Advection-Dominated Accretion Flows

Insu Yi\textsuperscript{1,2,3,4}

\textsuperscript{1}Institute for Advanced Study, Princeton, NJ 08540, USA; yi@ias.edu
\textsuperscript{2}Center for High Energy Astrophysics and Isotope Studies, Research Institute for Basic Sciences
\textsuperscript{3}Physics Department, Ewha University, Seoul, Korea; yi@astro.ewha.ac.kr
\textsuperscript{4}Korea Institute for Advanced Study, Seoul, Korea

Abstract. We review basic properties of advection-dominated accretion flows (ADAFs) and their applications to astrophysical systems ranging from Galactic binary systems to galactic nuclei. A new classification scheme for low-luminosity, X-ray bright galactic nuclei is highlighted. Some outstanding unresolved issues are discussed.

1. Introduction

It has recently been recognized that rotating accretion flows with low radiative efficiency are applicable to a wide range of astrophysical systems including Galactic X-ray transients and active galactic nuclei. In the recently discussed models for such sources, the radiative efficiency becomes low because the gravitational binding energy dissipated during infall of accreted matter is not efficiently radiated away due to low gas densities but stored within the flows and radially advected inward. Viscous torque transports angular momentum while generating heat through dissipation. Radiative efficiency is basically determined by the fraction of the viscously dissipated energy which goes into radiation (Narayan & Yi 1995b). In the optically thin limit, which occurs when the density of the accretion flow falls below a certain limit, the radiative efficiency becomes small due to the long time scales for electron-ion energy exchange process and relevant cooling processes. The first discussions of such flows are found in Ichimaru (1977), Rees et al. (1982), and references therein.

The accretion flows around compact objects could be classified into several types (e.g. Narayan et al. 1998b, Frank et al. 1992). (i) Geometrically thin disks radiate away a large fraction of the viscously dissipated energy and due to high density, the optical depth for outgoing radiation is large. The disk temperature is relatively low and hence the internal pressure support is small, which results in the small geometrical thickness (e.g. Frank et al. 1992 for review). In these flows, the cooling rate $Q^-$ is balanced by the heating rate $Q^+$, $Q^+ = Q^-$ while the electron temperature $T_e$ equals the ion temperature $T_i$. (ii) Two-temperature, Shapiro-Lightman-Eardley (1976) type accretion flows also
maintain the energy balance as $Q^+ = Q^-$ while $T_i \sim T_{\text{vir}} \gg T_e$ where $T_{\text{vir}}$ is the usual virial temperature. These flows are thermally unstable (Piran 1978, cf. Rees et al. 1982). (iii) Slim disk or optically thick advection-dominated accretion flows (ADAFs) occur when accretion rates typically exceed the Eddington rate and the optical depth is large. In this limit, photons diffuse out on a time scale longer than the radial inflow time scale (Abramowicz et al. 1988 and references therein). These flows have $Q^+ > Q^-$ and $T_i = T_e$. (iv) Optically thin ADAFs (Ichimaru 1977, Rees et al. 1982, Narayan & Yi 1994, 1995b, Abramowicz et al. 1995) are relevant for substantially sub-Eddington accretion rates. The optical depth in these flows are typically less than unity and the radiative cooling rate is much smaller than the heating rate, i.e. $Q^+ > Q^-$, while the ion temperature $T_i \sim T_{\text{vir}}$ is much higher than the electron temperature $T_e$. These flows are most interesting in systems which have relatively low luminosities and high emission temperatures (Narayan & Yi 1995b).

In this review, we focus on the optically thin ADAFs and discuss their applications to various astrophysical systems.

2. Basics of ADAFs: A Simple Version

2.1. Basic Equations

Following Narayan & Yi (1994, 1995ab), we adopt the following notations for a steady, axisymmetric, rotating accretion flow; $R =$ cylindrical radius from the central star, $\rho =$ gas density, $H =$ thickness or vertical scale height of the flow, $v =$ radial velocity, $\Omega =$ angular velocity, $\Omega_K = (GM/R^3)^{1/2} =$ Keplerian angular velocity, $c_s =$ isothermal sound speed, $\nu =$ kinematic viscosity coefficient, $T =$ temperature of the gas, and $s =$ specific entropy of the gas. Then, the basic conservation equations for mass, radial momentum, angular momentum, and energy become respectively,

\begin{align*}
\rho RH v &= \text{constant}, \\
\frac{v}{R} \frac{dv}{dR} - (\Omega^2 - \Omega_K^2) R &= -\frac{1}{\rho} \frac{d}{dR} (\rho c_s^2), \\
\rho RH \frac{d(\Omega R^2)}{dR} &= \frac{1}{dR} \left( \nu \rho R^3 H \frac{d\Omega}{dR} \right), \\
\rho v T \frac{ds}{dR} &= q^+ - q^- \equiv f q^+ 
\end{align*}

where $q^+ = \nu \rho R^2 (d\Omega/dR)^2 =$ viscous dissipation rate per unit volume, $q^- =$ radiative cooling rate per unit volume, and $\rho v T (ds/dR) = q^{\text{adv}} =$ radial advection rate per unit volume. Therefore, the energy equation becomes simply $q^{\text{adv}} = q^+ - q^- = q^+$ which defines the advection fraction $f$. When advection cooling per unit volume dominates, $q^- \ll q^+, f \approx 1$. Even when $f$ differs substantially from unity, most of dynamical calculations assume a constant $f$ rather than solving the full energy equation (or assumes a simple cooling such as bremsstrahlung). Viscosity coefficient is specified by the $\alpha$ prescription (e.g. Frank et al. 1992) $\nu = \alpha c_s H = \alpha c_s^2 / \Omega_K$ where $\alpha$ is a constant often assumed to be in the range $0.01 - 1$. 2
The accretion flows are classified into three types according to the relative importance of the terms in the energy equation (Narayan et al. 1998). (i) \( q^+ \approx q^- \gg q^{adv} \): energy balance is maintained between viscous heating and radiative cooling, which corresponds to high-efficiency flows such as geometrically thin disks. (ii) \( q^+ \approx q^{adv} \gg q^- \): radiative loss is negligible and luminosity is very low. (iii) \( q^+ \ll q^- \approx -q^{adv} \): viscous heating is negligible and thermal energy of the flow is converted to radiation as in the cooling flows. Infall of matter is driven by pressure loss as gas cools.

2.2. Self-Similar Solution

The dynamical equations derived admit a self-similar solution for \( f = \text{constant} \) as shown in Narayan & Yi (1994) and Spruit et al. (1987);

\[
v = - \left[ \frac{(5 + 2\epsilon')}{3\alpha^2} g(\alpha, \epsilon') \right] \alpha v_K, \tag{5}
\]

\[
\Omega = \left[ \frac{2\epsilon' (5 + 2\epsilon')}{9\alpha^2} g(\alpha, \epsilon') \right]^{1/2} \frac{v_K}{R}, \tag{6}
\]

\[
c_s^2 = \left[ \frac{2(5 + 2\epsilon')}{9\alpha^2} g(\alpha, \epsilon') \right] v_K^2, \tag{7}
\]

where \( v_K = R\Omega_K = (GM/R)^{1/2} \), \( \epsilon' = \epsilon/f \), \( \epsilon = (5/3 - \gamma)/(\gamma - 1) \), and \( g(\alpha, \epsilon') = \left[ 1 + \frac{18\alpha^2}{(5+2\epsilon')^2} \right]^{1/2} - 1 \). The specific heat ratio \( \gamma = 4/3 - 5/3 \).

For \( \alpha = 0.01 - 0.3 \), \( \alpha^2 \ll 1 \) and \( f \approx 1 \) gives

\[
v \approx - \left( \frac{9\gamma - 9}{9\gamma - 5} \right) \alpha v_K, \tag{8}
\]

\[
\Omega \approx \left[ \frac{2(15 - 9\gamma)}{3(9\gamma - 5)} \right] \Omega_K \leq \Omega_K, \tag{9}
\]

\[
c_s^2 \approx \left( \frac{6\gamma - 6}{9\gamma - 5} \right) v_K^2 \tag{10}
\]

The self-similar solution reveals the basic properties of the ADAFs. (i) The radial accretion time scale is much shorter than that of the thin disk. (ii) Sub-Keplerian rotation occurs due to large internal pressure support. (iii) Vertical scale height \( H \sim c_s/\Omega_K \sim R \). Moreover, the positive Bernoulli parameter indicates that ADAFs are prone to outflows although there have not been any self-consistent inflow/outflow solutions (Narayan & Yi 1995a). The vertically integrated equations and the self-similar solutions do not introduce any serious errors in flow dynamics since the height integration is a good approximation (Narayan & Yi 1995a).
2.3. Cooling and Heating Mechanisms

In ADAFs, due to low radiative cooling efficiency of electrons and rather ineffective ion-electron coupling, which is taken to be the Coulomb coupling, ions are nearly virialized (Narayan & Yi 1995b)

\[ T_i \sim 2 \times 10^{12} \beta r^{-1} K. \]  

(11)

Electrons’ energy balance is maintained as (Mahadevan & Quataert 1997)

\[ \rho T_e v \frac{ds}{dR} = \rho v \frac{de}{dR} - kT_e \frac{dn}{dR} = q^{ie} + \delta q^+ - q^- \]  

(12)

where \( kT_e v (dn/dR) = q^{\text{compress}} \) is the compressive heating and \( \delta q^+ \) is the direct viscous heating on electrons with \( \delta \sim m_e/m_p \sim 10^{-3} \) (Nakamura et al. 1996). \( q^{ie} > q^{\text{compress}} \) for \( \dot{m} \gtrsim 0.1 \alpha^2 \) and \( \rho v (de/dR) \approx q^{ie} - q^- \) describes the electron energy balance. \( q^{ie} < q^{\text{compress}} \) occurs when \( \dot{m} \lesssim 10^{-4} \alpha^2 \) and the energy balance for electrons becomes \( \rho v (de/dR) \approx q^{\text{compress}} \), which is appropriate only for very low luminosity systems. \( \delta q^+ < q^{\text{compress}} \) is realized only for \( \delta \lesssim 10^{-2} \) which implies that the direct viscous heating is uninteresting in most practical cases (e.g. Mahadevan & Quataert 1997).

For convenience, we introduce some physical scalings; mass \( m \equiv M/M_\odot \), radius \( r = R/R_s \) (\( R_s = 2GM/c^2 = 2.95 \times 10^5 m \) cm), accretion rate \( \dot{m} = \dot{M}/\dot{M}_{\text{Edd}} \) (\( \dot{M}_{\text{Edd}} = L_{\text{Edd}}/0.1 c^2 = 1.39 \times 10^{18} m g/s \) where \( L_{\text{Edd}} \) is the Eddington luminosity). The equipartition magnetic field \( B^2/8\pi = (1 - \beta)pc_s^2 \) with \( \beta = 0.5 \) and \( f = 1 \) where \( \beta \) is the ratio of magnetic to total pressure. The self-similar solution gives the physical quantities in terms of the dimensionless variables defined here.

\[ v \approx -1 \times 10^{10} \alpha r^{-1/2} \text{ cm/s}, \]  

(13)

\[ \Omega \approx 3 \times 10^4 \alpha^{-1} r^{-3/2} \text{ s}^{-1}, \]  

(14)

\[ c_s^2 \approx 1 \times 10^{20} r^{-1} \text{ cm}^2 / \text{s}^2, \]  

(15)

\[ n_e \approx 6 \times 10^{19} \alpha^{-1} m^{-1/2} r^{-3/2} \text{ cm}^{-3}, \]  

(16)

\[ B \approx 8 \times 10^8 \alpha^{-1/2} m^{-1/2} r^{-5/4} \text{ G}, \]  

(17)

\[ \tau_{es} \approx 24 \alpha^{-1} r^{-1/2} \dot{m}, \]  

(18)

\[ q^+ \approx 5 \times 10^{21} m^{-2} r^{-4} \text{ erg/s/cm}^3, \]  

(19)

where \( n_e \) is the electron number density and \( \tau_{es} \) is the electron scattering depth.

Various (electron) cooling processes give rise to distinct spectral components (Narayan & Yi 1995b). The integrated cooling rate \( Q = \int dV q \) where \( q \) is the cooling rate per unit volume and \( \int dV \) denotes integration over the entire accretion flow. The total electron cooling rate is

\[ Q_e^- = Q_{\text{sync}} + Q_{\text{Compt}} + Q_{\text{brem}} \]  

(20)

where \( Q_{\text{sync}} \) is the synchrotron cooling rate which gives rise to spectral emission components in radio, IR, or optical/UV depending on the mass \( m \) and accretion...
rate $\dot{m}$. $Q_{\text{Compt}}$ is the Compton cooling rate which is mainly responsible for optical/UV/soft X-ray emission. $Q_{\text{brems}}$ is the bremsstrahlung cooling contributing to X-ray and soft gamma-ray emission. If ADAFs in the inner regions around accreting compact objects are surrounded by the optically thick disks, optical/UV emission from cool disks is expected. Energetic protons in ADAFs may result in high energy gamma-rays (Mahadevan et al. 1997). Similar radiation processes in zero angular momentum spherical accretion have been extensively discussed (e.g. Melia 1992 and references).

### 2.4. Critical Quantities

ADAFs exist when accretion rates fall below a certain critical rate $\dot{M}_{\text{crit}} = \dot{m}_{\text{crit}} \dot{M}_{\text{Edd}}$. Such a critical rate arises because there exists a maximum accretion rate above which heating could be balanced by radiative cooling without any necessity of advective cooling (Rees et al. 1982, Narayan & Yi 1995b, Abramowicz et al. 1995, Narayan et al. 1998). In the case of the single temperature case, i.e. $T_e = T_i \propto r^{-1}$, assuming an optically thin, bremsstrahlung cooling (good for $r > 10^3$), we get $q^+ \propto m^{-2}r^{-4}\dot{m} \propto \dot{m}$ and $q^- = q_{\text{brems}} \propto \rho^2 T_e^{1/2} \propto \rho^2 \propto \alpha^2 r^{-1/2}$. In the case of the single temperature with $T_e = T_i \neq T_{\text{vir}}$, assuming synchrotron and Compton cooling, the critical accretion rate becomes $\dot{m}_{\text{crit}} \propto 10^{-4}\alpha^2$. These well motivated critical rates are too low to be of practical interest (Esin et al. 1997). Esin et al. (1997) found that the critical rate deduced from observed spectral transition in soft X-ray transients is much higher than the above rates.

In the two-temperature ADAFs, the bottleneck in energy transfer from ions to electrons define another critical rate which is good for $r \lesssim 10^3$ (e.g. Narayan et al. 1998). That is, using $q^+ \propto \dot{m}$ and $q^- \propto q_{\text{sync}}^{\text{iv}} \propto \dot{m}^2$ and equating the two rates, $q^+ = q_{\text{sync}}^{\text{iv}}$, gives a critical accretion rate $\dot{m}_{\text{crit}} \approx 1 \times 10^3 (1 - f) e^2 \beta^{-1/2} \alpha^{-1/2} / \dot{M}$. Alternatively, $t_{\text{iv}} = t_{\text{acc}} \approx R / v$ gives $\dot{m}_{\text{crit}} \approx 0.3\alpha^2$. In sum, $\dot{m}_{\text{crit}} \propto \alpha^2$ for $r \lesssim 10^3$ and $\dot{m}_{\text{crit}} \propto \alpha^2 r^{-1/2}$ for $r \gtrsim 10^3$, which depicts the actual radial dependence of the critical accretion rate.

It is interesting to point out that there exists a critical $\alpha$ (Chen et al. 1995). For $\alpha < \alpha_{\text{crit}} \sim r$, $\dot{m}_{\text{crit}}$ exists while for $\alpha > \alpha_{\text{crit}} \sim r$, $\dot{m}_{\text{crit}}$ doesn’t exist.

### 2.5. Some Recent Works on Heating Ions

Bisnovatyi-Kogan & Lovelace (1997) recently claimed that large electric fields parallel to magnetic fields can accelerate electrons and hence bypassing the bottleneck in energy transfer from ions to electrons, which could rule out ADAFs as a possible accretion flow type. However, such a possibility is realized only when a substantial resistivity on microscopic scale exists. This requires a small magnetic Reynolds number. Their proposal to use the macroscopic turbulent resistivity is not applicable on microscopic scales as pointed out by Blackman (1998) and Quataert (1998). Blackman (1998) argued that the Fermi acceleration by large scale magnetic fluctuations associated with MHD turbulence may lead to preferential ion heating and hence two-temperature plasma. This heating is not applicable to non-compressible Alfvenic turbulence which is most likely to be more important than the compressive mode. For weak magnetic fields substan-
tially weaker than equipartition fields, Alfvénic component of MHD turbulence is dissipated on scales of proton Larmor radii (Gruzinov 1998, Quataert 1998). This mechanism favors ion heating and two-temperature plasma. For strong fields (i.e. near equipartition), the Alfvénic turbulence cascades to scales much smaller than proton Larmor radii and can directly heat electrons, which could cast doubt on ADAFs with equipartition strength magnetic fields. That is, there is a possibility that equipartition plasma doesn’t allow two-temperature plasma. This issue may ultimately be resolved by observed spectra.

2.6. ADAF Luminosity

In ADAFs, the observed radiative luminosity $L_{\text{ADAF}} = \int L(E)dE \sim \int q^{\epsilon}dV$ or in terms of the ADAF radiative efficiency $\eta_{\text{ADAF}}$, $L_{\text{ADAF}} = \eta_{\text{ADAF}} M c^2$ where $\eta_{\text{ADAF}} = \eta_{\text{eff}} \times 0.2 \dot{m} \alpha^{-2} \propto \dot{m} \propto M/M$ (Narayan & Yi 1995b). In contrast, the thin disk luminosity $L_{\text{thin disk}} \sim \eta_{\text{eff}} M c^2$ with $\eta_{\text{eff}} \sim 0.1$ (e.g. Frank et al. 1992).

3. Some Physical Issues

3.1. Global Solutions

The self-similar solution considered so far applies to regions far from the inner and outer boundaries where physical scales of the system demand the dynamical equations to deviate from self-similarity. A comprehensive summary of the Newtonian case is found in Kato et al. (1998). The task of finding the global solution is to find an eigenvalue $j$ (specific angular momentum accreted by the central object) with proper boundary conditions. The proper boundary conditions are (i) outer thin disk matching inner ADAF, (ii) sonic point for the transonic ADAF, and (iii) vanishing torque at the inner boundary. Solving the radial momentum equation, angular momentum equation, and energy equation, along with the implicit continuity equation for $v$, $\Omega$, $c_s$, $\rho$ and eigenvalue $j$, global solutions are found. The self-similar solution is a good approximation for a wide range of radii between the inner and outer boundaries. The pseudo-Newtonian potential case has been solved by Matsumoto et al. (1985), Narayan et al. (1997), Chen et al. (1997). The major findings are as follows. (i) For $\alpha \lesssim 0.01$, inefficient angular momentum transport results in slow radial accretion and a wide radial zone of super-Keplerian rotation. In the super-Keplerian rotation region, a thick torus-like structure with funnel around the rotation axis forms and the pressure profile shows a maximum, which is reminiscent of the ion torus model (e.g. Rees et al. 1982). (ii) For $\alpha \gtrsim 0.01$, efficient angular momentum transport and rapid accretion $v \propto \alpha v_K$ occur. There exist no pressure maximum and the accretion flows are quasi-spherical.

In the relativistic case with a spinning black hole (Abramowicz et al. 1996, Peitz & Appl 1997, Gammie & Popham 1998, Popham & Gammie 1998), the Newtonian and pseudo-Newtonian results are largely confirmed for $r \gtrsim 10$. For $r \lesssim 10$, however, significant differences and spin effects are seen. Detailed calculations of emission spectra have been carried out by Jaroszynski & Kurpiewski (1997).
It has been an issue whether shocks form in transonic accretion flows. In the steady calculations, shocks are not seen and in the time-dependent calculation of Igumenshchev et al. (1996), no shocks have been seen. Mannoto et al. (1996)'s time-dependent calculation shows shock-like steepening in density waves, which could be responsible for ADAF variabilities.

3.2. Stability

The geometrically thin disks could be thermally and viscously unstable under certain circumstances (e.g. Frank et al. 1992). ADAFs are primarily stable against thermal and viscous perturbations: (i) ADAFs are stable against long wavelength perturbations (Narayan & Yi 1995b, Abramowicz et al. 1995). (ii) ADAFs are marginally stable against short wavelength perturbations in the single temperature case (Kato et al. 1996, 1997, Wu 1997). (iii) ADAFs are stable against short wavelength perturbations both thermally and viscously in the two temperature case (Wu & Li 1996). The stability of ADAFs have led to an argument that the thermally unstable Shapiro-Lightman-Eardley disk (Piran 1978) is a spatially transitional accretion flow linking the thermally unstable outer thin disk to the stable inner ADAF. In the time-dependent calculations, small wavelength perturbations in the single temperature ADAFs have been observed to grow while moving inward. The growth rate is however not rapid enough to affect the steady global structure (Mannoto et al. 1996).

4. X-ray Transients

4.1. Black Hole Systems

Detailed spectral fitting has been tried for black hole systems such as A0620-00, V404 Cyg (Narayan et al. 1996), Nova Mus 1991, Cyg X-1, GRO J0422+32, GRO J1719-24 (Esin et al. 1997), and 1E1740.7-2942 (Vilhu et al. 1997). The main physical parameters used in the spectral fitting are $M$, $\dot{m}$, $\alpha$, $\beta$, and $r_{tr}$, where the last one is the radius of accretion flow transition from outer thin disk to inner ADAF. The spectral fitting assumes that (i) the outer thin disk is joined to the inner ADAF and (ii) the outer thin disk is unstable against disk instability. The instability causes the heating/cooling waves to propagate inward, resulting in delays between different emission components (Lasota et al. 1996, Hameury et al. 1997). The transition radius $r_{tr}$ is crucial in determining spectra (Lasota et al. 1996, Narayan et al. 1996,1998, Esin et al. 1998).

4.2. Thin Disk - ADAF Transition

ADAFs exist when $q^- < q^+$ which in principle determines $r_{tr}$. In the single temperature, bremsstrahlung dominated case,

$$q_{vis}^+ \propto m^{-2} \dot{m}^{-4} r^{-4} \propto r^{-4}$$

$$q_{brem}^- \propto \rho^2 T^{1/2} \propto \alpha^{-2} m^{-2} \dot{m}^{-7/2} r^{-7/2}$$

and $q^+ = q^-$ gives (Honma 1996)

$$r_{tr} \sim 3 \times 10^2 (\alpha^4/\dot{m}^2).$$
In more realistic cases, details of transition are unknown. For instance, Honma (1996) considers the radial turbulent heat diffusion near the interface between the thin disk and the ADAF. Spectral fitting gives a different result on $r_{tr}$ (Esin et al. 1997). The issue of the accretion flow transition still remains unresolved.

4.3. Neutron Star Systems

Neutron star transients have not been well studied using the ADAF models mainly due to the lack of sufficiently well-developed physical understanding of accretion flows near the neutron star surface. The radiation feedback from the soft radiation emitted by the stellar surface makes the spectral calculations considerably more complicated (Narayan & Yi 1995b, Yi et al. 1996). Moreover, it is unclear whether the spectral transition similar to that of black hole systems exists in neutron star systems.

4.4. Accretion-Powered X-ray Pulsars

Spectral transition in neutron star systems could be potentially very difficult to detect if thermalization near the neutron star surface occurs rapidly (e.g. Yi et al. 1996). Accretion-powered X-ray pulsars could however provide an observable signature of the temporal accretion flow transition. In accretion-powered pulsars such as 4U 1626-67, GX 1+4, and OAO 1657-415, abrupt torque reversals have been observed, which are extremely difficult to explain within the conventional models involving disk-magnetospheric interaction (Chakrabarty et al. 1997). The difficulties arise mainly due to (i) short reversal time scales, (ii) nearly identical spin-up and spin-down rates, (iii) small X-ray luminosity changes, (iv) significant spectral transition reported in 4U 1626-67 (Vaughan & Kitamoto 1997), and (v) torque-flux correlation observed in GX 1+4 (Chakrabarty et al. 1997).

The observed phenomenon is well accounted for by the accretion disk transition triggered by a gradual, small amplitude modulation of mass accretion rate. The sudden reversals occur at a rate $\sim 10^{16-17} g/s$ when the accretion flow makes a transition from (to) a primarily Keplerian flow to (from) a substantially sub-Keplerian, radially advective flow (Yi et al. 1997, Yi & Wheeler 1998). The proposed transition model naturally shows that (i) the transition time scale is likely to be shorter than days, (ii) the required accretion rate change is at the level of a few $\times 10$ percent, and (iii) the abrupt reversal is a signature of a pulsar system near spin equilibrium with small mass accretion rate modulations near the critical accretion rate on the time scale of years. Other possible explanations for the spectral transition and the torque-flux correlation have been suggested (Nelson et al. 1997, van Kerkwijk et al. 1998) with varying difficulties in explaining the above observational facts. The accretion flow transition is similar to those in black hole transients and cataclysmic variables, which strongly suggests a common physical origin (Yi et al. 1997). The transition time scale is $t_{thermal} \sim (\alpha \Omega_K)^{-1}$ or $t_{vis} \sim R/\alpha c_s \sim (R/H) t_{thermal}$, where the latter time scale becomes $\sim 10^3 s$ for $\alpha \sim 0.3$, $R \sim 10^9 cm$, $\dot{M} \sim 10^{16} g/s$. This time scale is short enough to induce transition on a time scale of a day or less.

If the accretion flow’s thermal pressure is a substantial fraction $\xi$ of the dynamical pressure, i.e. $c_s^2 \sim \xi R^2 \Omega_K^2$, the thickness of the flow $H \sim \xi^{1/2} R$, the radial accretion speed $v_R \sim -\alpha \xi H \Omega_K$, and the angular rotational frequency
\[ \Omega \sim (1 - 5\xi/2 - \alpha^2\xi^2/2)^{1/2} \Omega_K = A\Omega_K \] where \( \xi \to 0 \) is the Keplerian limit and \( A < 1 \) for ADAFs and \( A = 1 \) for thin Keplerian disks. When \( A < 1 \) occurs, \( R'_c = A^{2/3}R_c \), where \( R_c \) is the Keplerian corotation radius and \( R'_c \) is the sub-Keplerian corotation radius. The accretion flow is truncated at a radius \( R'_o \) and using \( N'_0 = AM(GM_*/R'_o)^{1/2} \) the torque on the star becomes

\[ \frac{N'_o}{N'_0} = \frac{71 - (8/7)(R'_o/R'_c)^3}{6A} \tag{24} \]

which pushes the pulsar’s spin toward an equilibrium spin \( P'_e = P_e/A > P_e \) where \( A \) denote quantities after transition from thin disk to ADAF. In this picture, torque reversal is expected if

\[ B_\ast \sim 5 \times 10^{11} L_{x, 36} \frac{1}{P_{\ast, 10} G} \tag{25} \]

where \( L_{x, 36} = L_x/10^{36} \text{erg/s} \) is the X-ray luminosity and \( P_{\ast, 10} = P_\ast/10s \) is the pulsar spin period. Observed quasi-periodic oscillation (QPO) periods tightly constrain the proposed model. Yi & Grindlay (1998) discusses some possible implications on spin-up of LMXBs to MSPs when the accretion flows are ADAFs in these systems.

### 4.5. Energetic Protons: Lithium Production

Energetic ions present in ADAFs are capable of nuclear spallation. Lithium production in ADAFs in binary systems and tori in galactic nuclei has been discussed by Ramadurai & Rees (1985), Jin (1990), Yi & Narayan (1997). Using the self-similar solution, the relevant physical quantities are as follows: number density of protons \( n_H \sim 6 \times 10^{20} m^{-1} \text{cm}^{-3} \), number density of \( \alpha \) particles \( n_\alpha \sim 5 \times 10^{19} m^{-1} \text{cm}^{-3} \), energy per nucleon \( E \sim 300 \text{MeV} \), and the radial accretion speed \( v_R \sim 2 \times 10^9 \text{cm/s} \). The production of \( ^7\text{Li} \) dominates and the production cross-section \( \sigma_+ (E) \sim 100(E/100 \text{MeV})^{-2} \text{mbarn} \) for \( E \geq 8.5 \text{MeV} \). Continuous production of \( ^7\text{Li} \) within the accretion flow leads to increase in Lithium abundance according to

\[ \frac{\Delta n_{\text{Li}}}{n_H} = \frac{1}{2} \sigma_+(E) v_\gamma \frac{n_\alpha^2}{n_H} \Delta t_{\text{flow}} \tag{26} \]

where \( \Delta t_{\text{flow}} = \Delta R/v_R \). The enrichment results in the terminal abundance \( n_{\text{Li}}/n_H = \int d(n_{\text{Li}}/n_H) \approx 0.1 \) or the total Lithium mass production rate \( \dot{M}_{\text{Li}} = 2 \times 10^{-8} m_\odot \text{yr}^{-1} \).

It is expected that \( ^7\text{Li} \) enrichment occurs around NSs and BHs containing ADAFs, which has been seen in recent precision spectroscopic measurements of \( ^7\text{Li} \) in V404 Cyg, A0620-00, GS 2000+25, Nova Mus 1991, Cen X-4 (Martin et al. 1994 and references therein). In contrast, WD systems show no such effect, which indicates that only at \( r \sim 1 \) relativistic energies are reached while \( r \sim 10^3 \) is the inner most radius in the WD systems, which gives \( E \ll 1 \text{MeV} \).

### 4.6. Thermalization of Particles

Although we have adopted the thermal temperatures for ions and electrons, it is not proven that particle energy distributions are adequately approximated by
thermal distributions. Recent investigations (e.g. Mahadevan & Quataert 1997, Quataert 1998, Blackman 1998) have suggested some limited information on this unresolved issue.

In most of the ADAF models, protons and ions are energized by viscous heating which is mostly unspecified. Alfvénic turbulence (which does not result in strongly non-thermal distributions) and Fermi acceleration (leading to power-law tails) have been considered. Coulomb collisions and synchrotron absorption do not lead to rapid thermalization of protons. So far it appears that the acceleration mechanism itself determines the proton energy distributions. Thermalization of electrons could occur more easily (Ghisellini & Svensson 1991). Coulomb collisions can thermalize electrons for \( \dot{m} \gtrsim 10^{-2} \alpha^2 \). Synchrotron self-absorption leads to thermalization when \( \dot{m} \gtrsim 10^{-5} \alpha^2 r \). The electron energy distributions directly affect radio emission spectra.

Recently Mahadevan (1998) pointed out that low frequency \( \nu \lesssim 10^9 \text{Hz} \) radio spectrum of Sgr A* could be contributed by electrons and positrons produced by charged pions from proton-proton collisions and that neutral pion production could account for tentative detection of gamma-rays in the direction of Sgr A*. In both cases, particle distributions need to be strongly non-thermal and the results depend very sensitively on high energy tails.

5. Galactic Nuclei

ADAFs may exist in galactic nuclei including the Galactic center source Sgr A*. For our discussions, we define \( m_7 = m/10^7 \), \( \dot{m}_{-3} = \dot{m}/10^{-3} \), and \( R_s = 2GM/c^2 = 3 \times 10^{12} m_7 \text{cm} \).

When the accretion flow is a thin disk, the luminosity \( L = \eta \dot{M} \alpha c^2 \) with the efficiency \( \eta = \eta_{eff} \sim 0.1 \) (e.g. Frank et al. 1992). Although the total luminosity is high, the emission temperature or the disk temperature \( T_{\text{disk}} \sim 6 \times 10^6 m_7^{-1/5} \dot{m}_{-3}^{3/10} r^{-3/4} K \) is too low to account for X-ray emission. The disk luminosity \( L_{\text{disk}} \sim 1 \times 10^{42} m_7 \dot{m}_{-3} \text{erg/s} \) is expected to occur in optical/UV/soft X-ray. X-ray emission and radio emission are usually accounted for by optically thin corona and radio jets. For the latter, the estimated radio power is \( L_{\text{jet}} \sim 1 \times 10^{42} \tilde{a}^2 \eta_{\text{jet}} m_7 \dot{m}_{-3} \text{erg/s} \) where \( \tilde{a} \) is the black hole spin parameter and \( \eta_{\text{jet}} \) is the jet radiative efficiency.

In ADAFs, relevant physical quantities scaled for galactic nuclei are equipartition magnetic field \( B \sim 1 \times 10^4 m_7^{-1/2} \dot{m}_{-3}^{1/2} r^{-5/4} G \), electron scattering depth \( \tau_{\text{es}} \sim 5 \times 10^{-2} \dot{m}_{-3} \), the ion temperature \( T_i \sim 2 \times 10^{12} r^{-1} K \), and the electron temperature \( T_e \sim 5 \times 10^9 K \). ADAFs are expected to exist when the mass accretion rate falls below \( \dot{m}_{\text{crit}} = \dot{M}_{\text{crit}} / \dot{M}_{\text{Edd}} \approx 0.3 \alpha ^2 \sim 10^{-3} - 10^{-2} \). Radio emission is easily explained by the synchrotron emission with the characteristic synchrotron emission frequency (Yi & Boughn 1998ab)

\[
\nu_{\text{sync}} \sim 1 \times 10^{12} m_7^{-1/2} \dot{m}_{-3}^{1/2} r^{-5/4} T_{e9}^2 \text{Hz} \tag{27}
\]

where \( T_{e9} = T_e / 10^9 K \sim 5 \). The highest synchrotron radio emission frequency is \( \nu_{\text{max}} = \nu_{\text{sync}}(r \sim 1) \sim 3 \times 10^{13} m_7^{-1/2} \dot{m}_{-3}^{1/2} \text{Hz} \) which comes from the inner most
region of ADAF near the black hole horizon. The radio luminosity

\[ L_R \sim \nu L^\text{sync}_\nu \sim 2 \times 10^{32} x_{M3}^{8/5} T_{e9}^{21/5} m_7^{-6/5} m_3^{-4/5} \nu_{10}^{7/5} \text{ erg/s} \]  

(28)

where \( x_{M3} = x_M/10^3 \) is a dimensionless synchrotron self-absorption parameter and \( \nu_{10} = \nu/10^{10} \text{ Hz} \). ADAFs predict such a relation - frequency relation. ADAFs predict such a relation

\[ \theta(\nu) \sim 2 m_7^{3/5} m_3^{-2/5} \nu_{10}^{-4/5} (D/10\text{Mpc})^{-1} \mu\text{as} \]  

(29)

where the angular size \( \theta(\nu) \) becomes \( \sim \text{mas} \) for distance scales \( D \sim 10\text{kpc} \).

Optical/UV/X-ray in ADAFs arise from inverse Compton scattering of radio synchrotron photons and hard X-rays are from bremsstrahlung and multiple Compton scattering. For \( \dot{m} < 10^{-3} \), X-ray emission is dominated by the bremsstrahlung emission, \( L_x \sim L^\text{brem}_x \sim \dot{m} \nu^2 \) and the radio luminosity

\[ L_R \propto m_7^{8/5} m_3^{6/5} \nu, \]  

gives the luminosity \( L^\text{COMP}_x \propto \dot{m}^{-7/5+N} \) with \( N \geq 2 \). In this case, we expect

\[ L_R \propto m_3^{8/5} L_x^{3/5} \]  

Yi & Boughn (1998ab) have derived and tested the radio/X-ray luminosity relation for ADAFs using \( L_x = L_x(2 - 10\text{keV}); \)

\[ L_R \sim 10^{36} m_7 (\nu/15\text{GHz})^{7/5} (L_x/10^{40} \text{erg/s})^x \text{ erg/s} \]  

(30)

where \( x \sim 1/5 \) for \( \dot{m} \lesssim 10^{-3} \) and \( x \sim 1/10 \) for \( \dot{m} > 10^{-3} \) or \( L_{R,adv}/L_{x,adv} \propto mL_x^{-1} \).

ADAFs are likely to drive jets/outflows (Narayan & Yi 1995a). If jets are powered by black hole’s spin energy (e.g. Frank et al. 1992),

\[ L_{R,jet}/L_{R,adv} \sim 4 \times 10^5 \bar{a}^2 \eta_{jet} m_7^{-1/5} m_3^{1/5} \]  

(31)

and \( L_{R,jet}/L_{x,adv} \propto \bar{a}^2 mL_x^{-1} \) are expected. That is, \( L_{R,jet} \gg L_{R,adv} \) for \( \bar{a} \gg 2 \times 10^{-3} \eta_{jet}^{-1/2} \). It is often argued that radio-loud nuclei have \( \bar{a} \lesssim 1 \). If a galactic nucleus contains a thin disk and a jet \( L_{R,jet}/L_{x,disk} \sim \bar{a}^2 \epsilon_{jet}/\eta_{eff} \sim O(1) \) is expected.

Characteristic ADAF emission spectra are determined primarily by \( \dot{m} \) and weakly affected by the black hole mass \( M \). Any combinations among \( L_x, L_R \) and \( M \) give useful information on the nature of emission from galactic nuclei.

5.1. Galactic Center Source Sgr A*

Galactic center radio source Sgr A* is a prime candidate for an ADAF (Rees 1982, Narayan et al. 1995, 1998, Mannoto et al. 1997). Sgr A*, which is at the dynamical center of the Galaxy, appears to contain a massive black hole with mass \( \sim (2.5 \pm 0.4) \times 10^6 M_\odot \). Wind accretion from nearby IRS 16 wind is expected to provide a mass accretion rate \( \dot{M} \gtrsim \text{a few} \times 10^{-6} M_\odot/\text{yr} \). With the conventional \( \sim 10\% \) efficiency, such a high accretion rate would correspond to a luminosity of \( 0.1 M c^2 \gtrsim 10^{30} \text{erg/s} \) which is some 3 to 4 orders of magnitude larger than the observed radio to gamma-ray luminosity of \( \lesssim 10^{37} \text{erg/s} \). Moreover, a standard thin disk would give peak emission in near infrared but
Menten et al. (1997)'s 2.2 micron upper limit rules out this possibility. These facts strongly suggest that an ADAF is present in Sgr A*.

Spectral fitting based on the observed $M$ and the estimated $\dot{m}$ adequately accounts for the observed emission seen from radio to hard X-ray. Bower & Backer (1998) measured the intrinsic source size $\lesssim 0.48$ mas at 7mm, which at 8.5kpc gives the linear size of 4.1AU and the lower limit on the brightness temperature $4.9 \times 10^9 K$. Both the size and temperature measurements are consistent with the ADAF models. Inverted radio spectrum with sharp cutoff along with no jet like elongations near Sgr A* are also consistent with the ADAF predictions. X-ray constraints are rather uncertain. ROSAT 0.8-2.5 keV luminosity is $\sim 1.6 \times 10^{34} erg/s$ (Predehl & Trumper 1994) and ASCA 2-10 keV luminosity is $\leq 4.8 \times 10^{35} erg/s$ (Koyama et al. 1996). Both constraints are easily satisfied by an ADAF.

5.2. NGC 4258

NGC 4258 almost certainly has a central black hole with mass $M = 3.5 \pm 0.1 \times 10^7 M_\odot$, which is concentrated within 0.13pc (for a distance of 6.4Mpc) of the dynamical center. This source has an observed 2-10 keV X-ray luminosity of $L_X = 4 \times 10^{40} erg/s$ (Makishima et al. 1994). Optical luminosity $L_{opt} \lesssim 10^{42} erg/s$ (Wilkes et al. 1994) provides an additional constraint. 22 GHz continuum emission (after subtraction of jet component) has not been detected with a $3\sigma$ upper limit of 220 $\mu$Jy (Herrnstein et al. 1998) or a luminosity upper limit $L_R(22GHz) < 2.4 \times 10^{35} erg/s$. The non-detection of the core at 22 GHz could imply that $\dot{m} \sim 10^{-2}$ and $r_t \sim 30$ if an ADAF exists in NGC 4258 (Gammie et al. 1998, cf. Lasota et al. 1996). Such a constraint is highly suspect due to a possibility of strong variabilities in radio emission from ADAFs (Blackman 1998, Ptak et al. 1998, Herrnstein et al. 1998).

5.3. M60, M87, NGC1068, and M31

Di Matteo & Fabian (1997b) has attempted to fit M60 emission spectra with an ADAF under the assumption that the accretion rate is close to the Bondi accretion rate with $M \sim 10^9 M_\odot$. Due to uncertainties and lack of flux measurements, a definitive conclusion as to whether an ADAF exists needs more data in X-ray or other wavebands. Reynolds et al. (1996) claimed that an ADAF model spectrum for a blackhole mass $M = 3 \times 10^9 M_\odot$ and the accretion rate $\dot{m} \sim 10^{-3}$ accounts for the observed fluxes of M87. However, such a conclusion is highly suspect because the extended radio emission component has not been properly removed. ADAFs themselves do not produce extended radio emission. NGC1068 is an obscured Seyfert with a very high obscuration-corrected X-ray luminosity. X-rays are most likely to be dominated by scattering and the X-ray luminosity may be too bright for ADAFs based on the estimated black hole mass $\sim a few \times 10^7 M_\odot$ (Yi & Boughn 1998ab). Yi & Boughn (1998b) also considered M31 which has a central black hole of mass $M = 3 \times 10^7 M_\odot$. The observed radio luminosity is too low for the observed X-ray luminosity if an ADAF is assumed around the black hole. It is highly likely that the X-ray luminosity is dominated by binary sources in the nucleus while the radio emission may be due to a very weak ADAF.
5.4. X-ray Bright Galactic Nuclei

Yi & Boughn (1998ab) and Franceschini et al. (1998) applied the ADAF model to a small sample of X-ray bright galactic nuclei which have black hole mass estimates. Since ADAFs are most relevant for low luminosity, hard X-ray sources, faint, hard X-ray galactic nuclei are likely hosts of ADAFs (Fabian & Rees 1995, Di Matteo & Fabian 1997a, Yi & Boughn 1998ab). Hard spectrum, faint X-ray sources could contribute significantly to the diffuse X-ray background. 50% of 2-10 keV XRB could be accounted for by ADAF sources with the comoving density of $3 \times 10^{-3}$ Mpc$^{-3}$ for $L_x \sim 10^{41}$ erg/s, which is comparable to the local density of the $L_*$ galaxy (Di Matteo & Fabian 1997a, Yi & Boughn 1998a). However, unless $M > 10^9 M_\odot$, $L_x > 10^{40}$ erg/s would be already too bright for ADAFs to account for the observed X-ray background. This is because at high luminosities X-ray emission is dominated by the Compton scattering which result in X-ray spectra much different from the background spectrum similar to the bremsstrahlung-dominated X-ray spectra. That is, in relatively bright ADAFs, high $\dot{m}$'s correspond to the Compton-dominated cooling regime. Although a significant clumping of ADAF gas would enhance bremsstrahlung over Compton (Di Matteo & Fabian 1997), such a possibility is difficult to realize.

So far, we have argued that ADAF models in galactic nuclei are testable due to distinguishing characteristics of ADAFs. ADAFs with low radiative efficiency and high temperature are likely for massive black holes accreting at accretion rates $\lesssim 10^{-2} M_{\text{Edd}}$. Hard X-ray emission and inverted spectrum radio emission from compact core are expected. There exists a characteristic radio/X-ray luminosity relation as shown by Yi & Boughn (1998ab). For known black hole masses, existence of hot ADAFs can be tested by radio/X-ray observations. Black hole masses could be estimated based on radio/X-ray luminosities. ADAF sources could however contain jets/outflows which can contribute to radio emission. Depending on the level of radio activity and existence of extended radio emission features, galactic sources could be classified (Yi & Boughn 1998ab). Such a classification can be quantified in a manner similar to that adopted for Galactic X-ray sources. In fact, there exist interesting spectral similarities between Galactic binary X-ray sources and galactic nuclei.

We define X-ray bright galactic nuclei (XBGN) as galactic nuclei with X-ray luminosities in the range $10^{40} \lesssim L_x \lesssim 10^{42}$ erg/s which is sub-luminous compared with the more powerful active galactic nuclei (AGN) which generally have $L_x \gtrsim 10^{43}$ erg/s. Most of XBGN are expected to overlap with emission line galaxies with $L_x \sim 10^{39} - 10^{42}$ erg/s. However, some of the low luminosity Seyferts with $L_x \gtrsim 10^{42}$ erg/s cannot be ruled out. Based on our discussions of ADAFs, for $L_x \sim 10^{41}$ erg/s,

$$L_R \sim 4 \times 10^{36} (M_{BH}/3 \times 10^7 M_\odot) \text{ erg/s}$$

(32)

at 20 GHz with the characteristic inverted radio spectrum $I_\nu \propto \nu^{2/5}$. These sources at distances $\sim 10(M/3 \times 10^7 M_\odot)^{1/2} \text{Mpc}$ should be detected as $\sim 1 mJy$ point-like radio sources (Yi & Boughn 1998a). If X-ray and radio are indeed from ADAFs, the black hole masses can be estimated (Yi & Boughn 1998b).

Yi and Boughn (1998ab) proposed the source classification based on the known black hole masses and the ADAF radio/X-ray luminosity relation. Adopting Sgr A*, NGC 4258, NGC 1068, NGC 1316, NGC 4261, and NGC 4594 as
fiducial sources, a statistically incomplete sample of XBGN are classified into radio-loud XBGN and radio-quiet XBGN. The former show that the observed radio luminosity $L_{R,\text{obs}} \sim L_{R,\text{jet}} \gg L_{R,\text{adv}}$ where $L_{R,\text{jet}}$ and $L_{R,\text{adv}}$ are the expected radio jet luminosity and ADAF radio luminosity, respectively. These sources are expected to show extended radio emission, unlikely to have strongly inverted radio spectra, and may have compact ADAF radio emission from compact cores separate from the extended emission components. The latter show that $L_{R,\text{obs}} \sim L_{R,\text{adv}}$ and that the dominating emission components are compact cores with inverted spectra.

Kellermann et al. (1998) and Falcke (1998) show that jet-like radio emission features are common among AGN and emission line galaxies. Surprisingly, even in radio-quiet sources, elongated radio emission features are sometimes seen on small scales, which could imply that some type of jet/outflow activities are very common in galactic nuclei regardless of their large scale radio activities. In order to resolve this issue, high resolution radio measurements for nearby (<10Mpc) sources are crucial. Hard X-ray emission and inverted spectrum, compact radio emission are very likely to be found closely correlated.

5.5. QSO Evolution

Yi (1996) has suggested that transition of accretion flows from thin disks to ADAFs at $\dot{m}_{\text{crit}} \approx 0.3\alpha^2$ could account for the observed sudden decline in the number of bright QSOs at redshift $z \sim 2$ and downward (see also Fabian & Rees 1995). Once ADAFs set in, the luminosity evolves according to

$$L = L_{\text{ADAF}} \approx 30\dot{m}^2 L_{\text{Edd}}$$

where $x \sim 2$ and $\dot{m} = \dot{M}/\dot{M}_{\text{Edd}} \propto \dot{M}/M$. The last expression implies that even when the mass accretion rate $\dot{M}$ is kept constant, $\dot{m}$ decreases merely due to the growth of black hole mass.

For instance, in a flat universe with no cosmological constant, for the initial black hole mass $M = M_i$ at $z = z_i$ and $\dot{M} = \text{constant}$, when $\dot{m} < \dot{m}_{\text{crit}}$ and $M/H_0 \gg M_i$,

$$L(z) \propto (1 + z)^{3(x-1)/2} \left[ (1 + z)^{-3/2} - (1 + z_i)^{-3/2} \right]$$

or

$$L(z) \propto (1 + z)^{K(z)}$$

with

$$K(z) = \frac{3(x-1)}{2} \left[ 1 + \frac{1}{[(1+z)/(1+z_i)]^{3/2} - 1} \right]$$

which shows that the luminosity declines with a power-law similar to that seen in observations (Yi 1996 and references therein). The epoch at which ADAFs set in for $\dot{m} = \dot{m}_i \sim 1$ (i.e. $\dot{m} = \dot{m}_{\text{crit}}$ first occurs)

$$1 + z_c = \left[ \left( \frac{t_{\text{Edd}}}{t_o} \right) \left( \frac{1}{\dot{m}_{\text{crit}}} - \frac{1}{\delta} \right) + (1 + z_i)^{-3/2} \right]^{-2/3}$$

where $\delta = \dot{M}t_{\text{Edd}}/M_i$ and $t_{\text{Edd}} = (\dot{M}_{\text{Edd}}/M)^{-1} = 4.5 \times 10^7 (\eta_{\text{eff}}/0.1) \text{yr}$. The observed sudden decline of QSOs at $z \sim 2$ is naturally accounted for.
5.6. Outflows
ADAFs are prone to outflows or jets (Narayan & Yi 1994,1995a) although a self-consistent inflow/outflow solution has not been found yet (cf. Xu & Chen 1997). It remains to be seen if a self-consistent inflow/outflow solution can account for compact and extended jet-like emission components in XBGN.

6. Some Unresolved Issues
ADAFs may exist in sources spanning many decades of masses of compact accreting sources. Some of the old outstanding issues, which are mostly concerning with low-luminosities and hard X-ray emission, are plausibly resolved by various versions of ADAF models. There exist however a number of unsolved problems in the ADAF framework.

(i) The physics of accretion flow transition, temporal and spatial, remains unclear. The spatial transition (i.e. from outer thin disk to inner ADAF) is better understood than the temporal transition to a certain extent. However, there has not been an adequate explanation for $r_t$. It remains unsolved why the disk flow makes a transition to ADAF with little, if any, change of accretion rate.
(ii) Even luminous systems show energetic X-ray emission which is absent in the thin disk models. It is often assumed that these sources have X-ray emitting coronae for which little physics is known. It is crucial to link the ADAF models to corona models with proper physical understanding. (iii) ADAFs emission could be highly variable. The issue of steady vs. non-steady ADAFs is directly related to the observed variabilities in ADAF candidate sources which show occasional non-detections.
(iv) A number of plasma astrophysical issues remain to be solved. Is the two-temperature flow physically allowed? Are particles rapidly thermalized? What is the correct strength of magnetic fields responsible for synchrotron emission? (v) Observationally, faint X-ray sources, which could be seen by high resolution, high sensitivity experiments such as AXAF, should be studied in great details. Ultimately, ADAF-related issues and the very question on the future of ADAFs’ are likely to be answered by observations.

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