THE STRUCTURE OF THE CENTRAL DISK OF NGC 1068: A CLUMPY DISK MODEL

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ABSTRACT

NGC 1068 is one of the best-studied Seyfert II galaxies, for which the black hole mass has been determined from the Doppler velocities of water masers. We show that the standard $a$-disk model of NGC 1068 gives disk mass between the radii of 0.65 and 1.1 pc (the region from which water maser emission is detected) to be about $7 \times 10^7 M_\odot$ (for $a = 0.1$), more than 4 times the black hole mass, and a Toomre $Q$-parameter for the disk is $\sim 0.001$. This disk is therefore highly self-gravitating and is subject to large-amplitude density fluctuations. We conclude that the standard $a$-viscosity description for the structure of the accretion disk is invalid for NGC 1068. In this paper, we develop a new model for the accretion disk. The disk is considered to be composed of gravitationally bound clumps; accretion in this clumped disk model arises because of gravitational interaction of clumps with each other and the dynamical frictional drag exerted on clumps from the stars in the central region of the galaxy. The clumped disk model provides a self-consistent description of the observations of NGC 1068. The computed temperature and density are within the allowed parameter range for water maser emission, and the rotational velocity in the disk falls off as $r^{-0.35}$.

Subject headings: accretion, accretion disks — galaxies: individual (NGC 1068) — galaxies: kinematics and dynamics — galaxies: Seyfert

1. INTRODUCTION

The discovery of emission from water masers in the central region of the Seyfert II galaxy NGC 1068 provides an opportunity to study the properties of the associated accretion disk. The mass of the central black hole is estimated to be $1.5 \times 10^7 M_\odot$ (Greenhill & Gwinn 1997) from the measurement of the Doppler velocities of the masing spots.

The nucleus of NGC 1068 is heavily obscured, and the luminosity of the central source is determined from the observed flux by modeling the dust obscuration and the scattering of photons by ionized gas in the nuclear region. According to a careful analysis carried out by Pier et al. (1994), the bolometric luminosity of NGC 1068 is estimated to be about $8 \times 10^{44}$ ergs s$^{-1}$; the luminosity is perhaps uncertain by a factor of a few.

We use the black hole mass and the bolometric luminosity to construct the standard viscous disk model for NGC 1068 (§ 2) and find that the disk is highly self-gravitating, thereby rendering the $a$-disk model inapplicable. The effect of the irradiation of the disk from the central source does not modify this conclusion (§ 2).

In § 3 we present a new model for the disk in NGC 1068 composed of gas clumps. The accretion in this case arises as a result of gravitational interaction among clumps (§ 3.1); the structure of the clumped-disk (CD) model of NGC 1068 is described in § 3.2.

The velocity of the masing spots appears to be falling off with distance from the center as $r^{-0.35}$ (Greenhill & Gwinn 1997), which is less rapid than the Keplerian power law of $-0.5$. One possible reason for this could be that the disk of NGC 1068 is sufficiently massive so as to modify the rotation curve. However, we show in § 3 that this is not so. Another possibility is that the flattening of the rotation curve is due to a central star cluster. The influence of the star cluster on the accretion rate and the disk structure is discussed in § 3.2.

2. STANDARD THIN-DISK MODEL FOR NGC 1068

The standard viscous accretion disk model for NGC 1068, including the irradiation of the disk from the central source, is presented below. Throughout this paper we take the mass of the black hole at the center of NGC 1068 to be $1.5 \times 10^7 M_\odot$ (Greenhill & Gwinn 1997) and the bolometric luminosity to be $8 \times 10^{44}$ ergs s$^{-1}$ (Pier et al. 1994).

The theory of thin accretion disk is well developed and is described in a number of review articles and monographs, e.g., Frank, King, & Raine (1992). When the flux intercepted by a disk is not small compared to the local energy-generation rate, such as that expected for the masing disk of NGC 1068, the incident flux must be included in determining the thermal structure of the disk.

The fractional luminosity intercepted by the disk depends on the disk geometry, the scattering of radiation by the coronal gas, etc. For instance, if the dominant source of radiation intercepted by the disk were the scattered radiation from an extended corona, then we might expect the incident flux at the disk to be roughly uniform. On the other hand, for a flaring or a warped disk the radiation intercepted from the source directly might dominate, and the incident radiation in this case depends on the inclination angle of the normal to the disk. We adopt the second model to analyze the disk structure; it is straightforward to consider other possibilities, but we do not pursue these since the main result of this section, viz., the standard $a$-disk model for NGC 1068 is highly unstable to gravitational perturbations, turns out to be independent of whether we include irradiation and radiative forces in the calculation of disk structure.

Let the luminosity of the central source be $L_c$. The flux incident at the disk, $F_{\text{in}}$, a distance $r$ from the central source,
is taken to be
\[ F_{\text{in}}(r) = \frac{L_i}{4\pi r^2} \frac{dH}{dr} = \frac{\beta H L_i}{4\pi r^2}, \tag{1} \]
where \( H \) is the vertical scale height, and the constant factor \( \beta \) is defined by \( dH/dr = \beta H/r \) and determines the fraction of the flux that is intercepted by the disk. The upward energy flux due to mass accretion rate \( \dot{M} \) is
\[ F_{\text{up}}(r) = \frac{3}{8\pi} \Omega^2 \dot{M}, \tag{2} \]
and the ratio of these fluxes \( F_{\text{in}}/F_{\text{up}} \approx \epsilon \beta H/R_{\text{osc}} \), where \( \epsilon \approx 0.1 \) is the efficiency of the conversion of the rest mass to energy by the central source and \( R_{\text{osc}} \) is the Schwarzschild radius of the black hole, \( \Omega = (GM_i/r^3)^{1/3} \) is the angular rotation speed, and \( M_i \) is the mass contained inside the radius \( r \). The effective temperature of the disk at \( r \) is given by
\[ \sigma T_{\text{eff}}^4(r) = F_{\text{up}} + F_{\text{in}} \equiv F_t. \tag{3} \]

The temperature at the midplane of the disk (\( T_c \)) can be calculated by considering the first moment of the radiative transfer equation and making use of the Eddington approximation, and it is given by
\[ T_c^4 = \frac{3}{4} T_{\text{eff}}^4 \left( \frac{F_{\text{up}}}{F_t} \right)^3 \left( \frac{F_0}{3} \frac{F_{\text{up}}}{F_t} - \frac{2}{3} \right), \tag{4} \]
where \( T_0 = \kappa \Sigma \) is the optical depth of the disk. The opacity \( \kappa \), for masing disks, with \( H \_2 \) density and temperature in the range \( 10^8 \text{--} 10^{10} \text{ cm}^{-3} \) and \( 100 \text{--} 1000 \text{ K} \), respectively, is dominated by metal grains and is given by (Bell & Lin 1994)
\[ \kappa \approx 0.1 T^{1/2} \text{ cm}^2 \text{ g}^{-1}. \tag{5} \]

The temperature in the disk midplane is not sensitive to the height where the energy is deposited so long as the optical depth of the disk to the incident radiation is much greater than one.

The solution to the hydrostatic equilibrium equation in the vertical direction, which includes the radiation pressure at the disk surface \( (2F_{\text{in}} + F_{\text{up}})/c \), yields
\[ \rho_c c_s^2 \approx \Omega^2 \Sigma H + \frac{2F_t}{c} \left( 1 - \frac{\tau_0}{\tau_{\text{in}}} \frac{F_{\text{up}}}{F_t} \right), \tag{6} \]
where \( \Sigma \) is the surface mass density, and \( c_s \) and \( \rho_c \) are the sound speed and gas density in the midplane of the disk. Taking the effective viscosity in the disk to be \( \alpha c_s H \), the mass accretion rate \( \dot{M} \) is given by
\[ \dot{M} = \frac{2\pi \alpha \Sigma c_s^2}{\Omega}. \tag{7} \]

Equations (4), (6), and (7) are three equations in three unknowns, viz. \( \Sigma, H, \) and \( T_c \), which are solved numerically, and the results are shown in Figure 1. Also shown in Figure 1, for comparison, are the solutions for \( \beta = 0 \), i.e., when the irradiance of the disk is neglected. The solution in the latter case can be obtained analytically and is given below:
\[ T_c \approx 200 \alpha^{-2/9} \dot{\dot{m}}^{8/9} M_i^{2/9} r^{-1} \text{ K}, \tag{8} \]
\[ n \approx 10^{10} \alpha^{-2/3} \dot{\dot{m}}^{1/3} M_i^{1/6} r^{-3/2} \text{ cm}^{-3}, \tag{9} \]
and the Toomre Q-parameter for the stability of the disk is
\[ Q = \frac{\Omega c_s}{\pi G \Sigma} = 2.4 \times 10^{-3} \alpha^{1/3} \dot{\dot{m}}^{-1/3} M_i^{1/6} r^{-3/2}, \tag{10} \]

where \( \dot{\dot{m}} \) is the accretion rate in terms of the Eddington rate, \( M_i = (10^7 M_\odot), \) and \( r \) is the radial distance from the center in parsecs. For the NGC 1068 system, \( M_i \approx 1.5 \) and \( \dot{\dot{m}} \approx 0.4 \). The inner and the outer radii of the masing disk are at 0.65 and 1.1 pc, respectively.

For the numerical results shown in Figure 1, we took the value of \( \beta \) such that the disk intercepts about 50% of the flux from the central source. In this case the disk temperature is dominated by the incident radiation, and the temperature in the disk midplane is close to \( T_{\text{eff}} \approx 511 \text{ K} \) at \( r = 1 \text{ pc} \). The structure of the disk (\( \Sigma, Q, H, T_c \)) in this case is almost independent of the opacity of the gas, and consequently it is unaffected by any uncertainty in \( \kappa \). The molecular mass of the masing disk, when irradiation is included, is \( \sim 7.0 \times 10^7 M_\odot \) (for \( \alpha = 0.1 \)), and the \( Q \) is about \( 10^{-3} \) (see Fig. 1). For the irradiation-dominated disk, the disk mass decreases with \( \alpha \) as \( \alpha^{-0.9} \), and the \( Q \) increases as \( \alpha^{0.9} \). Since \( Q \) should be greater than 1 for stability, we see that the masing disk of NGC 1068 is highly unstable to gravitational perturbation. A decrease in the irradiation flux makes the disk more unstable.

The main conclusion of this section is that the standard \( \alpha \)-disk model for the NGC 1068 masing disk is inconsistent. The disk, according to this model, is highly unstable and should fragment into clumps. A consistent analysis in this case should include the effect of self-gravity and clump interaction, which is described below.

3. A MODEL FOR CLUMPY SELF-GRAVITATING DISKS

A disk with a small value for \( (1 - Q) \) is likely to develop a spiral structure that can transport angular momentum outward. However, when \( Q < 1 \), as in the case of NGC 1068...
(see § 2), the disk is likely to break up into clumps, and the accretion rate is determined by the gravitational interaction among these clumps. Accretion in self-gravitating disks was considered by Paczyński (1978), Lin & Pringle (1987), Shlosman & Begelman (1987), Shlosman, Begelman, & Frank (1990), and has been investigated more recently by Kumar (1999) in some detail for clumpy disks. In § 3.1 we derive the results that we need to construct a model for CDs, and its application to NGC 1068 is discussed in § 3.2.

3.1. Velocity Dispersion and Accretion Rate in a Clumpy Disk

Consider a disk consisting of clumps of size \( l_c \), and take the mean separation between clumps to be \( d_c \). The tidal radius of a cloud, the distance to which the gravity of the cloud dominates over the gravity of the central mass, is

\[ \frac{d_c}{r_c} \approx \frac{M_c}{M} \]

where

\[ M_c = \frac{\pi r_c^2 \sigma_c l_c^3}{d_c^2} = \frac{\pi r_c^2 \sigma_c l_c^3}{d_c^2} \]

\[ M_c \]

and \( \sigma_c \) is the surface mass density of the cloud at radius \( r_c \).

The mean time interval for a clump to undergo strong gravitational interaction with another clump, with impact parameter \( d \gtrsim d_c \), can be shown to be

\[ t_{\text{int}} \approx \frac{M_c}{r_c^2 \sigma_c l_c^3} \]

The time for a clump to undergo strong gravitational encounter with another clump is

\[ t_{\text{grav}} \approx \frac{M_c}{r_c^2 \sigma_c l_c^3} \]

where

\[ M_c = \frac{\pi r_c^2 \sigma_c l_c^3}{d_c^2} \]

\[ M_c \]

and \( \sigma_c \) is the surface mean mass density of the disk at radius \( r_c \).

The mean time interval for a clump to undergo strong gravitational interaction with another clump, with impact parameter \( d \gtrsim d_c \), can be shown to be

\[ t_{\text{int}} \approx \frac{M_c}{r_c^2 \sigma_c l_c^3} \]

and the change in the displacement amplitude of the epicyclic oscillation of a clump as a result of gravitational interaction is

\[ \delta r \sim \frac{d_c}{r_c} \]

where

\[ d_c \approx \frac{\pi r_c^2 \sigma_c l_c^3}{d_c^2} \]

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The expression for \( \delta r \) is similar to that given by Gammie, Ostriker, & Jog (1991) for the velocity dispersion of molecular clouds in the Galaxy.

We assume that the kinetic energy of the relative motion of colliding clouds is dissipated and their orbit is circularized. Colliding clumps might coalesce so long as the size of the resulting object does not exceed the maximum length for gravitational instability, \( l_{\text{max}} \sim (r M_d)/M_d \); clouds of larger size are susceptible to fragmentation because of the rotational shear of the disk, which limits their size.

The characteristic time for a clump to fall to the center can be estimated using equation (11) and is given by

\[ t_r \approx \frac{G m_c M \Omega}{r^2} \]

and the associated average mass accretion rate is

\[ \dot{M} \approx \Omega(r) \frac{M_c}{r} \]

This result is applicable so long as the density contrast between the clump and the interclump medium is about a factor of 2 or larger, and \( l_c \) is of order \( d_c \).

For \( l_c \ll d_c \), clouds undergo several gravitational encounters with other clouds before undergoing a physical collision, and during this time the random velocity of clouds continues to increase. The rate of increase of velocity dispersion in two-body gravitational interaction is given by

\[ \frac{d^2 v}{dt^2} \approx \frac{G M_c M \Omega}{r^2 v^3} \]

We assume that the scattering gives rise to isotropic velocity dispersion, and so the scale height for the vertical distribution of clumps is

\[ \sim v \Omega \]

The time for a clump to undergo physical collision, assumed to be completely inelastic, is

\[ t_{\text{col}} \sim \Omega^{-1}(d_c/l_c)^2 \]

and therefore the velocity dispersion is

\[ \bar{v}^2 \approx \frac{G m_c}{l_c} \]

This corresponds to the Safronov number being equal to 1. Since the cloud collision time is much greater than \( \Omega^{-1} \), the effective viscosity is suppressed by a factor \( (\Omega_{\text{col}})^2 \) compared to the case where collision frequency is greater than \( \Omega \) (e.g., Goldreich & Tremaine 1978), i.e.,

\[ \nu_e \sim \bar{v}^2 (t_{\text{col}}/\Omega^2) \]

and thus the mass accretion rate is given by

\[ \dot{M} \approx M \Omega \left( \frac{M_d}{M_c} \right) \frac{l_c}{r} \approx \Omega M \left( \frac{M_d}{M_c} \right) \frac{l_c}{r} \]

This accretion rate is smaller than given by equation (15) by a factor of \( (d_c/l_c) \).

We can parameterize the effect of unknown size distribution and separation between clumps on the accretion rate and the velocity dispersion by a dimensionless parameter \( \eta \) and rewrite equations for \( M \) and \( \dot{v} \), in the following form for future use:

\[ \dot{M} = \eta \Omega(r) \frac{M_c}{r \pi M_d} \]

and

\[ \dot{v} = \eta \Omega(r) \frac{M_c}{r \pi M_d} \]

For \( l_c \sim d_c \sim l_{\text{max}} \), we see from equations (16) and (18) that \( \eta \sim 1 \). The effective kinetic viscosity in this case is approximately

\[ \Omega c \sim Q^{-2} \]

where \( Q \sim l_{\text{max}}/H \) is the Toomre Q-parameter, \( H \sim \bar{v}_c/\Omega_c \), same as in the Ansatz suggested by Lin & Pringle (1987) for a self-gravitating disk. We note that the relative velocity of collision between clouds is smaller than their orbital speed by a factor of the ratio of the total mass to the molecular mass, and as long as the cooling time for postcollision gas is less than the time

\[ t_c \approx \frac{G M_c M \Omega}{r^2} \]
between collisions, the clouds are not smeared away because of heating and the disk remains clumpy.

### 3.2. Application to NGC 1068

Equation (19) can be recast in the following form:

\[
\frac{M_r^2}{M_r + M_c} = \psi, \tag{21}
\]

where

\[
\psi = \frac{\pi^2}{\eta^{2/3}} \frac{r M_r^{2/3}}{G^{1/3}}.
\]

The solution to this equation in the case where the central mass dominates is

\[
M_r \approx \psi^{1/2} \left( M_c^{1/2} + \psi^{1/2} \right) \tag{23}
\]

and from equation (20) we find the velocity dispersion of clouds to be

\[
\bar{v}_r \approx \eta^{2/3} (GM_r^{1/3}) \tag{24}
\]

For NGC 1068, the luminosity \( \sim 8 \times 10^{44} \) erg s\(^{-1}\), \( M_r \sim 1.5 \times 10^7 M_\odot \), and \( M_c \sim 8 \times 10^{24} \) g s\(^{-1}\). Substituting these numbers in the above equations, we find

\[
M_r \approx 1.3 \times 10^6 M_\odot r^{1/2} \eta^{-1/3} \tag{25}
\]

and

\[
\bar{v}_r \approx 8 \text{ km s}^{-1} \eta^{2/3}, \tag{26}
\]

where \( r \) is measured in parsecs. Note that the disk mass in this model is more than an order of magnitude smaller than the standard z-disk model discussed in § 2. The velocity dispersion of clumps is larger than the sound speed and is of order the observed scatter in the velocity of masing spots in NGC 1068.

The number density of molecules and the gas temperature can be obtained by solving the hydrostatic and the thermal equilibrium equations for blobs in the vertical direction in the presence of incident flux (see § 2). We find the thermal temperature of blobs to be about 510 K at \( r = 1 \) pc (see Fig. 1), and the density scale height \( H \sim 1.5 \times 10^{16} \) cm. The mean number density of \( H_2 \) molecules \( n \sim M_r/(4\pi r^2 H \mu) \sim 5 \times 10^{18} \text{ cm}^{-3} \). Thus both the number density and the temperature are within the allowed range for water maser emission.\(^2\)

The size of clumps is \( \lesssim r(M_r/M_c) \sim 0.1 \) pc, and their mass is \( \sim 10^3 M_\odot \). Since the Jeans mass is \( \sim 10 M_\odot \), the clumps could have some star formation activity. However, the efficiency of star formation is usually quite low, of order a few percent even for the giant molecular clouds, so we do not expect the gas in the clumps to be converted into stars in the short lifetime of the order of 10\(^9\) yr or less for the clumps in the masing disk of NGC 1068.

Note that the velocity dispersion of blobs is independent of \( r \) (eq. [26]), so the scale height for the vertical distribution of blobs increases as \( r^{3/2} \), provided that the velocity dis-

\(^2\) The allowed range for number density for water maser emission is \( 10^8 - 10^{14} \text{ cm}^{-3} \), and the temperature range is 200–1000 K.

3 The density enhancement of the shocked gas depends on the strength of the magnetic field in the clumps and is of order unity for an equipartition magnetic field.

of clouds is nearly isotropic; this increase of disk thickness with \( r \) is more rapid than in the standard z-disk model considered in the last section. At a distance of 2 pc from the center of NGC 1068, the scale height is about 0.3 pc (for \( \eta = 5 \)), so the flaring disk blocks a significant fraction of the radiation from the central source.

Clouds colliding at a relative speed of 8 km s\(^{-1}\) raise the gas temperature to approximately 10\(^4\) K, which is too hot for maser emission. However, the cooling time of gas at the density of \( \sim 10^9 \) cm\(^{-3}\) is of order 100 yr, which is short compared to both the orbital time and the collision time of order 10\(^4\) yr, ensuring that the disk remains cold and that a steady state solution exists.

The slope of the rotation curve in the disk follows from the use of equation (25):

\[
\frac{d \ln V_{\text{rot}}}{d \ln r} = -0.5 + \frac{1}{2} \frac{d \ln M_r}{d \ln r} = -0.5 + \frac{\psi^{1/2}}{4 M_r^{1/2}} \tag{27}
\]

The disk mass for the NGC 1068 system, in our model, is small \((M_r/M_c \approx 0.1)\), and so the slope of the rotation curve is very close to \(-0.5\).

A slower falloff of the rotation curve requires a more rapid increase of the total mass \((M_r)\) with \( r \) than in the model discussed above. This could arise, for instance, if there is a star cluster at the center. The rotation curve falls off as \( r^{-0.35} \) when the mass of the star cluster within a parsec of the center of NGC 1068 is about 8 \( \times 10^6 M_\odot \). Thatte et al. (1997) have carried out near-infrared speckle imaging of the central 1" of NGC 1068 and conclude that about 6% of the near-infrared light from this region is contributed by a star cluster. They estimate that the stellar mass within 1" (\( \sim 50 \) pc) of the nucleus of NGC 1068 is approximately 6 \( \times 10^6 M_\odot \); if the density in this cluster falls off as \( r^{-2} \), as in a singular isothermal sphere, then the expected stellar mass inside 1 pc is about 1.2 \( \times 10^7 M_\odot \). We note that the stellar mass within 100 pc of our own Galactic center is estimated to be in excess of 5 \( \times 10^8 M_\odot \) (Kormendy & Richstone 1995; Genzel, Hollenbach, & Townes 1994), and the mass enclosed within radius \( r \) is seen to increase almost linearly with \( r \) for \( r \gtrsim 2 \) pc. Thus the possibility of a similar stellar mass cluster at the center of NGC 1068 is not surprising.

We discuss below the effect a star cluster has on the structure of the accretion disk.

#### 3.2.1. Effect of a Star Cluster on Accretion Disk

Let us consider that the stellar mass within a radius \( r \) of the center is \( M_s(r) \). The disk mass, as before, is taken to be \( M_r(r) \), and the central mass is \( M_c \). The use of equation (23), which still applies with \( M_c \) replaced by \((M_c + M_s)\), implies that the disk structure is not much affected at small radii where \( M_s(r) \ll M_c \). At larger radii, where the stellar mass becomes comparable to or exceeds \( M_c \), the mass of the gas disk \((M_s)\) increases with \( r \) as \( r^{3/2}(M_c + M_s)^{1/2} \), which is somewhat more rapid than the case of \( M_s = 0 \) considered in
§ 3.1. However, for the masing disk of NGC 1068, $M_\text{s} \lesssim M_\text{c}$ for $r \lesssim 1.6$ pc, and the effect of the stellar cluster on the disk structure, in particular the number density of molecules and the gas temperature, is small.

The velocity dispersion of clouds due to gravitational encounters is proportional to the local surface mass density $\bar{\sigma}$ of gas and is unaffected by the stellar cluster (see eq. [24]). Therefore, the disk thickness ($H$) at first increases with radius as $r^{3/2}$ and then flattens out when $M_s$ starts to dominate the total mass:

$$H \approx \frac{\bar{v}_r}{\Omega(r)} \approx \frac{\eta r^3 \bar{\sigma}(r)}{M_s}.$$  \hspace{1cm} (28)

For $\bar{\sigma} \propto r^{-3/2}$, expected of the CD model with constant $M$, $H/r \propto r^{1/2}$; however, a more rapid decrease of $\bar{\sigma}$ with $r$ leads to a corresponding decrease of $H/r$.\(^4\)

The accretion rate calculated in § 3.1 is modified because of the dynamical friction suffered by clumps in the disk by the star cluster; the orbits of stars in the cluster are gravitationally perturbed by the clumps, so that on average the density of stars behind a clump is greater than the density in the front. We estimate below the accretion rate that arises as a result of the frictional drag exerted by stars. We assume that clumps are not stretched out by the tidal force of stars, which is valid so long as the cloud mass density is greater than the mean mass density associated with the star cluster, i.e., $n_{\text{H}_2} \gtrsim M_\text{s}/(m_{\text{H}_2} r_d^3)$ or $n_{\text{H}_2} \gtrsim 2 \times 10^8 \text{ cm}^{-3} (M_\text{s}/10^7 M_\odot)/r/1 \text{ pc}^{-3}$; this condition is satisfied for NGC 1068.

The dynamical friction timescale for a clump to fall to the center is (e.g., Binney & Tremaine 1987)

$$t_{\text{df}} \approx \frac{r^2 M_v}{G m_s M_* \ln \Lambda},$$  \hspace{1cm} (29)

where $m_s$ is clump mass, $\Lambda \approx r v^2_{\text{rot}}/(G m_s)$, and $v^2_{\text{rot}} = GM_*/r$. Taking the clump size to be the largest length scale for gravitational instability in a shearing disk, i.e., $l_s \approx \pi^3 G \bar{\sigma}/2$, we find $m_s \approx \pi M_\text{s}/l_s^3$. Thus $\ln \Lambda \approx 3 \ln (M_\text{s}/M_\odot)$ is about 7 for the NGC 1068 system. Substituting these into the above equation, we find

$$t_{\text{df}} \approx \frac{\Omega^{-1}}{10} \left( \frac{M^4}{M_* M^2_\text{s}} \right).$$  \hspace{1cm} (30)

For $M_\star \approx M_\odot/3$ and $M_\star/M_* \approx 0.1$, values applicable to the NGC 1068 system, the dynamical friction time is of the same order as the timescale for clumps to fall to the center owing to a gravitational encounter with other clumps (see eq. [22]). Thus we see that a star cluster of modest mass at the center of NGC 1068 can both give rise to a sub-Keplerian rotation curve, as perhaps observed by the water maser, and also because of its dynamical friction on clumps remove angular momentum of the gaseous disk resulting in an accretion rate that is of the same order as needed to account for the observed luminosity (the total $\dot{M}$ is the sum of the accretion rate owing to gravitational interaction between clumps and the dynamical friction drag exerted by stars).

It was pointed out by Ostriker (1983) that even a smooth disk is subject to frictional drag from a star cluster. He showed that the characteristic drag time for removing angular momentum of a gaseous disk is of the same order as the relaxation time for the star cluster. The accretion rate due to this process, in a form applicable to the NGC 1068 system, is given by

$$\dot{M} \sim M_* \Omega I_0 \left( \frac{r_s^2}{r} \right) \left( \frac{M_*}{m_\star} \right),$$  \hspace{1cm} (31)

where $r_s$ is the radius and $m_\star$ is the mass of stars in the cluster, and $I_0 \sim 10$ is a dimensionless quantity that depends on the ratio of the escape velocity at the surface of a star to the stellar velocity dispersion. We see that the accretion rate resulting from friction drag from stars on a smooth disk is much smaller than the rate resulting from the frictional drag on clumps calculated earlier.

4. CONCLUSION

We find that the structure of the masing disk of NGC 1068, as determined using the standard $\alpha$-viscosity prescription, is both inconsistent with the observations and is self-contradictory. In particular, the disk mass according to this model is $7 \times 10^7 M_\odot$, which is much greater than the black hole mass, and the Toomre $Q$-parameter has an extremely low value of $\sim 10^{-3}$ (for $\alpha = 1$). This makes the disk highly unstable to gravitational perturbations and suggests that it breaks up into clumps, a conclusion that is inconsistent with the smooth disk assumption of standard $\alpha$-disks.

We have described in this paper a different model for the disk. We consider the disk to consist of gas clumps that undergo gravitational interactions with one another leading to an inward accretion of gas. The mass of the masing disk of NGC 1068, according to this model, is about $10^6 M_\odot$, and the velocity dispersion of clumps is about 10 km s$^{-1}$, which is in rough agreement with the observed velocity dispersion of masing spots. The temperature and the density of clumps in this model are approximately 510 K and $5 \times 10^8$ cm$^{-3}$, respectively, which are within the allowed parameter range for maser emission.

However, the rotational velocity of clumps, in this model, is close to the Keplerian value, whereas the observations of masing spots indicate a slower falloff of their rotational velocity (Greenhill & Gwinn 1997). One obvious way these results can be reconciled is if the mass in stars within 1 pc of the center of NGC 1068 is of the order of the black hole mass. Speckle observations of the central 1" region of NGC 1068 in near-infrared in fact suggest that the stellar mass contained within 50 pc of the center is about $5 \times 10^6 M_\odot$ (Thatte et al. 1997), which is sufficient to explain the flattening of the observed rotation curve in the masing disk. The dynamical friction exerted by this star cluster on the clumps in the disk removes angular momentum at a rate that is of the same order as needed for the nearly Eddington accretion rate for the system. Thus the clump-clump gravitational interaction and the dynamical friction drag force on clumps together determine the disk structure, which we find to be consistent with observations.

The CD model considered in this paper for the disk of NGC 1068 should also apply to any active galactic nuclei at a distance of about 1 pc or more from the center, depending on the accretion rate and the luminosity, where the disk becomes self-gravitating and the standard $\alpha$ model is too
inefficient to account for the accretion rate (Shlosman & Begelman 1989).

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