Radiatively Driven Mass Accretion onto Galactic Nuclei by Circumnuclear Starbursts

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ABSTRACT

We examine the physical processes regarding the radiatively driven mass accretion onto galactic nuclei due to the intensive radiation from circumnuclear starbursts. The radiation from a starburst not only contracts an inner gas disk by the radiation flux force, but also extracts angular momenta due to the relativistic radiation drag, thereby inducing an avalanche of the surface layer of disk. To analyze such a mechanism, the radiation-hydrodynamical equations are solved with including the effects of radiation drag force as well as radiation flux force. As a result, it is found that the mass accretion rate due to the radiative avalanche is given by \( \dot{M}(r) = \eta \frac{L_* c^2}{c^2} \left( \frac{r}{R} \right)^2 \left( \frac{\Delta R}{R} \right) \left( 1 - e^{-\tau} \right) \) at radius \( r \), where the efficiency \( \eta \) ranges from 0.2 up to 1, \( L_* \) and \( R \) are respectively the bolometric luminosity and the radius of a starburst ring, \( \Delta R \) is the extension of the emission regions, and \( \tau \) is the face-on optical depth of a disk. In an optically thick regime, the rate depends upon neither the optical depth nor the surface mass density distribution of the disk. The present radiatively driven mass accretion may provide a physical mechanism which enables mass accretion from 100 pc scales down to \( \sim \) pc, and it can eventually link to an advection-dominated viscous accretion onto a massive black hole. The radiation-hydrodynamical and self-gravitational instabilities of the disk are briefly discussed. In particular, the radiative acceleration possibly builds up a dusty wall, which shades the nucleus in edge-on views. It provides another version for the formation of an obscuring torus.

Key words: accretion, accretion disks — black hole physics — galaxies: active — galaxies: nuclei — galaxies: Seyfert — galaxies: starburst — hydrodynamics — quasars: general — radiative transfer
1 Introduction

In a standard picture of active galactic nuclei (AGNs), a viscous accretion disk of $\lesssim \text{pc}$ surrounding a putative supermassive black hole is thought to be the central engine for the enormous energy output. The mass supply onto such tiny regions is an issue of growing interest. The accretion driven by a non-axisymmetric wave or a bar is a possible mechanism (Shore & White 1982; Shlosman, Frank, & Begelman 1989), but it is found so far that the bar-driven accretion is just effective in the galactic nuclear regions beyond $\sim 100 \text{ pc}$ (Wada & Habe 1995). Hence, another mechanism that enables the mass accretion from 100 pc down to pc scales is required. Recently, Umemura, Fukue, & Mineshige (1997) have proposed a possible physical mechanism in the relevant scales, that is, the mass accretion driven by radiation from circumnuclear starbursts.

In recent years, there have been accumulated observational evidences for starbursts in circumnuclear regions of AGNs. Among them, it is remarkable that molecular gas of at least $10^{11} \text{M}_\odot$ and dust of $\sim 10^8 \text{M}_\odot$ are found towards a QSO (BR1202-0725) at redshift $z = 4.69$ (Ohta et al. 1996; Omont et al. 1996). This, taking into account the Universe age of $\sim 1\text{Gyr}$ at $z = 4.69$ and the inferred metallicity, suggests that the QSO inhabits a massive host galaxy undergoing starbursts. Also, the recent HST images of nearby luminous QSOs have shown that luminous QSO phenomena occur preferentially in luminous host galaxies, often being ellipticals (McLeod & Rieke 1995; Bahcall et al. 1997; Hooper, Impey, & Foltz 1997). In addition, AGN events appear to be frequently accompanied by starbursts (Soifer et al. 1986; Scoville et al. 1986; Heckman 1991; Rieke 1992; Scoville 1992; Fillipenko 1992). The IRAS survey has revealed that at bolometric luminosities greater than $10^{12} \text{L}_\odot$, the space density of starburst galaxies are just comparable to that of QSOs in the same luminosity range (Soifer et al. 1986). Millimeter observations of the IR galaxies indicate that all are exceedingly rich in molecular gas with up to $2 \times 10^{10} \text{M}_\odot$ in their central regions (Scoville et al. 1986). The correlations between X-ray and CO luminosities as well as between CO and far-infrared luminosities in Seyfert galaxies also suggest that more powerful AGNs inhabit more active star-forming galaxies (Yamada 1994). Such increasing evidences suggest a possible physical connection between AGN events and starburst activities (Norman & Scoville 1988; Rowan-Robinson & Crawford 1989; Terlevich 1992; Scoville 1992; Terlevich 1992; Perry 1992; Taniguchi 1992; Rowan-Robinson 1995).

As for the structures of starburst regions, recent high-resolution observations of Seyfert nuclei, including those from HST, have revealed circumnuclear starburst regions with radial extension of a few tens of pc up to kpc, which often exhibit ring-like features (Wilson et al. 1991; Forbes et al. 1994; Mauder et al. 1994; Buta, Purcell, & Crocker 1995; Barth et al. 1995; Maoz et al. 1996; Storchi-Bergman, Wilson & Baldwin 1996; Leitherer et al. 1996). The rings frequently contain a number of extremely compact star clusters of $< 5 \text{ pc}$, which are composed of hundreds of O-type stars. Recently, an unexpectedly large ($\sim 100 \text{ pc}$) disk of gas and dust is discovered by the HST around a bright nucleus in an active galaxy NGC 4261 (Jaffe et al 1993; Ferrarese, Ford, & Jaffe 1996). This disk is surrounded by a fat torus composed of early-type stars.
Umemura, Fukue, & Mineshige (1997, hereafter referred to as Paper I) have considered the radiative effects on nuclear disks due to circumnuclear starbursts, and proposed the possibility of a radiative avalanche, that is, the mass accretion driven by the relativistic radiation drag. In paper I, the effects of radiation drag are analyzed solely by solving the azimuthal equation of motion. In this paper, we explore extensively the radiative avalanche with including radial motion driven by radiation flux force. Also, the radiation-hydrodynamical instability of the nuclear region is analyzed with including vertical motion as well, so that the radiative building-up of a dusty obscuring wall is predicted. Some related topics such as the supernova effects, the generation of seed magnetic fields, and a possible link to an advection-dominated viscous flow are discussed. In §2, the basic equations which govern the radiation-hydrodynamical disk evolution are provided. In §3, the timescales for a radiative avalanche, the disk evolution, and the resultant mass accretion rate are presented. In §4, the radiation-hydrodynamical/gravitational instabilities are analyzed. In §5, we discuss the effects of stellar evolution, the generation of magnetic fields, and a link between a radiative avalanche and an advection-dominated accretion onto a black hole. §5 is devoted to the conclusions.

2 Basic Equations

2.1 Radiation Fields by Starbursts

Here, we suppose a geometrically thin ring of starburst regions for simplicity, so that the assumption allows us to explore the radiative effects analytically. The effects of the extension of the starburst regions are briefly discussed. Under the assumption, the radiation fields by the starburst ring with the radius $R$ and the bolometric luminosity $L_*$ are given as follows (see also Paper I): The radiation energy density at $r(< R)$ is given by

$$E = \frac{L_*}{4\pi c(R^2 - r^2)}.$$  \hspace{1cm} (1)

The radiation flux in radial directions is given by

$$F^r = \frac{L_*}{8\pi R^2} \frac{2x(x^2 + 1)G_1 - 3xG_2}{(x^2 + 1)^{5/2}},$$  \hspace{1cm} (2)

where $x \equiv r/R$, and $G_1$ are $G_2$ are respectively $G_1 = G[3/4, 5/4, 1, 4x^2/(x^2 + 1)^2]$, and $G_2 = G[5/4, 7/4, 2, 4x^2/(x^2 + 1)^2]$ with $G$ being the Gaussian hypergeometric function. The solution is well approximated by

$$F^r \simeq -\frac{L_*}{4\pi R^2} \left[ \frac{r}{2R} + \frac{r^3}{\pi R^2(R - r)} \right].$$  \hspace{1cm} (3)

[This is slightly different from equation (2) in Paper I, but this is a more accurate approximation.] The rotation of the ring generates the toroidal flux at $r$ to be

$$F^\varphi = \frac{3EVr}{2R}.$$  \hspace{1cm} (4)
where $V$ is the rotation velocity of the ring. The radiation stress tensors are given as

$$(P^{rr}, P^\varphi\varphi) = \left(\frac{E}{2}, \frac{E}{2}\right). \quad (5)$$

### 2.2 Radiation Hydrodynamical Equations

As a relativistic result of radiative absorption and subsequent re-emission, radiation fields exert drag force on moving matter in resistance to its velocity. This radiation drag extracts angular momentum from a rotating disk, thereby allowing the gas to accrete onto the center. This is known to be the Poynting-Robertson effect in the solar system (Poynting 1903; Robertson 1937). The radiation drag usually makes effects of $O(v/c)$, but it could provide a key mechanism for momentum/angular momentum transfer in environments of intensive radiation fields (Loeb, 1993; Umemura, Loeb & Turner 1993; Fukue & Umemura 1994; Umemura & Fukue 1994; Tsuribe, Fukue, & Umemura 1994; Tsuribe, Umemura & Fukue 1995; Fukue & Umemura 1995; Umemura, Fukue & Mineshige 1997; Tsuribe & Umemura 1997).

Specific radiation forces which are exerted on moving fluid elements with velocity $v$ are given as

$$f_r = \frac{\chi c}{c} (F^r - v_r E - v_r P^{rr} - v_\varphi P^\varphi r)$$

and

$$f_\varphi = \frac{\chi c}{c} (F^\varphi - v_\varphi E - v_\varphi P^{\varphi\varphi} - v_r P^r \varphi), \quad (6)$$

(Mihalas & Mihalas 1984) respectively in radial and azimuthal directions. Here $\chi$ is the mass extinction coefficient, which is given by $\chi = (\kappa + n_e \sigma_T)/\rho$, where $\kappa$ is the absorption coefficient, $n_e$ is the electron number density, $\sigma_T$ is the Thomson scattering cross section, and $\rho$ is the mass density. Coupled with the radiative quantities presented in the previous section, the last term of each component of specific forces turns out to be $O(vV/c^2)$, and thus is neglected in the present analyses. Using equations \((4)\) and \((5)\), we have the radiation hydrodynamical equations to $O(v/c)$ as

$$\frac{dv_r}{dt} = \frac{v_\varphi^2}{r} - \frac{1}{\rho} \frac{dp}{dr} - \frac{\Phi}{c} + \frac{\chi}{c} \left(F^r - \frac{3}{2} Ev_r\right), \quad (8)$$

$$\frac{1}{r} \frac{d(rv_\varphi)}{dt} = \frac{3\chi E}{2c} \left(\frac{r}{R} V - v_\varphi\right), \quad (9)$$

where $p$ is the gas pressure, and $\Phi$ is the external force potential (e.g., gravity). The azimuthal equation of motion \((9)\) implies that the radiation flux force (the first term on the right-hand side) makes fluid elements tend to co-rotate with the stellar ring, whereas the radiation drag (the second term on the right-hand side) works to extract the angular momenta from the elements. It is noted that both terms are $O(v/c)$ in this equation, in contrast to the radial flux force of $O(1)$ in equation \((8)\).
3 Radiative Avalanche

3.1 Timescales

As shown in Paper I, the radiation drag time scale is given by

\[ t_\gamma \equiv \frac{8\pi c^2 R^2}{3\chi L_*} \]  \hspace{1cm} (10)

First, we consider an ionized disk. The total number per unit time of ionizing photons from starburst regions is roughly estimated as

\[ \dot{N}_{UV} \sim f_{UV} L_/h\nu_L \sim 5.4 \times 10^{-5} s^{-1} (L_* / 3 \times 10^{12} L_\odot), \]

with \( \nu_L \) being the Lyman limit frequency, and \( f_{UV} \) is a factor which represents the contribution to UV light by starbursts (assumed to be \( f_{UV} \sim 0.1 \) here), while the total recombination rate to excited levels of hydrogen is

\[ \dot{N}_{\text{rec}} = n_e \alpha^{(2)} M_g / m_p \rightsim 1 \times 10^{56} s^{-1} (M_g / 10^8 M_\odot) (R_g / 100 \text{pc})^{-3}, \]

where \( M_g \) and \( R_g \) are the mass and the radius of the gas disk and \( \alpha^{(2)} \) is the recombination rate coefficient to all excited levels of hydrogen. The optical depth of ionized gas is roughly estimated as

\[ \tau \sim M_g \sigma_T / \pi R_g^2 m_p, \]  \hspace{1cm} (11)

This is fairly long compared to the duration of AGN phase or starburst phase inferred from observations (\( \sim 10^8 \) yr).

In practice, a large amount of dust is detected in luminous IR galaxies (e.g. Scoville et al. 1991). From theoretical points of view, the timescale of photodesorption of dust by UV is longer than the Hubble time for refractory grains like graphites or silicates, and also the thermal sputtering rate is completely negligible for such grains at the temperature of \( \sim 10^4 \text{K} \) (Burke & Silk 1974; Draine & Salpeter 1979). If a disk contains dust grains, then the accretion timescale could be much shorter because of its large opacity. The optical depth is estimated to be

\[ \tau \sim 5 \times 10^2 (M_g / 10^8 M_\odot) (R_g / 100 \text{pc})^{-2} (f_{dg} / 10^{-2}), \]

where \( f_{dg} \) is the dust-to-gas mass ratio (Mathis, Rumpl, & Nordsiek 1977). If the relaxation between gas and dust is sufficiently rapid, then the dust accretes with accompanying gas. The relaxation time between gas and dust is estimated by

\[ t_{\text{relax}} = \frac{3 m_d (2\pi)^{1/2} (kT)^{3/2}}{8\pi c^4 m_p^{1/2} n_p c^2 Z_d^2 \ln \Lambda}, \]  \hspace{1cm} (12)

where \( m_d \) is the mass of dust grain, \( n_p \) is the number density of protons, \( Z_d \) is the dust charge in units of proton charge, and \( \ln \Lambda \) is the Coulomb logarithm (Spitzer 1962). Tentatively
assuming the thickness of gas disk to be one tenth of the radius, we evaluate \( n_p = n_e \simeq 1.3 \times 10^6 \text{cm}^{-3}(M_g/10^9M_\odot)(R_g/100\text{pc})^{-3} \). Then, for spherical dust grains,

\[
    t_{\text{relax}} = 3.1 \times 10^{-2} \text{yr} \left( \frac{n_p}{10^6\text{cm}^{-3}} \right)^{-1} \left( \frac{T}{10^4\text{K}} \right)^{-1/2} \left( \frac{a_d}{0.1\mu\text{m}} \right) \left( \frac{\rho_s}{\text{g cm}^{-3}} \right) \left( \frac{Z_d}{40} \right)^{-2},
\]

where \( a_d \) is the grain radius and \( \rho_s \) is the density of solid material within the grain (e.g., Spitzer 1978). Hence, dust is tightly coupled with gas in a typical rotation timescale of \( \sim 10^5 \) yr, and therefore the disk can be regarded to be composed of fluid with mass extinction as

\[
    \chi = \frac{n_e \sigma_T + n_d \sigma_d}{\rho_g + \rho_d}, \quad (14)
\]

where \( n_d, \sigma_d, \) and \( \rho_d \) are the number density, cross-section, and mass density of dust, respectively, and \( \rho_g \) is the mass density of gas. If taking the dust-to-gas mass ratio of \( \sim 10^{-2} \) that is inferred in the Solar neighborhood, then the opacity ratio is

\[
    \frac{n_d \sigma_d}{n_e \sigma_T} = 1.9 \times 10^3 \left( \frac{a_d}{0.1\mu\text{m}} \right)^{-1} \left( \frac{\rho_s}{\text{g cm}^{-3}} \right)^{-1} \left( \frac{f_{d g}}{10^{-2}} \right). \quad (15)
\]

Thus, we find \( n_d \sigma_d \gg n_e \sigma_T \) and \( f_{d g} \ll 1 \) in probable situations. Then, we can evaluate \( \chi \simeq n_d \sigma_d/\rho_g \), so that

\[
    t_{\gamma} \simeq 2.4 \times 10^6 \text{yr} \left( \frac{L_*}{3 \times 10^{12}L_\odot} \right)^{-1} \left( \frac{R}{100\text{pc}} \right)^2 \left( \frac{f_{d g}}{10^{-2}} \right)^{-1} \left( \frac{a_d}{0.1\mu\text{m}} \right) \left( \frac{\rho_s}{\text{g cm}^{-3}} \right). \quad (15)
\]

Hence, for intensive starbursts, the drag timescale could be shorter than the duration of AGN phase or starburst phase.

### 3.2 Disk Evolution

Suppose that an initially quiescent gas disk is abruptly ignited by a starburst at the radius \( R \) with luminosity \( L_* \). Since the radiation drag efficiency decreases exponentially with the optical depth (Tsuribe & Umemura 1997), just an optically thin surface layer avalanches toward a nucleus. In equation (8), the radial radiation flux force, \( \chi F^r/c \), is obviously greater by \( O(c/v) \) than the drag force. So, the radiation drag makes a negligible contribution in the radial equation of motion. On the other hand, in azimuthal directions (see equation 3), both the radiation flux force and the radiation drag are \( O(v/c) \), and therefore they work simultaneously. Thus, the disk evolution is expected to proceed in two steps: First, the optically thin surface layer (or the whole of an optically thin disk) is contracted by the radial radiation flux force and is quickly settled in quasi-rotation balance. Then, the layer sheds angular momentum due to the radiation drag, and consequently undergoes an avalanche.

With the radial equation of motion (8), disregarding pressure gradient force as well as the drag force, the condition of rotation balance is given by

\[
    \frac{v_r^2}{r} = \frac{d\Phi}{dr} - \frac{\chi}{c} F^r. \quad (16)
\]
Here, we assume a gravitational potential such that
\[ r \frac{d\Phi}{dr} \equiv V^2 \left( \frac{r}{R} \right)^{2n}. \]  
(17)
The rotation velocity of the stellar ring is given by
\[ V = \left( \frac{GM_{\text{dyn}}}{R} \right)^{1/2}, \]
where \( M_{\text{dyn}} \) is the dynamical mass contained within \( R \). If \( F^r \) is negligibly small, then \( n = -1/2 \) simply corresponds to Keplerian rotation, \( n = 0 \) to flat rotation, and \( n = 1 \) to rigid rotation.

Equation (16) is rewritten with (17) into a dimensionless form as
\[ \tilde{v}_\phi^2 = x^{2n} + \Gamma x^2 \left[ \frac{1}{2} + \frac{x^2}{\pi(1-x)} \right], \]  
(18)
with using \( x \equiv r/R \), \( \tilde{v}_\phi \equiv v_\phi/V \), and \( \Gamma \equiv L_*/L_E \), where \( L_E \) is the Eddington luminosity corresponding to \( M_{\text{dyn}} \),
\[ L_E = \frac{4\pi c G M_{\text{dyn}}}{\chi}. \]  
(19)
For a dusty disk, the Eddington luminosity is numerically given as
\[ L_E = 1.68 \times 10^{11} L_\odot \left( \frac{M_{\text{dyn}}}{10^{10} M_\odot} \right) \left( \frac{f_{\text{dg}}}{10^{-2}} \right)^{-1} \left( \frac{a_d}{0.1 \mu\text{m}} \right) \left( \frac{\rho_s}{\text{g cm}^{-3}} \right). \]  
(20)
The azimuthal equation of motion (9) can be also rewritten in a dimensionless form as
\[ \frac{d(x\tilde{v}_\phi)}{dt} = \frac{x(x - \tilde{v}_\phi)}{(1 - x^2)}, \]  
(21)
where \( \tilde{t} \equiv t/t_\gamma \).

To begin with, we see the first stage, i.e., the radial contraction phase. As shown in equation (18), the radial compression by the flux force is expected to be strongest on the outer edge of disk. Because the radial flux does not extract angular momentum, the angular momentum conservation determines the radius of rotation equilibrium. Since \( \tilde{v}_\phi = 1 \) at the outer edge of the disk before a starburst, the equilibrium radius after a starburst is determined by the condition of \( \tilde{v}_\phi = 1/x \). Combined with equation (18), the resultant equilibrium radii \( r_{eq} \) for the material that is originally located at the outer edge are shown versus \( \Gamma \) for three types of rotation potentials in Figure 1. The \( r_{eq} \) is a slowly decreasing function with \( \Gamma \) because of angular momentum barrier. The asymptotic behavior for \( \Gamma \ll 1 \) is \( r_{eq}/R = 1 - \frac{\Gamma}{2\pi(n+1)} \)^{1/2}, while \( r_{eq}/R = (2/\Gamma)^{1/4} \) for \( \Gamma \gg 1 \). Equation (18) shows that the radial radiation flux enhances the rotation velocity, so that a shear flow may be excited on the boundary between the optically thin surface layer and the optically thick deep stratum. The fractional difference of rotation velocity, which is defined by \( \Delta v_\phi/v_\phi \equiv (\tilde{v}_\phi - x^n)/x^n \), is also shown at \( r_{eq} \) in Figure 1. \( \Delta v_\phi/v_\phi \) is found to be weakly dependent upon \( \Gamma \), and we have asymptotic solutions as
\[
\frac{\Delta v_\phi}{v_\phi} = \begin{cases} 
\left[ \frac{(n+1)\Gamma}{2\pi} \right]^{1/2} & \text{for } \Gamma \ll 1 \\
\Gamma^{\frac{n+1}{2}} & \text{for } \Gamma \gg 1.
\end{cases}
\]  
(22)
If $\Gamma = O(1)$ (which is indeed the case for active galaxies), $\Delta v_\phi/v_\phi = 0.2 - 0.6$ in the range of $n = -1/2$ to $1$. This shear flow may generate turbulence due to the Kelvin-Helmholtz instabilities, or the viscosity could thermalize the shear flow energy.

Next, we consider the avalanche phase. If the radiation is very weak ($\Gamma \ll 1$), equation (21) coupled with (18) is reduced to a differential equation as

$$x' = \frac{dx}{dt} = \frac{(1 - x^{n-1})}{(n+1)x^n(1 - x^2)}. \quad (23)$$

This equation can be analytically solved for some values of $n$. For instance, the Keplerian rotation potential ($n = -1/2$) leads to the solution

$$\tilde{t} = -\sqrt{x} + \frac{1}{\sqrt{3}} \tan^{-1}\left(\frac{1 + 2\sqrt{x}}{\sqrt{3}}\right) - \frac{1}{2} \log\left(\frac{x}{1 + x + \sqrt{x}}\right) - 0.154. \quad (24)$$

If we take a Mestel disk ($n = 0$), then the solution is

$$\tilde{t} = 1 - x - \log x. \quad (25)$$

The rigid-rotation potential results in $x = \text{const}$. In other words, no accretion is driven in a rigid-rotation potential under the circumstances of very weak radial radiation flux (see also Paper I). If $\Gamma = O(1)$, however, the effects of radial flux force is important not only in radial contraction but also in the angular momentum transfer. When the radial flux force is included, the accretion is governed by

$$x' = \frac{\bar{v}_\phi(x - \bar{v}_\phi)}{(1 - x^2)h(x)}. \quad (26)$$

from equations (18) and (21), where

$$h(x) = (n+1)x^{2n-1} + \Gamma x \left[1 + \frac{x^2(6 - 5x)}{2\pi(1 - x)^2}\right]. \quad (27)$$

As shown below soon, this equation allows the mass accretion even in a rigid-rotation potential.

### 3.3 Mass Accretion Rate

If we suppose a geometrically thin, optically thick disk, then the upper surface is irradiated by radiation from the upper half of starburst regions and the lower surface is irradiated by the lower half. Taking this into account, we can estimate the mass accretion rate as

$$\dot{M} \equiv -\int 2\pi r v_r \rho dz 
\simeq -(1 + e^{-\tau/2}) \int_0^{\tau/2} \frac{2\pi r v_r}{\chi} e^{-\tau} d\tau_1, \quad (28)$$
where \( z \) is the vertical coordinate, \( d\tau_1 \equiv \chi \rho dz \), and \( \tau \) is the total face-on optical depth. This can be rewritten in a dimensionless form using the same notations as the previous subsection,

\[
\dot{m}_\gamma = \frac{3}{4} \left( 1 + e^{-\tau/2} \right) \int_0^{\tau/2} xx' e^{-\tau_1} d\tau_1 = -\frac{3}{4} xx'(1 - e^{-\tau}),
\]

(29)

where \( \dot{m}_\gamma \equiv \dot{M}/(L_*/c^2) \) and \( x' \) is given by (26). In an optically thick limit,

\[
\dot{m}_\gamma = -\frac{3}{4} xx'.
\]

(30)

Hence, the mass accretion rate is independent of the extinction coefficient as well as the density distribution in an optically thick disk. In comparison, in an optically thin regime, \( \dot{m}_\gamma = -3xx'\tau/4 \). This implies that the mass accretion rate is smaller in an optically thin disk than in an optically thick disk. This is because in optically thin media a number of photons penetrate the gas disk without collision with matter, and hence not all the photons from the starburst regions are available for radiation drag.

The fact that the radial flux raises the rotational velocity brings two effects on the mass accretion rate which mutually compete: One is a positive effect that the radiation drag force becomes more efficient due to enhanced \( v_\varphi \), and the other is negative effect that the enhanced \( v_\varphi \) leads to the increase of specific angular momentum at fixed \( r \), so that the accretion is decelerated. In Figure 2, for a wide variety of \( \Gamma = L_*/L_E \), the resultant mass accretion rate is shown as a function of radii in an optically thick regime. In this figure, a thin solid curve shows the estimation in the case of the negligible effects of the radial flux, i.e.,

\[
\dot{m}_\gamma = \frac{3}{8} \frac{1 - n}{1 + n} x^2
\]

(31)

(Paper I). The thick curves are the present results, where filled circles represent the equilibrium radii of outer-edge shown in Figure 1. As shown in Figure 2a, for a Keplerian disk, the negative effect overwhelms the positive one around the outer-edge of the disk. On the other hand, the positive effect is predominant for a Mestel-like disk (\( n = 0 \)) as shown in Figure 2b. In the vicinity of the ring (\( x \to 1 \)), we have an asymptotic solution as \( \dot{m}_\gamma = 3/4 \), being independent of \( \Gamma \) as well as \( n \). In the opposite limit, i.e., \( x \to 0 \), the solutions are not dependent on \( \Gamma \) for \(-1 < n < 1 \), and then the mass accretion rate is basically determined by (31).

It is worth stressing that the radial flux can allow the accretion even if the disk is governed by a rigid-rotation potential (\( n = 1 \)). As shown in Figure 2c, the induced mass accretion is not sensitive to \( \Gamma \), and the maximal rate is \( \dot{M} \sim 0.2L_*/c^2 \) in the range of \( \Gamma = 10^{-3} \) to \( 10^2 \). At small radii, there is an asymptotic solution as

\[
\dot{m}_\gamma = \frac{3}{8} x^2 \left[ 1 - \left( 1 + \frac{\Gamma}{2} \right)^{-1/2} \right].
\]

(32)

This approaches a limiting value \( 3x^2/8 \) for \( \Gamma \to \infty \), which coincides with (31) for \( n = 0 \).
In the long run, the mass accretion rate has turned out to be almost insensitive to the rotation law and given by

\[
\dot{M} = \eta \left( \frac{L_*}{c^2} \right) x^2 (1 - e^{-\tau})
\approx 0.2 \dot{M}_\odot \text{yr}^{-1} \eta \left( \frac{L_*}{3 \times 10^{12} L_\odot} \right) x^2 (1 - e^{-\tau}),
\]

(33)

where \( \eta \approx 0.2 - 1 \). Such radial dependence as \( x^2 \) implies that an avalanche occurs in a homologous manner. The radial dependence slightly deviates from \( x^2 \) near the outer edge of the disk. Especially, it is faster than \( x^2 \) in a rigid-rotation potential, so that some amount of mass would accumulate around \( r_{eq} \).

Although we have assumed a ring-like configuration for starburst regions, the regions may be extended over a wide range of radii and heights. However, equation (3) shows that the radiation from the inner edge of the star-forming regions would be most contributive as long as the surface brightness falls faster than \( \Sigma_L(R) \propto (1 - r^2/R^2) \), which is basically constant for \( R \gg r \). In other words, the regions of peak surface brightness would be always most responsible for the radiative avalanche. Also, in an optically thick regime, the effects of viewing-angles from starburst regions to a disk might reduce the mass accretion rate, depending upon the configuration of starburst regions (Taniguchi 1997). If we suppose a torus of emission regions with the half thickness of \( \Delta R \), then the rate should be multiplied roughly by \( \Delta R/R \) due to a geometrical dilution effect (Ohsuga et al. 1997).

4 Stability of Disk

4.1 Radiation-Hydrodynamical Instability

If we suppose a geometrically thin, optically thick disk, the upper surface is irradiated downward, while the lower surface is irradiated upward. Hence, the disk is rather confined by radiation force. But, the outer edge of an optically thick disk (of course, an optically thin disk as well) could be vertically unstable, especially if the starburst luminosity is super-Eddington. As far as a Kepler disk or a Mestel disk is concerned, the gravity is always dominant at small radii [see (3)]. Thus, the inner disk remains stable, whereas the outer edge could be blown away by radiative acceleration. The radius inside which the disk is stable is given by the root of an equation as

\[
x^{2n-3}(1 - x) = \Gamma \left[ \frac{1}{2} + \frac{x^2}{\pi(1-x)} \right].
\]

(34)

The solutions are shown in Figure 3 for \( n = -1/2, 0, \) and 1. We find that the radius is sensitively dependent upon \( \Gamma \) if \( \Gamma \gtrsim 1 \). To see the fate of the blown-out gas in three-dimensional space, we compare the radiation flux force with the gravitational force in Figure 4, illustratively for \( n = -1/2 \). On the hatched sides for each value of \( \Gamma \), the radiative force overwhelms the gravity in the vertical directions. The dotted lines are corresponding to
the unstable branches below which gas is accelerated downward by the gravity, while above which gas blows upward due to radiative force. The solid lines are stable branches. When \( \Gamma > 2 \), no stable branch exists and therefore the blown-out gas may result in a superwind (a “blizzard” phase). [If we assume the initial mass function which is deficient in low-mass stars as observed in starburst regions (see also §4.1), \( \Gamma > 2 \) is corresponding to early stages of \( < 5 \times 10^7 \text{yr} \) after the coeval star-formation.] If the gravity by the bulge stars significantly decelerates the wind, the spurted-out gas may fall onto the inner disk with losing angular momentum by a fraction of \( \sim R/V t_\gamma \) (cf. Meyer & Meyer-Hofmeister 1994). When \( \Gamma < 2 \), the stable branch composes a wall surrounding the starburst ring. In particular, the wall becomes of torus shape if \( \Gamma < 1 \). Taking account of the fact that the radiation force becomes ineffective for optically-thick gas, the optical depth of the formed wall should be of \( O(1) \). The radiation emitted originally in UV is absorbed by this wall and re-emitted possibly in IR wavelength. However, the radiative avalanche is almost independent of the spectrum of the radiation, but just depends upon the bolometric luminosity. Hence, the emission from the wall could induce an indirect radiative avalanche. Also, such a radiatively built-up obscuring wall may diminish the geometrical dilution, although the direct avalanche may be strongly subject to the geometrical dilution if the thickness of the starburst regions is as small as \( 10 \text{pc} \) (Taniguchi 1997).

The radiatively built-up obscuring wall may be a different version of the formation of an obscuring torus from those proposed so far (e.g. Krolik & Begelman 1988; Yi, Field, & Blackman 1994). In the present mechanism, the angular extension of the obscuring wall is a decreasing function of time according as the starburst luminosity fades out. Thus, the classification of AGN types by anisotropic viewing could be comprehended in the contexts of the evolution of circumnuclear starbursts. The further detailed confrontation with observations will be made elsewhere.

Finally, even an inner gravity-dominant disk could be unstable to warping through a similar mechanism proposed by Pringle (1996, see also Maloney et al. 1996), although the radiative acceleration is reduced by a factor of \( \sim 1/\tau \). If a radiative backreaction torque can drive a warp, the illumination to the warped disk is augmented, thereby increasing the accretion.

4.2 Gravitational Instability

We consider three components of gravity sources, that is, a central massive black hole, a massive disk surrounding the black hole, and the galactic bulge. The surface density distribution of the disk is assumed to be Mestel-like, and the bulge is to be uniform. To see the stability, we evaluate the Toomre’s \( Q \) to find

\[
Q = \frac{2c_s}{V f_g} \left( \frac{f_{\text{BH}}}{x} + 2f_g + 4f_{\text{bulge}}x^2 \right)^{1/2},
\]

where \( V = (GM_{\text{dyn}}/R)^{1/2} \) is the virial velocity for the total dynamical mass \( M_{\text{dyn}} \) within \( R \), and \( f_{\text{BH}}, f_g \) and \( f_{\text{bulge}} \) are respectively the mass fraction of the black hole, the disk, and
the bulge to the whole system. Here, we have assumed that the disk radius is comparable to $R$. If the gravity of the black hole or the bulge is predominant, then $Q > 1$ as readily seen, so that the disk can be stabilized. However, if the cooling is effective to give $c_s \ll V$, then the disk can be gravitationally unstable to a high degree. In some cases, the presence of an appreciable velocity dispersion could help to weaken the instability of the disk. For instance, the nuclear disk in Arp 220 is stabilized in outer regions of $\gtrsim 400$pc due to a large velocity dispersion of $\sim 90$km s$^{-1}$, while the inner region of $\lesssim 400$pc is marginally unstable (Scoville, Yun, & Bryant 1997). Anyhow, if the disk is unstable in some regions, it can fragment into lumps, presumably forming stars. Since the radiation drag works just as a surface effect for optically thick lumps, the effects of the drag force would be negligible for them. Nonetheless, if at least a portion of $1/\tau$ is left in the form of diffuse gas (it is not unlikely because $\tau \gtrsim 10^3$), it may be subject to effective drag. Even if all the matter was processed into stars, the stellar mass loss would supply diffuse gas. Also, the newly formed stars may exert radiative force on the innermost gas disk.

5 Discussion

5.1 Sequential Starbursts

Starbursts might not be coeval, but could be sequential inward. There are a lot of possibilities for the sequential starbursts. The radial radiation flux can compress the outer edge of gas disk, where star formation may be triggered. If the UV radiation produces the D-type ionization front at the outer edge, massive stars could form there (Elmegreen & Lada 1977). If sequential explosions of supernovae occur in an OB association embedded in a plane-stratified gas disk, they build up an elongated hot cavity (superbubble), and consequently the shock-heated gas spurts out from the disk (Shapiro & Field 1976; Tomisaka & Ikeuchi 1986; Norman & Ikeuchi 1989). The superbubble may also compress the outer edge of the disk by shocks, and new stars may be born there. Furthermore, the successive interaction of starburst winds with the inner gas disk might be possible through a similar mechanism to that considered by Unemura & Ikeuchi (1987). Anyhow, if an compressed gas ring is ignited by star-formation, it may induce a further radiative avalanche on the inner gas disk.

5.2 Effects of Stellar Evolution

Recently, it is reported that in starburst regions the stellar initial mass function (IMF) is deficient in low-mass stars, with the cutoff of about $m_1 = 2M_\odot$, and the upper mass limit is inferred to be around $m_u = 40M_\odot$ (Doyon, Puxley, & Joseph 1992; Charlot et al. 1993; Doane & Mathews 1993; Hill et al. 1994; Blandle et al. 1996). Here, we consider the effects of stellar evolution. We assume a Salpeter-type IMF for a mass range of $[m_1, m_u]$ as $\phi(m_\ast) = \phi_0(m_\ast/M_\odot)^{-1-s}$, where $s = 1.35$ and $m_\ast$ is mass of a star. We employ simple relations of the stellar mass to luminosity as $(l_\ast/L_\odot) = (m_\ast/M_\odot)^q$ with
Then, we obtain the total stellar mass and luminosity respectively as
\[ M_q = 3 \times 10^{10} \text{yr} \left( \frac{m_*}{M_\odot} \right)^{-\omega} \text{with } \omega = 2.7 \text{ (Lang 1974)}. \]

Then, we obtain the total stellar mass and luminosity respectively as \( M_* = 1.46 \phi_0 M_\odot \) and \( L_* = 2.2 \phi_0 (82t_7^{0.87} - 1) L_\odot \), where \( t_7 \) is the elapsed time in units of \( 10^7 \text{yr} \) after the ignition of a starburst. These give a large luminosity-to-mass ratio in early stages of evolution as \( L_* / M_* \approx 130t_7^{-0.87} (L_\odot / M_\odot) \), Also, if we assume that stars of > \( 4M_\odot \) are destined to undergo supernova explosions and release the energy radiatively with the efficiency of \( \epsilon_{SN} \) to the rest mass energy, the temporally averaged supernova luminosity is \( L_{SN} = \epsilon_{SN} \dot{M}_{SN} c^2 = 20 \phi_0 t_7^{-0.87} L_\odot \), where \( M_{SN}(t) \) is the total mass of SNs which exploded until \( t \), and \( \epsilon_{SN} = 10^{-4} \) is assumed. Thus, the ratio of \( L_{SN} \) to \( L_* \) is \( L_{SN} / L_* = 0.12 / (1 - 0.01t_7^{0.87}) \). Hence, the luminosity is always dominated by stellar luminosity until the supernova explosions are terminated at \( t_7 = 26 \).

### 5.3 Generation of Seed Magnetic Fields

Although the Compton drag is not predominant in a dusty disk, it is significant in the sense that magnetic fields could be generated, because the Compton drag force is exerted virtually on free electrons (Mishustin & Ruzmaikin 1971; Zeldovich, Ruzmaikin & Sokoloff 1983). The equation of motion for electrons is

\[ - \left( \frac{e}{m_e c} \right) (eE + v_e \times B) - \nu_{e\gamma} v_e + \nu_{ep} (v_p - v_e) = 0, \quad (36) \]

where \( m_e \) is the electron mass, \( v_e \) and \( v_p \) are the electron and proton velocities, \( E \) and \( B \) are the induced electric and magnetic fields, \( \nu_{e\gamma} = m_p / m_e c^2 = 3 \sigma_T L_\odot / 8 \pi m_e c^2 R^2 \), and \( \nu_{ep} = 4.8n_e T^{-3/2} \ln \Lambda \text{ (s}^{-1} \text{)} \) with \( \ln \Lambda \) being the Coulomb logarithm (\( \approx 30 \)). Since \( \nu_{e\gamma} \) is estimated as \( \nu_{e\gamma} \simeq 1.2 \times 10^{-14} \text{s}^{-1} \left( L_\odot / 3 \times 10^{12} L_\odot \right) \left( R_g / 100 \text{pc} \right)^{-2}, \nu_{e\gamma} \ll \nu_{ep}, \) and therefore \( v_e \) is approximately equal to the fluid velocity \( (v_e \simeq v) \). Then, Maxwell equations give

\[ \frac{\partial B}{\partial t} = \frac{m_e c \nu_{e\gamma}}{e} \nabla \times v - \frac{m_e c^2 \nu_{ep}}{4 \pi n_e e^2} \nabla^2 B. \quad (37) \]

Here, one may ignore the second term of the right-hand side for fields of galactic scale, because it arises from the Coulomb interactions (Zeldovich, Ruzmaikin & Sokoloff 1983). Thus, \( B \) is evaluated by

\[ B = \frac{m_e c}{e} \int \nu_{e\gamma} \nabla \times v dt. \quad (38) \]

If the gas disk is in Kepler rotation, then \( |\nabla \times v| \sim v_p / r \), so that

\[ |B| \approx 4.6 \times 10^{-13} \text{Gauss} \left( \frac{L_*}{3 \times 10^{12} L_\odot} \right) \left( \frac{R}{100 \text{pc}} \right)^{-2} \left( \frac{t_{SB}}{10^8 \text{yr}} \right) \left( \frac{M_{dyn}}{10^{10} M_\odot} \right)^{1/2} \left( \frac{r}{\text{pc}} \right)^{-3/2}. \quad (39) \]

Remembering that the radial radiation flux produces shear on the disk surface, these seed magnetic fields could be quickly enhanced due to the combination of differential rotation, turbulence, and magneto-rotational shear instabilities (Balbus & Hawley 1991; Hawley, Gammie & Balbus 1995; Matsumoto & Tajima 1995; Matsumoto et al. 1996), which may be relevant to inner viscous flow.
5.4 Link to an Advection-Dominated Flow

Recently, the solution of an advection-dominated viscous flow has been developed in the prescription of $\alpha$-viscosity (Abramowicz et al. 1995; Chen et al. 1995; Narayan & Yi 1994, 1995; Narayan 1996; Narayan, Kato, & Honma 1997; Chen, Abramowicz, & Lasota 1997). In the present radiation-hydrodynamical model, the strong shear flow which is comparable to rotation velocity is induced on the disk surface, and hence a hot corona with roughly virial temperature may be generated. This is a favorable situation where an advection-dominated accretion is raised. The advection-dominated solution exists when the mass accretion rate is relatively low as $\dot{M} \lesssim \dot{M}_{\text{crit}}$. If a fiducial efficiency of $10\%$ is assumed (Frank, King, & Raine 1992), then the critical rate is given as $\dot{M}_{\text{crit}} \simeq 3\alpha^2L_{E,\text{Th}}(r/100r_g)^{-1/2}/c^2$ at $r \geq 100r_g$ and constant at $r < 100r_g$, where $\alpha$ is the viscosity parameter ($\alpha < 1$, Shakura & Sunyaev 1973), $r_g$ is the Schwarzschild radius of the central black hole and $L_{E,\text{Th}}$ is the Eddington luminosity corresponding to the Thomson scattering. The energy conversion efficiency is $\epsilon \equiv \dot{L}_{\text{adv}}/\dot{M}c^2 \sim 0.1(\dot{M}/\dot{M}_{\text{crit}})$, where $\dot{L}_{\text{adv}}$ is the energy release from an advection-dominated accretion onto a putative massive black hole. When equating $\dot{M}_{\text{crit}}$ with (33), we find a transition radius where the mass accretion by the radiative avalanche links to an advection-dominated flow. The transition radius is given by

$$x_{\text{tr}} \equiv \frac{r_{\text{tr}}}{R} = 2.5 \times 10^{-2}\eta^{-2/5} \left(\frac{\alpha}{0.1}\right)^{4/5} \left(\frac{M_{\text{BH}}}{10^8M_\odot}\right)^{3/5} \left(\frac{L_*}{3 \times 10^{12}L_\odot}\right)^{-2/5} \left(\frac{R}{100\text{pc}}\right)^{-1/5}$$

(40)

where $M_{\text{BH}}$ is the mass of a central black hole. [It is noted that more precisely we should employ the total dynamical mass in relevant scales instead of $M_{\text{BH}}$ (Mineshige & Umemura 1996, 1997), although $M_{\text{dyn}} \sim M_{\text{BH}}$ if $x_{\text{tr}} \ll 1$.] Thus, if $R \sim 100$pc, the radiative avalanche may link to an advection-dominated accretion around $\sim 1$pc.

6 Conclusions

We have investigated fundamental physical processes of radiatively driven mass accretion due to circumnuclear starbursts by solving the radiation-hydrodynamical equations of motion including radiation drag force as well as radiation flux force. The present results are summarized as follows.

1. The mass accretion rate via a radiative avalanche is almost insensitive to a gravitational potential, if the starburst luminosity is comparable to or greater than the Eddington luminosity. In particular, the radial flux force allows the disk gas to accrete even in a rigid-rotation potential. The mass-accretion rate at radius $r$ is given by $\dot{M}(r) = \eta(L_*/c^2)(r/R)^2(\Delta R/R)(1-e^{-\tau})$, where $\eta$ ranges from 0.2 up to 1, $L_*$ and $R$ are respectively the bolometric luminosity and the radius of a starburst ring, $\Delta R$ is the extension of the emission regions, and $\tau$ is the face-on optical depth of a disk. In an optically thick regime, the rate depends upon neither the optical depth nor surface mass density distribution of a disk. The timescale of the radiative avalanche could be shorter than $10^8$yr for a dusty disk.
2. The radiation-hydrodynamical instability could prevent the disk surface from accretion especially in early stages ($\lesssim 5 \times 10^7$ yr) of stellar evolution. In later stages, it would result in forming an obscuring wall of torus shape. The wall indirectly induces a radiative avalanche with diminishing the geometrical dilution ($\Delta R/R$). Also, such a radiatively built-up wall shades the nucleus in edge-on views, so this is another version of the formation of an obscuring torus. In the present mechanism, the angular extension of the obscuring wall is a decreasing function of time according as the starburst luminosity fades out. Thus, the classification of AGN types by anisotropic viewing could be comprehended in the contexts of the evolution of circumnuclear starbursts. This picture will be confronted with observations in further detail elsewhere.

3. For an ionized surface of disk, since the Compton drag works virtually on free electrons, magnetic fields could be generated. Although the magnitude is as small as an order of $\sim 10^{-12}$ Gauss, it could be amplified quickly by the combination of differential rotation, turbulence, and magneto-rotational instabilities, which may be relevant to the inner viscous flow.

4. Since the strong shear flow induced by radiation force would result in hot coronae with roughly the virial temperature, a radiative avalanche is likely to link to an advection-dominated accretion at a point where the mass accretion rate falls to the critical rate, e.g., $\lesssim 1$pc. The link to an advection-dominated flow is possibly related to X-ray properties in ultra-luminous IRAS galaxies.

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Figure 1: The equilibrium radius, $r_{eq}$, of the outer edge of a disk after contraction by the radial radiation flux is plotted against $\Gamma \equiv L_\ast / L_E$ (upper panel), where $L_\ast$ is the bolometric luminosity of a starburst ring and $L_E$ is the Eddington luminosity corresponding to the gravitational mass within the ring radius, $R$. Here, three types of gravitational potentials are assumed, i.e., $n = -1/2$ (Keplerian rotation potential), $n = 0$ (flat-rotation potential), and $n = 1$ (rigid-rotation potential). Also, the fractional difference between an enhanced rotation velocity by the radiation flux and an original rotation velocity of gas disk, $\Delta v_\phi / v_\phi$, is shown at $r_{eq}$ (lower panel).
Figure 2: (a) The mass accretion rate by a radiative avalanche is shown in units of $L_*/c^2$ as a function of radius normalized by the ring radius, $R$, for $\Gamma \equiv L_*/L_E = 10$ (dot-dashed curve), 1 (dotted), 0.1 (dashed), and $10^{-3}$ (solid). Here, an optically thick disk is assumed in a gravitational potential of Keplerian rotation ($n = -1/2$). Filled circles represent the equilibrium radii of the outer-edge shown in Figure 1. A thin solid curve is the previous estimation with negligible effects of the radial flux: $\dot{M}/(L_*/c^2) = (3/8)(r/R)^2[(1-n)/(1+n)]$ (see Paper I). In the vicinity of the ring ($x \rightarrow 1$), all solutions approach an asymptotic value of $\dot{M} = 3L_*/4c^2$, being independent of $\Gamma$ as well as $n$. 
Figure 2: (b) Same as Figure 2a, but for $n = 0$ (flat-rotation potential).
Figure 2: (c) Same as Figure 2a, but for $n = 1$ (rigid-rotation potential). Although the rigid rotation ($n = 1$) results in no mass accretion if a radial flux is negligibly small, the radial flux for $\Gamma \geq 10^{-3}$ can generate an appreciable accretion.
Figure 3: The radius inside which the disk is vertically stable is shown against $\Gamma \equiv L_s/L_E$ for $n = -1/2$, 0, and 1. Outside the radius, an optically thin disk or the outer edge of an optically thick disk could be blown out by the radiation flux force.
Figure 4: Vertical radiation flux force is compared with the vertical gravitational force for $n = -1/2$ (a Keplerian disk) in cylindrical coordinates $(r, z)$. On the hatched sides for each value of $\Gamma \equiv L_*/L_E$, the radiation force overwhelms the gravity. The dotted lines are corresponding to the unstable branches below which gas is accelerated downward by the gravity, while above which gas blows upward due to radiative force. The solid lines are stable branches. The attached ages are that of starburst regions corresponding to each value of $\Gamma$. 