Magnetically elevated accretion disks in active galactic nuclei: broad emission line regions and associated star formation

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ABSTRACT

We propose that the accretion disks fueling active galactic nuclei are supported vertically against gravity by a strong toroidal ($\phi$-direction) magnetic field that develops naturally as the result of an accretion disk dynamo. The magnetic pressure elevates most of the gas carrying the accretion flow at $R$ to large heights $z > 0.1R$ and low densities, while leaving a thin dense layer containing most of the mass — but contributing very little accretion — around the equator. We show that such a disk model leads naturally to the formation of a broad emission line region through thermal instability. Extrapolating to larger radii, we demonstrate that local gravitational instability and associated star formation are strongly suppressed compared to standard disk models for AGN, although star formation in the equatorial zone is predicted for sufficiently high mass supply rates. This new class of accretion disk models thus appears capable of resolving two longstanding puzzles in the theory of AGN fueling: the formation of broad emission line regions and the suppression of fragmentation thought to inhibit accretion at the required rates. We show that the disk of stars that formed in the Galactic Center a few million years ago could have resulted from an episode of magnetically elevated accretion at $\sim 0.1$ of the Eddington limit.

Key words: accretion, accretion discs — black hole physics — Galaxy: nucleus — galaxies: active — quasars: general

1 INTRODUCTION

Simulations of accretion disks, both local (in the “shearing box” approximation: e.g., Brandenburg et al. 1995; Davis et al. 2010; Simon et al. 2012) and global (Beckwith et al. 2011; O’Neill et al. 2011), have long suggested that the operation of the magnetorotational instability (MRI) — essential for local angular momentum transport (Balbus & Hawley 1998) — can also lead to the generation of a large-scale toroidal magnetic field through a dynamo process. The strength of the dynamo field is enhanced by the presence of net magnetic flux (vertical field) threading the disk, which also enhances the rate of angular momentum transport (Hawley et al. 1995). In any case, the resulting toroidal field is much stronger than the imposed or stochastic poloidal field. Recent shearing box simulations (Bai & Stone 2013; Salvesen et al. 2016a) have shown that a vertical field with a magnetic pressure exceeding about 0.1% of the central gas pressure leads to a dynamo field whose pressure dominates the disk everywhere. We call such disks, which are supported against the vertical component of gravity by magnetic pressure, magnetically elevated disks.

The implications of magnetic support for accretion disk phenomenology are profound. Magnetically elevated disks are expected to be much thicker and, for the most part, less dense than standard accretion disks. Consequently, they should be less susceptible to gravitational instability and fragmentation into stars (Pariev et al. 2003; Begelman & Pringle 2007; Gaburov et al. 2012), resolving 1

1 Magnetically elevated disks are distinct from “magnetically arrested disks” (MAD: Narayan et al. 2003; Igumenshchev 2008), which require much larger poloidal fluxes.
a major stumbling block to the application of standard accretion disk theory to active galactic nuclei (AGN) (Kolya
dol et al. 1980; Shlosman & Begelman 1985; Shlosman & Begelman 1989; Goodman 2003). Their larger scale heights, faster in-
flow rates, and larger color corrections are in line with a variety of observations of X-ray binaries and cataclysmic variables (Begelman & Pringle 2005), while the coupling be-
tween poloidal flux and disk structure suggests intriguing new possibilities for explaining state transitions and mag-
netized winds Begelman & Armitage 2013; Begelman et al. 2013).

It is somewhat surprising that disks should be able to maintain the strong levels of magnetization seen in simulations. Horizontal magnetic fields are buoyant and should be continuously escaping, as simulations indeed show (Miller & Stone 2004). The structure and energetics of mag-
netically elevated disks can be modeled analytically in terms of a competition between the creation of toroidal field by MRI and its loss through buoyancy and work done on the gas while escaping Begelman et al. 2013). Because the field is con-
tinuously replenished by the dynamo process associated with MRI, which operates far from the equatorial plane, the magnetic pressure \( p_B \) in an elevated disk declines very gradu-
ally with height, while the density drops off rapidly outside a narrow core region where it is approximately uniform. This means that a disk that is moderately magnetically dom-
inated near the midplane, becomes overwhelmingly domi-
nated by magnetic pressure at a few scale heights above the core region. Although most of the mass in such a disk is confined to the equatorial layer, most of the dissipation and resulting accretion occur far from the midplane, where the density is very low. Shearing box simulations appear to agree qualitatively, and to some extent quantitatively, with this model Salvesen et al. 2016a.

One aspect of the model that has not been tested nu-
merically is the extent to which MRI continues to oper-
ate far from the equator. Applications of the model to
date Begelman & Pringle 2005; Begelman et al. 2013) have adopted the criterion derived by Pessah & Papalois (2003), which states that MRI switches off when the Alfvén speed \( v_{A0} \), associated with the toroidal field exceeds the geometric mean of the Keplerian speed and the gas sound speed. This limits the disk height to \( H \sim (c_s/c_K)^{1/2} R \), where \( c_s \) is the gas sound speed and \( c_K \) is the Keplerian speed, respectively. This limit, however, was derived from a local linear stability analysis assuming no radial or vertical stratification of the magnetic pressure or any other fluid quantity. These must be present if the disk is magnetically supported. In the ab-

ence of results incorporating global and nonlinear effects, we will parametrize our ignorance by treating \( H/R \equiv \xi \) as a constant.

In this paper we apply the Begelman et al. 2013 model for the structure of magnetically elevated disks to the fuel-
ing of AGN. In section 2 we review the model and our cur-
rent best estimates of parameters needed for the application, based on existing local simulations. In section 3 we show that such disks can undergo thermal instability, conducive to the formation of broad emission line regions with the observed properties. We discuss disk self-gravity in section 4, where we estimate threshold mass supply rates for fragmentation to set in, and star formation rates if it does. We argue that the unique properties of magnetically elevated disks may lead to substantial star formation coexisting with accretion in luminous AGN. We discuss our results and conclusions in section 5.

2 MAGNETICALLY ELEVATED DISK MODELS

Conditions on the equatorial plane of a magnetically ele-
vated disk consist of the central magnetic pressure, \( p_{B0} \), gas+radiation pressure, \( p_g \), and density, \( \rho_0 \). There is also the pressure of the poloidal magnetic field, \( p_p = B_p^2/8\pi \), which is assumed to be independent of \( z \). The suite of sim-
ulations performed by Salvesen et al. 2016b) suggests that these parameters are connected by the approximate relation

\[
\rho_0 \approx \frac{10}{p_{B0}^{1/2} p_p^{1/2}} \quad (1)
\]

for \( p_p \gtrsim 10^{-5} p_0 \). Equivalently, the plasma beta parameter on the equator is given by

\[
\beta_0 \equiv \frac{p_0}{p_{B0}} \approx 10^{-2} \frac{p_{B0}}{p_p} \quad (2)
\]

Even if magnetic pressure strongly dominates on the equa-
tor (\( \beta_0 \ll 1 \)), the density and gas pressure scale height on the equatorial plane, \( H_0 \), is given by its usual value in the absence of magnetic support. Anticipating that gas pres-
sure dominates over radiation pressure on the central plane of even the most luminous magnetically elevated disks, we have

\[
H_0 \sim \theta_0^{1/2} \chi^{3/2} \frac{r_g}{x}
\]

where \( \theta_0 \equiv kT_0/\mu c^2 \) for a central temperature \( T_0 \) and mean particle mass \( \mu \) (taken to be 0.6\( m_p \) in our simulations, \( r_g \equiv GM/c^2 = 1.5 \times 10^{8} \) m s cm is the gravitational radius of an \( M = 10^8 M_\odot \) black hole, and \( x \equiv R/r_g \).

According to the Begelman et al. 2013 model, the magnetic pressure varies with height according to

\[
p_B \approx p_{B0} \left( 1 + \frac{2}{4H_0^3} \right)^{-\beta_0/(1+\beta_0)} \approx p_{B0} \left( \frac{z}{2H_0} \right)^{-2\beta_0/(1+\beta_0)} \quad (4)
\]

for \( z \gg H_0 \). In addition to any constraint due to the strong toroidal field — which we have parametrized by assuming that MRI persists up to \( z \approx \xi R -- MRI is also subject to quenching at the original Balbus & Hawley 1991 limit, \( k_x v_A > \sqrt{3} \), where \( k_x \) is the vertical wavenumber and \( v_A = (2p_0/\rho)^{1/2} \) is the Alfvén speed associated with the poloidal field. At \( z \gg H_0 \), this condition is effectively equivalent to \( p_0 \gtrsim p_{B0} \), which is never satisfied if \( \beta_0 \ll 1 \). On the other hand, MRI will be quenched in the equatorial layer, \( z \lesssim H_0 \), if \( p_p > p_0 \), corresponding to \( \beta_0 \ll 0.1 \) (see equation [2]). Because the disk is geometrically thick and MRI-active throughout, it is reasonable to suppose that it is relatively effective in dragging in magnetic flux Lubow et al. 1994; Guilet & Ogilvie 2012), and we will therefore assume that \( \beta_0 \) tends to evolve toward this minimum value. That is to say, in our quantitative estimates below we will assume that \( \beta_0 \approx 0.1 \) at all radii, although we note that most numerical values are insensitive to the value of \( \beta_0 \) provided it is \( \ll 1 \).

We set the flux dissipated on one side of the equator to the energy liberated by accretion:

\[
F_d \approx \alpha \Omega \int_0^R \rho_0 dz \approx \frac{3}{8\pi} \frac{GM \dot{M}}{R^3}
\]

(5)
where $\alpha$ is the Shakura & Sunyaev (1973) viscosity parameter, plausibly $\sim 0.3$ in this highly magnetized limit (Salvesen et al. 2016a) and $\Omega$ is the Keplerian angular velocity. Setting $\beta_0 = 0.1$, we solve equation (5) to obtain

$$p_B \approx 3.0 \times 10^{9} \xi^{-0.8} \alpha_0 m_8^{-2} m \theta_0^{-0.1} x \times 26 \mathrm{erg cm}^{-3} ,$$

where $\dot{m} = M/1M_\odot \mathrm{yr}^{-1}$. The central density is given by

$$\rho_0 = \frac{3 \rho_0 \beta_0 \mu}{kT_0} \approx 3.3 \times 10^{-13} \xi^{-0.8} \alpha_0 m_8^{-2} m \theta_0^{-1.1} x \times 26 \mathrm{g cm}^{-3} .$$

(6)

While the equatorial conditions depend on the thermodynamics of the flow through $\theta_0$, conditions at $z \sim H$, where most of the accretion takes place, are independent of this parameter. Using equation (4), we obtain

$$p_B(H) \approx 3.4 \times 10^9 \xi^{-1} \alpha_0 m_8^{-2} m x^{-2.5} \mathrm{erg cm}^{-3}$$

(8)

and

$$\rho(H) = \frac{1}{\Omega \Omega H} \frac{d \rho}{d z}(H) \approx 6.9 \times 10^{-13} \xi^{-3} \alpha_0 m_8^{-2} m x^{-1.5} \mathrm{g cm}^{-3} .$$

(9)

3 THE BROAD EMISSION LINE REGION

Broad emission lines (BELs) in AGN are thought to arise from clouds or filaments of photoionized gas with a small column density ($\lesssim 0.1$) and substantial column density ($\gtrsim 10^{23} \mathrm{cm}^{-2}$), located at distances $x \sim 10^2 - 10^5$ from the black hole (Davidson 1972). Theoretical photoionization studies have shown that there is a strong observational selection effect for BELs to form under "optimally emitting" conditions of particle density and ionizing photon flux (Baldwin et al. 1981), which translates into a tight correlation between distance from the black hole and luminosity. The theoretical question of why AGN have broad-line regions (BLRs) thus boils down to whether filaments with the optimal density, covering factor and column density exist at the right distance. In this section we argue that magnetically elevated disks will naturally produce such filaments through a thermal instability.

Optimally emitting BEL clouds are characterized by densities $n \sim 10^{10} \rho_{10} \mathrm{cm}^{-3}$ and a broader range of dimensionless ionization parameters

$$U \equiv \frac{\dot{N}_{\text{ion}}}{4 \pi R^2 c n}$$

(10)

with typical values $U \sim 10^{-2}$, where $\dot{N}_{\text{ion}}$ is the production rate of photons capable of ionizing hydrogen. It is difficult to measure $\dot{N}_{\text{ion}}$ directly, but reverberation studies of the broad H$\beta$ line give a dimensionless radius of the BLR,

$$x_{\text{BLR}} \approx 6.2 \times 10^4 L_{45}^{1/2} m_8^{-1}$$

(Bentz et al. 2004), where $10^{45} L_{45} \mathrm{erg s}^{-1}$ is the bolometric luminosity and we adopt the bolometric correction $k_{5100 \lambda} = 8.1$ from Rumsey et al. 2012.

Low-density gas exposed to the AGN continuum will equilibrate to the inverse Compton temperature

$$T_{\text{IC}} = \frac{1}{4 k u} \int u_{\nu} h v d \nu = 10^7 T_{\text{IC}} ,$$

(12)

where $u$ and $u_{\nu}$ are the total and spectral radiation energy density, respectively. But if the density becomes high enough, two-body cooling processes will drive the gas thermally unstable, so that a fraction of it will cool to $10^5 \mathrm{K}$ and form a two-phase medium capable of producing strong UV and optical emission lines (McCray 1979; Krolik et al. 1981).

For $T_{\text{IC}} \gtrsim 1$ the principal collisional cooling mechanism is thermal bremsstrahlung, and the condition for isobaric thermal instability can be written

$$\frac{\rho}{u} > 3.2 \times 10^{-17} T_{\text{IC}}^{1/2} \mathrm{erg cm}^{-3}$$

(13)

(Begelman & McKee 1989), where $\rho$ and $u$ are in cgs units. Using $\rho(H)$ from equation (9) and assuming an accretion efficiency $0.1 \epsilon_{-1}$, this implies that the uppermost layers of the magnetically elevated disk should go thermally unstable for

$$x > x_{\text{th}} \approx 1.0 \times 10^7 \xi^6 \epsilon_{-1} \alpha_0 m_8 T_{\text{IC}}^{3} .$$

(14)

A necessary condition for the formation of a BLR is then $x_{\text{BLR}} > x_{\text{th}}$, or, equivalently,

$$\xi < 0.3 \epsilon_{-1} \alpha_0 m_8^{-3/4} L_{45}^{1/2} ,$$

(15)

This condition is extremely insensitive to all parameters, and is readily satisfied for elevated disk models truncated at $\xi \gtrsim 0.1$, which also provide adequate covering fraction to produce the observed line intensities. Given that at least half the mass in the two-phase regions should be in line-emitting clouds, we obtain column densities

$$N_{\text{BLR}} > 9.3 \times 10^{23} T_{\text{IC}}^{1/2} L_{45}^{1/2} \mathrm{cm}^{-2} ,$$

(16)

which are also adequate to produce the observed lines.

Finally, we check the ability of the hot phase to confine the line-emitting gas at densities inferred to exist in the BLR. Assuming that the two-phase system remains close to the thermal stability threshold (13), we obtain a hot-phase pressure given by

$$p_{\text{hot}} \approx 1.5 \times 10^{-18} T_{\text{IC}}^{2} u_{\text{BLR}}^{2} \rho_{\text{hot}} \mathrm{erg cm}^{-3} ,$$

(17)

where both the density of the hot phase, $\rho_{\text{hot}}$, and the radiation energy density at $x_{\text{BLR}}$, $u_{\text{BLR}} \approx 0.3 \mathrm{cm}^{-3}$, are expressed in cgs units. Since $p_{\text{hot}} < \rho(H)$ measured at $x_{\text{BLR}}$, this gives a lower limit on the hot phase pressure,

$$p_{\text{hot}} > 0.6 \xi^3 \epsilon_{-1} \alpha_0 m_8^{1/2} x_{\text{IC}}^{3} T_{\text{IC}}^{-1/4} L_{45}^{-1/4} \mathrm{erg cm}^{-3} ,$$

(18)

If the line-emitting gas, with a temperature $10^4 T_4 \mathrm{K}$, is in pressure balance with the hot phase this implies a lower limit on the particle density,

$$n > 4.5 \times 10^{11} \xi^3 \epsilon_{-1} \alpha_0 m_8^{1/2} T_{\text{IC}}^{2} T_4^{-1} L_{45}^{-1/4} \mathrm{cm}^{-3} ,$$

(19)

consistent with the inferred densities of the broad-line gas if $\xi \gtrsim 0.1$.

4 SELF-GRAVITY AND STAR FORMATION

We next extrapolate the elevated-disk model to radii where self-gravity and fragmentation of the disk becomes an issue in standard accretion disk models of AGN. Low gas density due to magnetic support is the key idea behind earlier suggestions that magnetically supported disks could evade
self-gravity in AGN (Pariev et al. 2003, Begelman & Pringle 2001, Gaburov et al. 2012). What is new is the realization that the density in a magnetically elevated disk should be strongly stratified with height, with accretion at high $z$ coexisting with a dense, equatorial layer where fragmentation and star formation may well occur. Our objective in this section is to assess the connections between these regions. Strong density stratification means that fragmentation and star formation, when it occurs at all, should be dominated by the equatorial layer. Therefore, we estimate the Toomre Q-parameter using the conditions on the midplane. The condition for gravitational instability is

$$Q \sim \frac{M}{2\pi \rho_0 R^2} < Q_{\text{crit}},$$

(20)

where $Q_{\text{crit}} = 1$ for a fluid dynamical disk but may differ in the presence of a strong toroidal field. (Riols & Latter 2016 find that $Q_{\text{crit}}$ may be as large as $\sim 10$ in two-dimensional MHD simulations of gravitoturbulent shear flows in which MRI is suppressed.) The value of $Q$ depends on the temperature of the equatorial layer, which could be affected by external irradiation (Shlosman & Begelman 1987) and other heating mechanisms (Ginsburg et al. 2016, and references therein). Following Shlosman & Begelman (1987) we assume that the equatorial layer comes into equilibrium with a fraction $\chi = 10^{-2} \chi_{-2}$ of the AGN flux $L/4\pi r^2$. Since the cooling timescale is shorter than $3/\Omega$ (Gammie 2001), instability leads to fragmentation at radii

$$x > x_{\text{frag}} \approx \frac{1}{2} \times 10^7 \xi^{16/19} \left(\frac{\alpha_{0.3}}{Q_{\text{crit}}}\right)^{20/19} \frac{m_{-1}^{11/19} m_{-29/38}}{m},$$

(21)

with a threshold temperature

$$T(x_{\text{frag}}) \approx 300 \xi^{-8/19} \left(\frac{\alpha_{0.3}}{Q_{\text{crit}}}\right)^{-10/19} m_{-4/19}^{12/19} \frac{R}{K},$$

(22)

where we have suppressed weak dependencies $\propto (\chi_{-2} \xi_{-1})^{11/38}$ and $\propto (\chi_{-2} \xi_{-1})^{2/19}$ in equations (21) and (22), respectively.

This estimate is valid inside the black hole’s gravitational radius of influence $r_{\text{BH}}$ within its host galaxy. Modeling the stellar potential of the galactic nucleus as that of an isothermal sphere with velocity dispersion $\sigma = 200 \sigma_{200}$ km s$^{-1}$, we have

$$x_{\text{BH}} = \left(\frac{c}{\sigma}\right)^2 = 2.3 \times 10^6 \sigma_{200}^{-2}.$$  

(23)

Outside $r_{\text{BH}}$, the enclosed mass increases linearly with $R$, and the equatorial density of a magnetically elevated disk carrying a fixed mass flux scales as $\rho_0 \propto R^{-2} T^{-1.1}$. (We note that $x = R/r_{\text{BH}}$ is effectively constant outside $r_{\text{BH}}$, since $r_{\text{BH}}$ is proportional to the enclosed mass.) This means that $Q_0 \propto T^{-1.1}$ with no explicit radial dependence, so the susceptibility to fragmentation increases outside $r_{\text{BH}}$ only if the temperature continues to drop. However, if the temperature levels off to a constant value $T = 100T_2$ K at $R \lesssim r_{\text{BH}}$, then no fragmentation will occur, at any radius, if $R_{\text{frag}} > r_{\text{BH}}$, corresponding to a lower limit on the local accretion rate needed to trigger fragmentation,

$$\dot{M} > 0.13^{4/5} \left(\frac{\alpha_{0.3}}{Q_{\text{crit}}}\right)^{4/5} \sigma_{200}^{11/10} T_2^{11/10} \frac{M_\odot}{\text{yr}^{-1}}.$$

(24)

and an associated AGN luminosity

$$L_{45} > 0.62 \xi^{4/5} \left(\frac{\alpha_{0.3}}{Q_{\text{crit}}}\right)^{4/5} \sigma_{200}^{11/10} T_2^{11/10}.$$  

(25)

We stress that this is the threshold for mass flux carried by the entire thick disk, not just the flux carried by the thin equatorial layer. We also note that this condition applies to the mass flux passing through the region around $r_{\text{BH}}$, which may be higher than the accretion rate reaching the black hole (in which case the luminosity in equation 25 would be lower).

If fragmentation starts in the equatorial layer, we assume that it proceeds all the way to star formation at a rate

$$\dot{\Sigma}_* = \epsilon_\star \Omega_0 \Sigma_0,$$

(26)

per unit area, where $\Sigma_0 = 2\rho_0 H_0$ is the surface density of the equatorial layer and $\epsilon_\star$ is the star formation efficiency. It is reasonable to suppose that feedback from massive stars drives turbulence which regulates the density at close to the critical value for fragmentation, implying $\rho_0 \sim \Omega^2 (2/\sigma_G Q_{\text{crit}})$ and $H_0 \sim \sigma_\star/\Omega$, where $\sigma_\star$ is the turbulent velocity dispersion. We assume that the MRI-driven dynamo continues to operate in the presence of this turbulence, but with the toroidal magnetic field now governed by the turbulent pressure,

$$p_{\text{m}} \sim \beta_0^{-1} \rho_0 \sigma_\star^2 \sim \frac{M v^2}{2\pi \rho_0 Q_{\text{crit}} R^2}.$$  

(27)



We can then use equation (4) to express $v_\star$ in terms of the accretion rate $M$, which is mostly flowing outside the turbulent star-forming layer:

$$v_\star \sim 3.1 (\xi \sigma_{200})^{-0.36} \left(\frac{\alpha_{0.3}}{Q_{\text{crit}}}\right)^{0.19} \left(\frac{\sigma_{200}}{\text{km s}^{-1}}\right)^{0.18} \frac{R}{r_{\text{BH}}}.$$  

(28)

We estimate the star formation efficiency by equating the energy injection rate by supernova remnants to the turbulent dissipation rate $\sim \Omega_0^2 \Sigma_0 v_\star^2$. The energy injection rate per unit area can be parametrized as

$$\dot{\epsilon}_{\text{SN}} \sim \eta_{\text{SN}} \dot{\Sigma}_* E_{\text{SN}} f_{\text{SN}},$$

(29)

where $\eta_{\text{SN}} = 0.01 \eta_{-2} M_\odot \text{yr}^{-1}$ is the number of Type II supernovae per solar mass of star formation, $E_{\text{SN}} = 10^{51} E_{51}$ erg is the initial kinetic energy injected by each supernova, and $f_{\text{SN}}$ is the fraction of that kinetic energy going into turbulence. In the limit of a weak magnetic field, energy injection occurs in the momentum-conserving limit when the blast waves have cooled (Hopkins et al. 2011). For ambient densities in the expected range $\sim 10^8$ cm$^{-3}$, the expansion speed of the blast wave at the start of the snowplow phase (which turns out to be remarkably insensitive to ambient density) is $v_{\text{cool}} \sim 10^3$ km s$^{-1}$, and $f_{\text{SN}} \sim (v_\star/v_{\text{cool}})^{-1/2} \lesssim 0.2$ is the fraction of energy remaining when the typical speed decays to $v_\star$ (Kim & Ostriker 2011). However, in the presence of a strong ambient magnetic field (with $\beta_0 < 1$), most of the supernova expansion energy should go first into magnetic energy, which is harder to radiate away. Thus, we expect $f_{\text{SN}}$ to be much larger than the momentum-conserving limit, perhaps by an order of magnitude. Writing $f_{\text{SN}} = 0.1 f_{-1}$, we obtain a star...
formation efficiency
\[ \epsilon_* \sim \frac{v_*^2}{\eta_{SN} E_{SN}/E_{SN}} = 2 \times 10^{-3} \frac{v_{1.10}^2}{\eta_{-2} E_{51} f_{-1}}. \]  
where \( v_{1.10} = v_1/10 \) km s\(^{-1}\).

Given the expressions for \( v_* \) and \( \epsilon_* \), we can integrate the star formation rate per unit area to obtain the total star formation rate outside \( R \),
\[ M_*(> R) = 2\pi \int_R \xi \rho R dR \frac{4.3\epsilon_* (R) v_*(R) M}{Q_{\text{crit}} R} \propto R^{-0.46}, \]
where we have assumed a power-law behavior for \( v_*(R) \) according to equation (28). We estimate the total star formation rate by evaluating equation (31) at \( R_{\text{frag}} \), with \( v_{1.10} = 0.012T(x_{\text{frag}})^{1/2} \) from equation (22). The result is
\[ \frac{M_*/M}{M} \sim 0.08(\eta_{-2} E_{51} f_{-1})^{-1} \xi^{28/19} \alpha_{0.3}^{-35/19} Q_{\text{crit}}^{5/19} m_8^{5/19} M_{\odot}^{27/38}. \]
Setting \( M_*/M \sim 1 \) effectively places an upper limit to the accretion rate able to reach the black hole — all mass supplied in excess of this limit would go into star formation. Given the fiducial values of the parameters, this limit is much higher than the limiting values predicted by standard accretion disk theory, and could plausibly support accretion in the most powerful known quasars with \( \dot{m} \) of several hundred. However, many of the parameters in this model are very uncertain, and it is important to recognize the effect on the predicted accretion and star formation rates if their values have been over- or underestimated. In particular, larger values of \( f_{SN} \), \( \alpha \), and \( Q_{\text{crit}} \) tend to suppress star formation, as does a top-heavy initial mass function (IMF) leading to a larger value of \( \eta_{SN} \). Such an IMF, with \( \eta_{-2} \) as large as 10, has been inferred for the young stars in the Galactic Center, from X-ray observations [Navakshin & Sunyaev 2003]. On the other hand, smaller values of the overall elevated disk thickness, as measured by \( \xi = H/R \), would promote star formation and provide more stringent constraints on the mass flux reaching the black hole.

5 DISCUSSION AND CONCLUSIONS

We have applied a simple model for the radial and vertical structure of magnetically elevated accretion disks (Begelman et al. 2015) to conditions in AGN and have shown that it provides a natural explanation for the existence and properties of the broad emission line region, while simultaneously providing a framework for understanding the relationship between accretion and star formation in the outer disk.

We argue that the BLR occurs as the result of thermal instability in the upper, accreting layers of the elevated disk, which are irradiated by the central engine. The picture of a two-phase BLR established through thermal instability has been around for decades (McCray 1974; Krolik et al. 1981), what is new is the role of the magnetic field in levitating the gas to heights where it can intercept a significant fraction of the incident radiation, and in regulating its mean density.

It is important to note that the magnetic field is not primarily responsible for confining the line-emitting gas in this picture, in contrast to the model proposed by Rees (1987) in which the BEL gas is fully confined by the pressure of the magnetic field. Here, the line-emitting and Compton-heated phases are in pressure balance with each other, with regions of hot and cold gas strung out along the predominantly toroidal magnetic field lines. Since the magnetic pressure is much larger than the thermal pressure, significant pressure fluctuations are possible transverse to the field, and corresponding density fluctuations have been seen in isothermal simulations of magnetically supported shearing boxes (Salvesen et al. 2016a).

Slightly further from the black hole, where the effective temperature drops below the sublimation temperature (~ 1500 K) for graphite, we expect the levitated gas to be dusty, with a covering factor and column densities consistent with the putative “dusty torus” that reproceses a significant fraction of the AGN continuum into the infrared. However, in contrast to the BLR, gas in this region may be strongly susceptible to radiation pressure forces due to the incident UV flux, and any model of this region will have to take such forces into account (Dorodnitsyn et al. 2012; Chan & Krolik 2014).

The same magnetically elevated disk models, extrapolated to still larger distances, provide a nuanced picture of the relationship between star formation and accretion in the AGN fueling process. Standard thin disk models predict the onset of fragmentation, and interruption of the accretion flow, when \( M_\text{BH} \) is larger than about \( 10^{-7} M_\odot \) yr\(^{-1}\) (Shlosman & Begelman 1987; Goodman 2003), corresponding to AGN luminosities \( \lesssim 10^{43} \) erg s\(^{-1}\). Simple one-zone models for magnetically supported disks, approximating \( \rho \sim \rho(H) \) (equation [9]), yield drastically higher thresholds for fragmentation (e.g., Begelman & Pringle 2007), but fail to take into account the strong vertical density stratification associated with the buoyant escape of magnetic field (Begelman et al. 2013). Incorporating this stratification, we find that the threshold for avoiding fragmentation altogether is not as high as in the one-zone case, but also that there is a new regime in which star formation in the equatorial layer coexists with unimpeded accretion at larger heights. Only when the ratio \( M_*/M \) approaches unity does star formation interfere with and possibly limit the amount of matter reaching the black hole. There seem to be reasonable parameters that permit values of \( M* \) as large as those powering the most luminous quasars.

Although our estimated star formation rates depend on a number of highly uncertain parameters, certain trends are clear. We find that little to no fragmentation should occur inside the black hole’s gravitational radius of influence \( r_{\text{BH}} \) for \( M_* \) less than a few percent of a solar mass yr\(^{-1}\), corresponding to \( L \sim 10^{44} – 10^{45} \) erg s\(^{-1}\). As \( M_* \) increases, star formation first sets in near or somewhat outside \( r_{\text{BH}} \), and migrates inward to radii \( R_{\text{frag}} \propto M^{-0.76} \). The total star formation rate within \( r_{\text{BH}} \) is dominated by conditions near the fragmentation radius, and increases steeply with the accretion rate \( \propto M^{1.71} \), until it becomes of the same order. Thus, star formation in the equatorial zone of a magnetically elevated disk could produce a substantial annulus of stars inside the sphere of influence, the mass of which could
be used to calibrate the luminosity and duration of a prior episode of activity. We can apply our model to the disk of young stars in the Galactic Center (GC), which has a rather well-defined inner edge (Paumard et al. 2006), corresponding to 1.2 × 10^{17} cm at a distance of 8 kpc. For a black hole mass $M = 4 \times 10^6 M_\odot$, this is well within the sphere of influence $r_{\text{BH}} = 5.3 \times 10^{18} \xi_0^{-2}$ cm, normalizing the velocity dispersion in the GC to 100 km s^{-1}. If we identify the inner edge of the stellar disk with $R_{\text{frag}}$, we obtain an accretion rate for the episode that formed the stars

$$\dot{M}_{\text{GC}} \sim 0.2 \xi_0^{3/2} \left( \frac{\alpha_3}{Q_{\text{crit}}} \right)^{40/29} M_\odot \text{yr}^{-1},$$

(33)

where $\xi_0 = 0.1$. For an accretion efficiency of 0.1 and all fiducial parameters equal to one, this would give a luminosity $\sim 10^{45}$ erg s^{-1}, corresponding to twice the Eddington limit for the GC black hole but, as we noted earlier, there are indications that $Q_{\text{crit}}$ might be considerably larger in a strongly magnetized disk (Riols & Latter 2016), which would depress $\dot{M}_{\text{GC}}$. (We also stress that $M_{\text{GC}}$ estimated here is a mass supply rate far from the black hole; the question of how much mass actually reached the black hole is separate.)

Using this accretion rate, we estimate a star formation rate

$$M_{\text{SF,GC}} \sim 6.5 \times 10^{-2} (\eta_{-2} E_3 f_{-1})^{-1} \xi_0^{-2/3} Q_{\text{crit}}^{-2/3} M_\odot \text{yr}^{-1},$$

(34)

For a Salpeter IMF, the total mass in stars formed during this accretion episode was $\sim 1.5 \times 10^4 M_\odot$, implying that the duration of the episode was

$$t_{\text{GC}} \sim 2.3 \times 10^5 (\eta_{-2} E_3 f_{-1})^{-0.42} \xi_0^{-0.52} Q_{\text{crit}}^{-2.62} \text{yr},$$

(35)

which, on the surface, is too short for supernova feedback to come into equilibrium. However, this expression is very sensitive to the value of $Q_{\text{crit}}$, and if we adopt $Q_{\text{crit}} = 3$ (but all other fiducial parameters equal to one), then $t_{\text{GC}}$ rises to 4 Myr, closer to the lifetime of a pre-supernova star. The evolutionary state of the GC stars suggests that they were formed within $\sim 1 - 2$ Myr, implying that supernova feedback was not fully developed during this episode of accretion and allowing for a shorter burst of accretion with a higher star formation rate. We note that if the IMF was very top-heavy, as suggested by Nakayshin & Sunyaev (2005), then the increase in $t_{\text{GC}}$ implied by a larger value of $\eta_{-2}$ would be compensated by the smaller total mass in stars present in the disk.

Magnetically elevated disk accretion is not expected to promote very large enhancements of star formation outside the sphere of influence. A growing body of observational evidence supports a correlation between star formation and AGN luminosity (or $L/L_{\text{Edd}}$), at rates that can reach $10^4$ times the accretion rate (Bonfield et al. 2011; Mullaney et al. 2012; Chen et al. 2013; Delvecchio et al. 2012; Bernhard et al. 2016). ALMA data suggest that the surface star formation rate can be highly concentrated (Wilson et al. 2014; Oteo et al. 2016), although current resolution does not extend to within a few hundred parsecs. Extrapolation of our star formation estimates to $R > r_{\text{BH}}$, however, suggest that $v_{\phi}$ levels off, and the integrated star formation rate increases only logarithmically with radius. A number of factors could render the magnetically elevated disk model — and our quasi-local star formation scenario — irrelevant at large $R$ (even if it applies at smaller radii), including insufficient poloidal magnetic flux, dominance of other mechanisms promoting accretion, such as global self-gravity or low intrinsic angular momentum of the gas supply, or feedback from the AGN itself (Silk 2013).

We stress that this feedback could be positive, enhancing star formation on large scales through several possible mechanisms. For example, the bow shock associated with a jet could pressurize and compress a ring of gas where the jet blows through, triggering star formation (Gaibler et al. 2012). A similar phenomenon holds for the more common nuclear AGN-driven winds in radio-quiet QSOs (Wagner et al. 2013). Widespread compression leading to induced star formation could also be driven by the interaction of the backflow from a jet cocoon with gas clumps in the inner kpc (Antonuccio-Delogu & Silk 2010). The latter phenomenon can self-regulate to give the canonical ratio $\sim 10^3$ of star formation rate to AGN accretion rate (Antonuccio-Delogu et al. 2016), and may also apply for less collimated flows. Indeed, observations confirm a connection between AGN-driven nuclear winds detected in X-ray absorption and fast molecular outflows (potentially hosting star formation) on large scales (Tombesi et al. 2013). These mechanisms can operate concurrently with effects due to magnetic elevation in the inner disk, that serve to reduce $Q$ and regulate star formation inside $r_{\text{BH}}$ to values considerably less than the AGN accretion rate.

Unfortunately, the association of global star formation with AGN activity may be hard to disentangle from magnetically-regulated processes occurring closer in. Once a substantial burst of star formation is triggered, the rapid formation of massive stars and supernovae, and associated negative feedback on gas feeding, will mask any correlation between the AGN fueling process — and associated nuclear star formation — and high rates of star formation on larger scales (Pitchford et al. 2016), in part also because of the different duty cycles associated with these phenomena. The best indication of magnetically elevated accretion may well be relic stellar disks that persist after the activity has subsided, as in the GC.

We caution that this simple model for magnetically elevated disks represents an extrapolation from local (shearing box) simulations (Bai & Stone 2013; Salvesen et al. 2016a). The few global simulations that demonstrate dynamo activity have generally agreed with the shearing box results, but involve either no imposed vertical field (Beckwith et al. 2011; O’Neill et al. 2011) or a weak vertical field (Suzuki & Inutsuka 2014), and therefore do not sample the parameter space of strongly magnetized disks. Even the local simulations have not reached equatorial magnetizations beyond $B_0 \sim 0.3$, although we suggest that $\beta_0 \sim 0.1$ is attainable, and perhaps natural. Given that elevated disks seem to require the presence of a strong vertical field (Salvesen et al. 2016a), it is important to understand how this field is either accumulated from the environment or generated locally through stochastic processes (e.g., Begelman & Armitage 2014). The presence of poloidal flux presents the possibility that a significant fraction of the liberated energy goes directly into a magnetocentrifugal
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The applications to AGN accretion discussed here depend on additional physical processes that are moderately well-understood in the hydrodynamical limit but which have hardly been studied in the highly magnetized (low-β) limit. Specifically, little research has been done on thermal instability or gravitational fragmentation in strongly magnetized disks, nor on supernova feedback in a strongly magnetized medium subject to strong cooling. The simulations needed to tackle these problems will not be cheap; however, significant progress should be possible with currently available computing resources in the near future.

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