is because they a priori assumed that the advection factor $f$ must be positive. They neglected the case where $f$ is negative.

The formation of the luminous hot accretion disk has two possible panels. If the initial temperature of the accretion gas at the outer boundary is low, the flow would prefer a cold thin disk first. In this case, the hot disk could be formed through the transition from the thin disk. In the present study we didn’t concern the exact physical mechanism of the transition. The

\*\* So $\dot{m}_1$ should be the accretion rate at which SLE model is most reasonable applicable since $f = 0$ is assumed in SLE. \*\*
general proposals put forward for the thin disk-ADAF transition, such as evaporation (Meyer & Meyer-Hofmeister 1994) and turbulent diffusive heat transport (Honma 1996), are still possible mechanisms. In addition, the secular instability present in the cold disk inner region (Lightman & Eardley 1974) could also be a promising mechanism. As argued by Thorne & Price (1975), this instability could swell this optically-thick radiation-pressure dominated region to a hot, gas-pressure dominated, optically thin region. Since this instability usually appears when the accretion is moderately high, $\dot{m} \gtrsim 0.1$ for example, this hot optically thin region is most possibly described by the luminous hot solution since ADAF doesn’t exist under such high accretion rates.

In this context, we would like to note that when the mass accretion rate is higher than the critical rate of ADAF $\dot{m}_1$ and when the initial temperature of the flow is low at the outer boundary, a sandwich model whereby a corona lies over a cold disk is also viable in addition to our new hot solution. Then a natural question is which of the two solutions would be applicable. Our answer is that it is determined by the details of the mechanism for the formation of the hot solution. If the hot solution is due to the evaporation of the cold disk and if the evaporation efficiency is not high, the sandwich solution will be chosen by the nature. If the hot solution is due to the secular instability, or the turbulent diffusive heat transport, or the evaporation with high efficiency, our new hot solution should be chosen. Note in the sandwich solution, there exists the possibility that the corona might be described by the luminous hot solution rather than ADAF.

If the temperature of the accretion matter at the outer boundary is high, then the hot accretion solution will form from the beginning of accretion. The new hot solutions are possibly very common in the centre of galaxies because: (1) the temperature of gas there is generally as high as $10^6 \sim 10^8 K$; (2) if the practical accretion rate is moderately large and if the viscous parameter $\alpha$ is not very high, this new hot solution is the only feasible hot solution since the critical accretion rate of ADAF, $\dot{m}_1 \sim \alpha^2$ (for diffusion-type viscous description), is very small, while the new hot solution can exist under an accretion rate as high as $\dot{m} \sim 1$ even though $\alpha$ is as low as 0.01 (see Fig. 3). For example, even for $\alpha = 0.1$, an ADAF exists only when $\dot{m} \lesssim 0.01$, while above this accretion rate, $0.01 \lesssim \dot{m} \lesssim 1$, the new hot solution holds.

We didn’t consider the possible existence of outflow in our model. Strong outflows are assumed to be present in ADAF due to its positive sign of the Bernoulli parameter and the inflow-outflow solution (ADIOS) was developed by Blandford & Begelman (1999) (However, see Nakamura 1998; Abramowicz, Lasota, & Igumenshchev 2000). The Bernoulli parameter in our luminous hot solution is found to be in general negative. However, outflows are still possible in our model when the vertical structure of the disk is taken into account, or other factors such as magnetic field is included, as shown by the numerical
simulation of Igumenshchev, Chen, & Abramowicz (1996). In this case, we expect the presence of the transition from the outer new hot solution to the inner ADAF in some situations.

It was recently found that in addition to the self-similar ADAF solution (Narayan & Yi 1994), there is another self-similar CDAF (convection-dominated accretion flow) solution which corresponds to a static envelope in which the mass accretion rate is very small (Narayan, Igumenshchev, & Abramowicz 2000; Quataert & Gruzinov 2000). This solution is possible only when the viscous parameter $\alpha$ and the direction or the efficiency of transport of the angular momentum by convection satisfy some conditions (Narayan, Igumenshchev, & Abramowicz 2000). We don’t know whether our luminous hot solution is relevant to CDAF since we don’t know whether the luminous hot solutions satisfy these two conditions. Almost all the analytical and numerical simulations up to date are for ADAF (e.g., Narayan, Igumenshchev, & Abramowicz 2000, Quataert & Gruzinov 2000; Igumenshchev, & Abramowicz 1999; Igumenshchev, Abramowicz, & Narayan 2000). Our luminous hot accretion solutions are physically quite different with ADAFs in the sense that the entropy of the flow in our solution decreases rather than increases towards the smaller radii. This will have great impact on the convection stability of the flow. Therefore, the extension of study to the new regime is needed.

5 PROMISING APPLICATIONS

Obviously, this new hot solution is characterized by a high efficiency. This together with its corresponding higher accretion rate compared to ADAF, indicates that this type of hot accretion disk could emit a large amount of hard X-ray radiation. This fact plus the arguments given above about the formation channels of the luminous hot accretion solution, make the solution a very promising model for the X-ray luminous sources. An example we note is the very high state of black hole X-ray binaries. Observations indicate that the luminosity of this state can reach close to Eddington luminosity, with the soft blackbody component and the hard nonthermal power-law component being comparable in flux. The power-law component does not show any evidence of a cutoff even out to a few hundred keV. The standard thin disk or the slim disk model can only explain the blackbody component. ADAF, the only viable hot accretion disk model before the present paper, can produce a hard power-law component, but its luminosity is too low compared with the blackbody component because of its low accretion rate and efficiency (Esin, McClintock, & Narayan 1997). Our new hot accretion disk model with $\dot{m} \sim 1$ is a promising model. The optically thick annulus formed within the transition radius is responsible for the soft component in the spectrum of the very high state, while the hard component is emitted by the hot corona above the annulus and the hot disk beyond the transition radius. The high luminosity of the hard component is due to the high efficiency and accretion rate while the extended power law hard tail is because the electron energy distribution has a high energy power law tail as a result of the magnetic reconnection and shock acceleration in the transition region. In addition, another signature of the very high state of BHXBs is their significant variability. The light curve seems to be composed of numerous random flares. This is typically the feature of magnetic reconnection, which must happens in the transition region as we argued in the previous section.

We conjecture that the high state of BHXBs corresponds to the luminous hot disk as well, but its accretion rate is smaller than in the very high state. The accretion rate may correspond to $\dot{m} \gtrsim \dot{m}_2$ in this case. This will result in a moderately lower soft component luminosity and much lower hard component luminosity compared to the very high state. In addition, the luminous hot solution with $\dot{m}_1 \lesssim \dot{m} \lesssim \dot{m}_2$ might corresponds to certain low states in which the X-ray luminosity is moderately high, approaching $\sim 10\% L_{\text{Edd}}$ (Nowak 1995). This is hard to be produced by an ADAF. Even though we assume $\alpha = 0.3$, the highest luminosity of an ADAF is $0.1 \approx \dot{m}_1 \times 10 \times 0.04 \approx 4\% L_{\text{Edd}}$.

On the galactic scale, Seyfert 1 galaxies are in general assumed to be the counterparts of the low state of BHXBs, and Narrow Line Seyfert 1 galaxies correspond to the very high state of BHXBs because of their respective very similar spectra and variability features. It is thus an absorbing and also promising project to work out a unified satisfactory explanation to these sources by our luminous hot accretion disk model. For example, similar with the very high state of BHXBs, the Narrow Line Seyfert 1 galaxies could be interpreted by our new hot accretion model with $\dot{m} \sim 1$. As shown by Mineshige et al. (2000), the extreme soft X-ray excess can be accounted for by the slim disk model with $\dot{m} \sim 1$. The optically thick annulus in our model corresponds to the slim disk in Mineshige et al. (2000). While the hard X-ray component whose flux is comparable to the soft X-ray component (Leighly 1999) can be accounted for by the hot accretion flow outside the annulus and the hot corona above the annulus.

ACKNOWLEDGMENTS

This work is supported by China Postdoctoral Fund and China 973 Project NKBRSF G19990754. I thank Mitch Begelman, Shin Mineshige, Ramesh Narayan, Eliot Quataert, Ronald Taam and Insu Yi for their helpful comments and/or discussions. The hospitality of Korea Institute for Advanced Study where this work began is acknowledged.
REFERENCES

Abramowicz, M.A., Chen, X., Kato, S., Lasota, J.-P., & Regev, O. 1995, ApJ, 438, L37
Abramowicz, M.A., Czerny, B., Lasota, J.P., & Szuszkiewicz, E., 1988, ApJ, 332, 646
Abramowicz, M.A., Lasota, J.-P., Igumenshchev, I. V. 2000, MNRAS, 314, 775
Begelman, M.C., Sikora, M., & Rees, M.J. 1987, ApJ, 313, 689
Blandford, R. D., & Begelman M. C. 1999, MNRAS, 303, L1
Chen, X., Abramowicz, M.A., Lasota, J-P, Narayan, R. & Yi, I. 1995, ApJ, 443, L61
Chen, X., & Taam, R. E. 1993, ApJ, 412, 254
Dermer, C.D., Liang, E.P., & Canfield, E. 1991, ApJ, 369, 410
Esin, A.A., Narayan, R., Ostriker, E., & Yi, I., 1996, ApJ, 465, 312
Esin, A.A., McClintock, J. E., & Narayan, R. 1997, ApJ, 489, 865
Honma, F. 1996, PASJ, 48, 77
Ichimaru, S., 1977, ApJ, 214, 840
Igumenshchev, I. V. & Abramowicz, I. A. 1999, MNRAS, 303, 309
Igumenshchev, I. V., Abramowicz, I. A., & Narayan, R. 2000, ApJ, 537, L27
Igumenshchev, I. V., Chen, X. M., & Abramowicz, M. A., 1996, MNRAS, 278, 236
Kato, S., Fukue, J., & Mineshige, S. 1998, Black-hole Accretion Disks (Kyoto University Press, Kyoto)
Leighly, K. M. 1999, ApJS, 125, 317
Lightman, A.P. & Eardley, D. M. 1974, ApJ, 187, L1
Mahadevan, R. 1997, ApJ, 477, 585
Manmoto, T., Mineshige, S., Kusunose, M. 1997, ApJ, 489, 791
Meyer, F.; Meyer-Hofmeister, E. 1994, A&A, 288, 175
Mineshige, S. Kawaguchi, T., Takeuchi, M., & Hayashida, K. 2000, PASJ, 52, 499
Nakamura, K. E. 1998, PASJ, 50 , L11
Nakamura, K. E., Kusunose, M., Matsumoto, R., & Kato, S. 1997, PASJ, 49, 503
Narayan, R., Igumenshchev, I. V., & Abramowicz, M. A. 2000, ApJ, 539, 798
Narayan, R., Kato, S., & Honna, F. 1997, ApJ, 476, 49
Narayan, R., Mahadevan, R., & Quataert, E. 1998, in “The Theory of Black Hole Accretion Discs”, eds. M.A. Abramowicz, G. Bjornsson, and J.E. Pringle, (Cambridge University Press)
Narayan, R. McClintock, J.E., & Yi, I. 1996, ApJ, 457, 821
Narayan, R., Mahadevan, R., Grindlay, J.E., Popham, R. & Gammie, C., 1998, ApJ, 492, 554
Narayan, R. & Popham 1993, Nature, 362, 820
Narayan, R. & Yi, I. 1994, ApJ, 428, L13
Narayan, R. & Yi, I. 1995, ApJ, 444, 231
Nowak, M.A. 1995, PASP, 107, 1207
Paczynski, B., & Wiita, P. J. 1980, A&A, 88, 23
Piran, T. 1978, ApJ, 221, 652
Pringle, J.E., Rees, M.J., Pacholczyk, A.G. 1973, A&A, 29, 179
Quataert, E., & Gruzinov, A. 2000, ApJ, 539, 809
Rees, M.J., Begelman, M.C., Blandford, R.D., & Phinney, E.S., 1982, Nature, 295, 17
Rybicki, G., & Lightman, A.P. 1979, Radiative Processes in Astrophysics (New York: Wiley)
Shakura, N. I. & Sunyaev, R. A. 1973, A&A, 24, 337
Shapiro, S.L., Lightman, A.P., & Eardley, D. M., 1976, ApJ, 203, 697
Thorne, K.S., & Price, R.H. 1975, ApJ, 195, L101
Yuan, F. 1999, ApJ, 521, L55
Yuan, F., Peng, Q.H., Lu, J.F., & Wang, J.M., 2000, ApJ, 537, 236
Zdziarski, A.A., 2000, Invited review for IAU Symp. 195, Highly Energetic Physical Processes and Mechanisms for Emission from Astrophysical Plasmas, P. C. H. Martens, S. Tsuruta, & M. A. Weber, eds., ASP, pp. 153-170

© 2000 RAS, MNRAS 000.