Non-blackbody Disks Can Help Explain Inferred AGN Accretion Disk Sizes

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Abstract

If the atmospheric density $\rho_{\text{atm}}$ in the accretion disk of an active galactic nucleus (AGN) is sufficiently low, scattering in the atmosphere can produce a non-blackbody emergent spectrum. For a given bolometric luminosity, at ultraviolet and optical wavelengths such disks have lower fluxes and apparently larger sizes as compared to disks that emit as blackbodies. We show that models in which $\rho_{\text{atm}}$ is a sufficiently low fraction of the interior density $\rho$ can match the AGN STORM observations of NGC 5548 but produce disk spectral energy distributions that peak at shorter wavelengths than observed in luminous AGN in general. Thus, scattering atmospheres can contribute to the explanation for large inferred AGN accretion disk sizes but are unlikely to be the only contributor.

In the appendix section, we present unified equations for the interior $\rho$ and $T$ in gas pressure-dominated regions of a thin accretion disk.

Key words: accretion, accretion disks – galaxies: active – galaxies: individual (NGC 5548) – quasars: general

1. Introduction

The active galactic nucleus (AGN) STORM project has measured the wavelength-dependent time delays with which the accretion disk of NGC 5548 reverberates in response to emission from very close to the black hole (Edelson et al. 2015). The delays scale approximately as $\tau_{\text{delay}}(\lambda) \propto \lambda$ (Faussauch et al. 2016, hereafter F16). This result has been interpreted as evidence for a disk that emits as a blackbody with a surface temperature profile $T_s(r) \propto r^{-1}$ (Starkey et al. 2017), since $\tau_{\text{delay}} \propto r/c$ and the peak wavelength of a blackbody spectrum follows $\lambda \propto T^{-1}$. In contrast, a blackbody disk that locally releases the energy generated by viscous dissipation has a slower decrease of $T_s$ toward larger radii, $T_s(r) \propto r^{-3/4}$ (Shakura & Sunyaev 1973; SS73; Novikov & Thorne 1973; NT73). A blackbody disk thus has $\tau_{\text{delay}}(\lambda) \propto \lambda^{4/3}$, compared to NGC 5548, its time delays would drop more quickly toward shorter wavelengths.

Furthermore, the lags in NGC 5548 imply a disk $\gtrsim 3$ times larger than expected for a standard thin disk with $L/L_{\text{Edd}} \sim 0.1$ (F16). AGN accretion disks with radii $2 \times$ or more larger than expected were first inferred from microlensing observations of quasars (e.g., Eigenbrod et al. 2008; Morgan et al. 2010; Blackburne et al. 2011, 2015). Larger than expected sizes have also been deduced from photometric variability of quasars in Pan-STARRS (Jiang et al. 2017) and DES (Mudd et al. 2017). Proposed explanations for this discrepancy (see Section 4 of Hall et al. 2014) include an inhomogeneous disk with a range of temperatures at each radius (Dexter & Agol 2011; see also Cai et al. 2017), a long timescale for the disk to respond to changing illumination (Gardner & Done 2017), and an increase in the mean measured lag due to reverberation from diffuse gas at large scales (Cackett et al. 2017).

Here, we explore another mechanism that can help explain the above observations: a disk that deviates from a blackbody such that as the radius decreases, $\lambda_{\text{peak}}$ shifts to an increasingly shorter wavelength than for a blackbody of the same total flux. Electron scattering in the disk atmosphere of photons generated by free–free emission within the disk can produce an emergent spectrum with just such a shift (Section 2). As originally discussed in NT73 Section 5.10 and SS73 Section 3a, free–free emission produces shorter-wavelength photons deeper in the disk. If scattering is non-negligible in the disk atmosphere, such photons are less likely to escape; the emergent spectrum becomes non-blackbody (e.g., Czerny & Elvis 1987). To locally radiate all energy generated by viscous dissipation, the disk must heat up: as scattering becomes sufficiently dominant (e.g., at low atmospheric densities), the peak of the spectrum shifts to shorter wavelengths and disks appear larger at a given wavelength than in the blackbody case (e.g., Rauch & Blandford 1991). A small spectral shift in this direction (of insufficient magnitude to match observations) can occur in the emission from disk regions where electron scattering dominates (Section 5.10 of NT73 and Section 3a of SS73). A further shift in this direction (still of insufficient magnitude to match observations) has been seen in previous work on disks with significant magnetic pressure support and thus low-density atmospheres (e.g., Blaes et al. 2006). This proposed explanation for large inferred accretion disk sizes is related to that of Abolmasov (2017), but distinct in that the electron scattering occurs in an accreting inflow in that work and in the disk atmosphere in this model. In Section 3, we explore which model parameters can produce a shift of the magnitude required to match the observations of NGC 5548, and we discuss our results in Section 4.

2. Two Non-blackbody Spectral Models

To examine how the spectrum from a disk with a scattering atmosphere varies with radius when gas pressure dominates, we consider two atmospheric density distributions: constant and exponential. In both cases, we assume the atmospheres to be isothermal with height.

2.1. Constant-density Atmosphere

For the form of the emergent specific flux $F_\lambda(r)$, we adopt the modified blackbody spectrum appropriate for isotropic scattering in a homogeneous medium (Rybicki & Lightman 1986; Equation (7))...
are the electron-scattering opacity and the frequency-dependent absorption opacity, respectively. The factor of \(\pi\) comes from integrating the (assumed isotropic) emergent specific intensity over all solid angles.

The value of \(T_S(r)\) is found by requiring local energy release. At each radius, the flux emitted over all wavelengths must equal the rate at which energy is generated per unit surface area: \[ F_s(r) \, d\nu \, = \, D(r) \, \text{where} \, D(r) \, \text{is given by Equation (12)}. \] For \(\kappa_{\nu, \text{abs}} \gg \kappa_{\text{es}}\), which occurs at sufficiently high densities or sufficiently low temperatures, \(F_s = \pi B_{\nu}(T_s)\) and \(T_S(r) = (D(r)/\sigma_{\text{sb}})^{1/4}\). For \(\kappa_{\nu, \text{abs}} \rightarrow 0\), which occurs at sufficiently low densities or sufficiently high temperatures, \(F_s \rightarrow 2\pi B_{\nu}(T_s)\sqrt{\kappa_{\nu, \text{abs}}/\kappa_{\text{es}}}\), which can deviate strongly from a blackbody spectrum.

### 2.2. Exponential Atmosphere

Zel’dovich & Shakura (1969, ZS69) derive the expression for the specific flux emerging from a semi-infinite isothermal scattering atmosphere with an exponential density distribution \(\rho(z) = \rho_1 e^{-(z-z_1)/h_1}\) where \(h_1\) is the atmospheric scale height and \(z_1\) is the height, above the disk plane, of the edge of the atmosphere. We take \(h_1 = r^2 h_k T_S/G M_p h\) where \(h = (k_BT^2)/(G M_p)^{1/2}\) is the scale height (see the text following Equation (2.28) in S37). Thus, we have

\[
h_1 = \left(\frac{k_B T^2}{G M_p} r\right)^{1/2} = \frac{T_S}{T_f}. \tag{2}
\]

We define

\[
k_{\nu, 1} \equiv 3\kappa_{\text{es}} \kappa_{\nu, \text{abs}} h_1^2 / \rho_1 \tag{3}
\]

which is independent of density because \(\kappa_{\nu, \text{abs}}\), the absorption opacity at the edge of the atmosphere, is proportional to the local density \(\rho_1\) (see Section 2.3). We can then write the emergent specific flux from this exponential atmosphere as

\[
E_s(r) = \frac{\pi B_{\nu}(T_S(r))}{1 + 3\kappa_{\text{es}} h_1 n_1 / 2 k_{\nu, 1}^{1/3}}. \tag{4}
\]

(Equation (A1.9) of ZS69, divided by 2 so that \(E_s \rightarrow \pi B_{\nu}\) when electron scattering is unimportant; see Rybicki & Lightman 1986). In that equation, we define

\[
n_1 \equiv \frac{K_{1/3} (2 k_{\nu, 1}^{1/2} / r_1^{1/2} )^{3/2}} {k_{\nu, 1}^{1/6} r_1^{1/2} K_{-1/3} (2 k_{\nu, 1}^{1/2} / r_1^{1/2} )^{3/2}} > 0 \tag{5}
\]

where the \(K_a()\) are modified Bessel functions of the second kind. When \(\rho_1\) is small the \(n_1\) term can dominate the denominator of \(E_s\) and yield a spectrum differing greatly from a blackbody spectrum.

In the limit \(k_{\nu, 1}^{1/2} / r_1^{1/2} \rightarrow \infty\), \(n_1 \rightarrow 0\) and \(E_s \rightarrow \pi B_{\nu}\). In the limit \(\rho_1 = 0\), \(n_1 \approx 70/51\) to <1% accuracy and we can use the definition of \(k_{\nu, 1}\) to write the limiting form of \(E_s\) as

\[
E_s \approx \frac{\pi B_{\nu} 102k_{\nu, 1}^{1/3}}{210 \kappa_{\text{es}} h_1} \approx \frac{\pi B_{\nu} 17k_{\nu, 1}^{1/3}}{35 \kappa_{\text{es}} h_1} \text{with} \tag{6}
\]

\[
\kappa_{\text{es}} h_1 \approx \left(\frac{512\kappa_{\nu, \text{abs}} h_1^2}{35\rho_1 k_{\nu, 1}^3 h_1^3}\right)^{1/3} \pi B_{\nu}(T_s(r)),
\]

which as expected is a factor of \(\approx 2^{1/3}\) lower than given at the end of Appendix 1 of ZS69.

### 2.3. Parameter Dependences

The non-blackbody frequency dependence of the spectrum emerging from the above atmospheres arises from the frequency dependence of the absorption opacity. We assume \(\kappa_{\nu, \text{abs}}\) is given by the sum of free–free and bound-free absorption opacities, and we approximate them as having the same frequency dependence (Meier 2012, Section 9.3.6.2). Defining \(x \equiv h\nu /k T_f\), the specific absorption opacity can then be written

\[
\kappa_{\nu, \text{abs}} = \kappa_{\text{ff}} + \kappa_{\text{bf}} \equiv \kappa_{\text{ff}}^{\kappa_{\text{bf}}} \rho T^{7/2}_S x^{-3}(1 - e^{-x}) \tag{7}
\]

such that

\[
\kappa_{\text{ff}}^{\kappa_{\text{bf}}} = (3.68 \times 10^{22} \text{ cm}^2 \text{ g}^{-1} \text{ (g cm}^{-3})^{-1} \text{ K}^{7/2} \times \left[X + Y + 1180 Z f(T)\right] (1 + X) \tag{8}
\]

where the H, He, and metal mass fractions are \(X = 0.71\), \(Y = 0.27\), and \(Z = 0.02\), respectively, for solar metallicity (Meier 2012, Section 9.3.), and \(f(T)\) is the fraction of metals that are not ionized (1 for neutral gas and 0 for fully ionized gas). The free–free opacity is dominated by electrons from H and He, while the bound-free opacity is dominated by electrons bound to metals. We take \(f(T) = 0.5\), which yields \(\kappa_{\text{ff}} = 8.04 \times 10^{33} \text{ cm}^2 \text{ g}^{-1}\) with the above mass fractions.

We also adopt \(\kappa_{\text{es}} = 0.2(1 + X) \text{ cm}^2 \text{ g}^{-1}\) (Meier 2012, Section 9.3.6.1); more metal-rich gas with a lower X has fewer electrons per gram and therefore a lower electron-scattering opacity.

We can now write out the full parameter dependences of the specific flux \(F_s(r)\) emerging from a constant-density atmosphere, again using \(x \equiv h\nu /k T_f\) for convenience:

\[
F_s(r) = \frac{(4k_{\nu, 1}^3/c^2h^2)[X^3/(e^x - 1)] T_S^3} {1 + \left[1 + 0.2x^2 T_f^{7/2}/\kappa_{\text{bf}}^{\kappa_{\text{bf}}} \rho(1 - e^{-x})[X + Y + 1180 Z f(T)]\right]^{1/3}} \text{with} \tag{9}
\]

where the \((1 + X)\) dependencies in \(\kappa_{\text{es}}\) and \(\kappa_{\text{bf}}^{\kappa_{\text{bf}}}\) have canceled out in the denominator.

The parameter dependences of the specific flux \(E_s(r)\) emerging from an exponential atmosphere can be written explicitly only in the low-density limit. That limit may not always apply, so we do not give the explicit form here.

In calculating \(\rho\) and \(T\) as part of determining \(E_s(r)\) or \(F_s(r)\), we adopt a mean mass per particle of

\[
\mu_p = \frac{m_p}{2X + 0.75Y + Z[0.56 - 0.5f(T)]} \tag{10}
\]
Figure 1. The wavelength-dependent relative disk reverberation time lag estimated from the response-function-weighted mean radius $r(\lambda) - r(1367 \text{ Å})$ (solid lines), overplotted on the wavelength-dependent time lag data for NGC 5548 (points and error bars, from F16; note that the $U$-band points just shortward of 4000 Å are known to be discrepant). Dotted lines show results for flux-weighted mean radii to illustrate the relative disk sizes expected in microlensing observations. We plot curves for blackbody disks (black) and scattering-atmosphere disks with constant density (red) and exponential density (cyan), all with $L = 0.069 L_{\text{Edd}}$. The density of the scattering atmosphere relative to the average interior density is $\log(\rho_{\text{atm}}/\rho) = -1$ (top), $-2$ and $-3$ (middle), and $-4$ and $-5$ (bottom). At $\rho_{\text{atm}}/\rho = 0$, the relative time lags lie above the blackbody disk value at 1000 Å ≤ $\lambda$ ≤ 6000 Å, and below it at other wavelengths. As $\rho_{\text{atm}}/\rho$ decreases, the relative time lags decrease further at $\lambda$ ≤ 1000 Å and increase at longer wavelengths until they are above the blackbody disk values at all wavelengths plotted (last three panels).
where \( m_p \) is the proton mass.

That expression is appropriate for gas that is completely ionized in H and He but in which metals range from neutral to fully ionized. For the parameters given above, \( \mu_p = 0.614 m_p \).

2.4. Calculating \( T_S \), the Emergent Spectrum, and Mean Measured Radii

The emergent \( F_\nu \) from a disk atmosphere depends on the opacity and thus the density in the atmosphere, which will be

Figure 2. Flux density \( f_\nu \) from an annulus at 60 \( R_{\text{sch}} \) in a disk with \( L = 0.069 L_{\text{edd}} \) around NGC 5548. Panel layout and color key are as in Figure 1. The middle right panel includes color-corrected blackbody curves with the same peak wavelengths as the constant-density and exponential models (dashed red–black with \( T_S = 3.07 \times 10^4 \) K and dashed cyan–black with \( T_S = 2.77 \times 10^4 \) K, respectively).
lower than the average interior density. We adopt an atmospheric density \( \rho_{\text{atm}}(r) \), which is a constant fraction of the interior density \( \rho(r) \). To find \( f_r(r) \) or \( E_r(r) \) thus requires knowing the gas density \( \rho(r) \), which, for gas pressure-dominated regions of an AGN accretion disk, requires knowing the disk interior temperature \( T(r) \). In such regions, we use the analytic expressions for \( T(r) \) and \( \rho(r) \) derived in the Appendix. For radiation pressure-dominated disk regions, we calculate \( \rho(r) \) and \( T(r) \) using Equations (2.11) and (2.12) of SS73, respectively.\(^\text{3}\) The transition between these two regions is the radius at which they yield equal values of \( \rho(r) \). We adopt radiation pressure-dominated values of \( \rho(r) \) and \( T(r) \) interior to that radius, and gas pressure-dominated values exterior to that radius.

Once a value for the atmospheric density \( \rho_{\text{atm}}(r) \) as a fraction of \( \rho(r) \) is adopted, the characteristic disk surface temperature \( T_s(r) \) is determined iteratively by requiring

\[
\int F_r(\rho_{\text{atm}}(r), T_s(r)) \, dr = D(r)
\]

for a constant-density atmosphere, or

\[
\int E_r(\rho_{\text{atm}}(r), T_s(r)) \, dr = D(r)
\]

for an exponential-density atmosphere. Once \( T_s(r) \) is known, the emergent spectrum as a function of \( \nu \) is calculated as the area-weighted integral of \( F_r \) or \( E_r \) over the disk. We present results for a range of atmospheric densities for both non-blackbody spectral models considered.

For comparison with microlensing and reverberation mapping observations, we assume a face-on disk and calculate the flux-weighted and response-function-weighted (Cackett et al. 2007) mean radii, respectively:

\[
r_f(\lambda) = \frac{\int_{r_{\text{in}}}^{r_{\text{out}}} r X_\lambda(r) \, dr}{\int_{r_{\text{in}}}^{r_{\text{out}}} X_\lambda(r) \, dr}
\]

and

\[
r_f(\lambda) = \frac{\int_{r_{\text{in}}}^{r_{\text{out}}} \psi_\lambda(r) \, r \, dr}{\int_{r_{\text{in}}}^{r_{\text{out}}} \psi_\lambda(r) \, dr}
\]

where \( X_\lambda(r) = B_\lambda(r) \) for a standard disk, \( F_r(r) \) for a constant-density scattering atmosphere (Section 2.1), and \( E_r(r) \) for an exponential scattering atmosphere (Section 2.2), and \( \psi_\lambda(r) \) is the response function (for a derivation, see Horne et al. in preparation). We assume instantaneous thermalization and re-emission of incident illumination. We refer to these

\(^3\) There is a typo in Equation (2.12) of SS73, which reads \( r^{-3/4} \) instead of the correct value of \( r^{-1/3} \); see Equation (5.9.10) of NT73.

3. Comparison to NGC 5548 Observations

To compare with the observations of NGC 5548, we adopt a cosmology with \( \Omega_m = 0.28 \), \( \Omega_{\Lambda} = 0.72 \), and \( H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1} \) (Komatsu et al. 2011), which yields a luminosity distance of 74.5 Mpc at its redshift of \( z = 0.017175 \) (De Rosa et al. 2015). We model a disk with viscosity parameter \( \alpha = 0.1 \) around a non-rotating black hole of mass \( M = (6.66 \pm 2.17) \times 10^7 \text{ M}_\odot \) (Pei et al. 2017). The Schwarzschild radius is \( R_{\text{Sch}} = 1.97 \times 10^{13} \text{ cm} = 657 \text{ light-seconds} \) and the Eddington limit is \( L_{\text{Edd}} = 4\pi G M \mu_c c / \sigma_T = 9.79 \times 10^{45} \text{ erg s}^{-1} \), where \( \mu_c = 2 m_p / (1 + X) \) is the mean mass per electron. The corresponding Eddington mass accretion rate is \( M_{\text{Edd}} = L_{\text{Edd}} / (0.1 \eta_{0.1} \cdot c^2) = 3.99 \times 10^{25} \text{ g s}^{-1} \), for our non-rotating black hole with \( r_{\text{in}} = 3 R_{\text{Sch}} \), we adopt the appropriate nonrelativistic value of \( \eta_{0.1} = 5/6 \) (Frank et al. 2002). NGC 5548 has bolometric luminosity \( \approx 10^{45.85} \text{ erg s}^{-1} \) (Woo & Urry 2002) and so we adopt \( L/L_{\text{Edd}} = 0.069 \).

We calculate \( X_\lambda(r) \) for \( 3 < r / R_{\text{Sch}} < 4000 \) for a range of atmospheric densities from \( \log(\rho_{\text{atm}}/\rho) = -0 \) to \(-5 \) for both scattering atmosphere models, and then calculate \( r(\lambda) \) for each. The transition between radiation- and gas-pressure-dominated regions occurs at \( r \approx 223 R_{\text{Sch}} \), as compared to \( r \approx 237 R_{\text{Sch}} \) predicted by SS73 Equation (2.17) or NT73 Equation (5.9.9) for the NGC 5548 parameters adopted in this section. The difference arises because the values of \( T(r) \) and \( \rho(r) \) in gas pressure-dominated regions differ slightly from what is assumed in the approximations used in deriving those equations (see the end of the Appendix).

In Figure 1, for comparison with Figure 5 of F16, we plot the wavelength-dependent relative disk reverberation time lag, estimated from the response-function-weighted mean radii \( r(\lambda) \) (solid lines) for blackbody disks (black) and scattering-atmosphere disks with constant density (red) and exponential density (cyan), all with \( L = 0.069 L_{\text{Edd}} \). At \( \rho_{\text{atm}}/\rho = 0 \), the relative time lags lie above the blackbody disk value at \( 100 \AA \lesssim \lambda \lesssim 6000 \AA \), and below it at other...
As \( \lambda \) decreases, the relative time lags decrease further at \( \lambda \lesssim 1000 \text{ Å} \) and increase at longer wavelengths until they are above the blackbody disk values at all wavelengths plotted (last three panels). As expected from Section 2, at sufficiently low \( \rho_{\text{atm}}/\rho \) the time lags at each radius approach limiting values for an exponential-density atmosphere and continue to increase for a constant-density atmosphere.

The dotted lines in Figure 1 show results for flux-weighted mean radii, to illustrate the relative disk sizes expected in microlensing observations. Such radii are smaller than response-function-weighted radii because the latter give high weight to radii that have not reached the Rayleigh–Jeans limit at a given \( \lambda \).

Scattering atmospheres with sufficiently low relative densities can explain the observed wavelength-dependent reverberation...
signal in NGC 5548. For response-function-weighted radii with log($\rho_{\text{atm}}/\rho$) = −4, for an exponential-density atmosphere we find $r_c(\lambda) \propto \lambda^{1.14 \pm 0.02}$, and for a constant-density atmosphere we find $r_c(\lambda) \propto \lambda^{0.00 \pm 0.02}$ at 1000 < $\lambda$ < 4000 Å and $r_c(\lambda) \propto \lambda^{0.82 \pm 0.02}$ at $\lambda$ > 4000 Å.

In Figure 2, we plot the flux density $f_\nu$ versus $\nu$ from an annulus at 60 $R_{\text{Sch}}$, to illustrate how the emergent spectrum changes with $\rho_{\text{atm}}/\rho$ for both scattering atmosphere models. The middle right panel includes color-corrected blackbody curves with peak wavelengths matching both models. Because both models have more flux at short and long wavelengths than the matching color-corrected blackbody curves, modeling disk emission with the latter (e.g., Davis et al. 2006; Done et al. 2012) will not fully reproduce emission from these scattering atmosphere models.

In the left panel of Figure 3, we plot $f_\nu$ for NGC 5548 versus expectations for disks with $\rho_{\text{atm}} = 10^{-6} \rho$. The constant-density model overpredicts $f_\nu$ by less than a factor of three at the wavelengths plotted. However, panel (5) of Figure 4 shows that the integrated $\lambda F_\lambda$ for such a disk peaks at much shorter wavelengths than for blackbody disks and is a poor match to the optical–UV–X-ray SED of NGC 5548 (Mehdipour et al. 2015). Note that panel (1) of Figure 4 shows the expected deviation of even a standard thin disk’s SED from a sum of blackbodies, as seen in Figure 5.10.1 of NT73 and Figure 3 of SS73. However, this deviation is due almost entirely to emission from radii $r < 20R_{\text{Sch}}$. The differences are small at larger radii; e.g., the SED from $r > 60R_{\text{Sch}}$ is shown in panel (1).

For objects where such SEDs conflict with observations, winds from the disk may reduce $M$ at small radii and thus the luminosity from the inner disk, preventing the SED from peaking at too short a wavelength (Slone & Netzer 2012; Laor & Davis 2014; see also Proga 2005). Alternatively, at small radii the disk may transition to a thick torus whose emission peaks in the far or extreme ultraviolet (e.g., Section 2 of Gardner & Done 2017). As a purely illustrative limiting case of the above, in the right panel of Figure 3 and the lower left panel of Figure 4, we show SEDs for models with log($\rho_{\text{atm}}/\rho$) = −4 but from disk radii $r > 60R_{\text{Sch}}$ only ($r > 11$ light-hours). These SEDs show the relative fluxes of the inner and outer disks at those wavelengths. As shown in Figure 7 (left panel), our model’s predicted mean relative disk reverberation time lags at $\lambda > 1367$ Å (solid lines) are not greatly changed if the emission from $r < 60R_{\text{Sch}}$ is excluded (dashed lines). Absolute time lag predictions, in contrast, are strongly dependent on the UV emission properties of the inner disk. Figure 7 (right panel) shows that our model’s predicted mean absolute disk reverberation time lags with emission from $r < 60R_{\text{Sch}}$ excluded (dashed lines) are closer to the AGN STORM observations of NGC 5548 than our model’s predictions including the full disk (solid lines). Thus, models that suppress optical/UV emission disk from radii $r < 60R_{\text{Sch}}$ to better fit the constraints of observed AGN SEDs.
may be able to simultaneously match inferred relative and absolute disk sizes.

4. Discussion

If the atmospheric density in an AGN accretion disk is sufficiently low, scattering in the atmosphere can produce a non-blackbody emergent spectrum. Disks with these scattering atmospheres have lower brightness temperatures and higher color temperatures at ultraviolet and optical wavelengths than the blackbody temperature of a disk of the same bolometric luminosity; thus, they can in principle reconcile disk theory sizes and flux sizes with microlensing sizes (see, e.g., Section 1 of Hall et al. 2014).

The AGN STORM observations of NGC 5548 are in many —but not all—aspects consistent with a disk with a constant-density scattering atmosphere with density $\rho_{\text{atm}} = 10^{-3} \rho$. In particular, although the observed optical–UV–X-ray SED of NGC 5548 is not well matched by that model (Figure 4), the observed optical–UV flux from NGC 5548 is lower than in all models at our adopted Eddington ratio of $L = 0.069 L_{\text{edd}}$ (Figure 3). That Eddington ratio is based on the bolometric luminosity from Woo & Urry (2002), which was found by summing multiwavelength flux values for NGC 5548 from the literature. Contamination of flux measurements by the host galaxy, particularly at infrared wavelengths, could mean that the true Eddington ratio is lower than we have assumed, which would help explain the lower optical–UV flux level (e.g., Gardner & Done 2017 adopt $L = 0.040 L_{\text{edd}}$).

As a consistency check, we note that the physical density corresponding to $\rho_{\text{atm}} = 10^{-3} \rho$ in our NGC 5548 disk models is $10^{-10} - 10^{-12} \text{g cm}^{-3}$ at $4-4000 \text{R}_{\text{Sch}}$ from the black hole, with a minimum at $8.5 \text{R}_{\text{Sch}}$. The vertical length scale corresponding to $\tau_{\text{es}} = 1$ at such densities is $h_{\text{es}} = 10^{10} - 10^{12} \text{cm}$. The atmospheric scale height/radius ratio $h_{\text{es}}/r$ peaks at 0.18 at $8.5 \text{R}_{\text{Sch}}$ and has a value $h_{\text{es}}/r < 10^{-2}$ at $40 < r/\text{R}_{\text{Sch}} < 4000$. A self-consistent atmosphere model must have $h_{\text{es}}/r \ll 1$, so our constant-density model results at $r < 40 \text{R}_{\text{Sch}}$ should be treated with caution. For an exponential atmosphere with $\rho_{\text{atm}} = 10^{-4} \rho$ and scale height $h_{\text{es}}$ (Equation (2)), we find at all radii that $h_{\text{es}}/r \lesssim 6 \times 10^{-5}$, so $h_{\text{es}} < h_{\text{es}}$. Thus, the semi-infinite exponential atmosphere model we use (Section 2.2) is a reasonable approximation. Compared to an infinite exponential atmosphere, it neglects a region of optical depth only $\tau = r_{\text{es}}/\rho_{\text{atm}} h_{\text{es}} < 1$.

We can compare our $\log(\rho_{\text{atm}}/\rho) = -4$ model predictions to the results of microlensing observations. For flux-weighted radii with $\log(\rho_{\text{atm}}/\rho) = -4$, for an exponential-density atmosphere we find $r_{f}(\lambda) \propto \lambda^{0.09 \pm 0.02}$, and for a constant-density atmosphere we find $r_{f}(\lambda) \propto \lambda^{0.84 \pm 0.02}$. In the Einstein Cross, Eigenbrod et al. (2008) found flux-weighted radii $r_{f} \propto \lambda^{1.2 \pm 0.3}$. In the quasar HE 1104–1805, Blackburne et al. (2015) found $r_{f} \propto \lambda^{0.04 \pm 0.02}$ and a size larger than expected by a factor of 4 at 1900 \AA, as compared to our prediction of a factor $\sim 3$ (Figure 5). In a sample of 11 quasars, Blackburne et al. (2011) found $r_{f} \propto \lambda^{0.17 \pm 0.15}$ and a size larger than expected by a factor of $10^{2.5}$ at 1736 \AA, as compared to our prediction of a factor $\sim 3$. The results from first two studies above are reasonably consistent with both the theoretical prediction of $r_{f} \propto \lambda^{0.73}$ and with our model results; see Hall et al. (2014) for further discussion of the Blackburne et al. (2015) results.

For $\rho_{\text{atm}}$ values low enough for scattering to help explain the inferred large sizes of AGN accretion disks, some mechanism must support the low-density atmosphere above the disk; magnetic pressure support is an obvious possibility. Blaes et al. (2006) have shown that accounting for magnetic pressure yields a more extended accretion disk atmosphere that shifts the emergent SED in a similar manner to our Figure 4; see also Begelman & Pringle (2007), Davis et al. (2009), and Tao & Blaes (2013). Such studies found photospheric densities of $\rho_{\text{atm}} \sim 10^{-3} \rho$ (Figure 6 of Blaes et al. 2006) and $\rho_{\text{atm}} \sim 2 \times 10^{-3} \rho$ (Figure 3 of Davis et al. 2009) in disks around stellar-mass black holes. It remains to be seen if disks around supermassive black holes can reach the lower atmospheric densities required for scattering atmospheres to explain large apparent AGN disk sizes.

The models discussed here predict that scattering atmospheres of different densities will yield different apparent disk sizes and $t_{\text{delay}}(\lambda) \propto \lambda^{3}$ slopes. X. In addition, at low atmospheric densities, these models predict sizes relative to blackbody disks that have maxima at $\lambda \lesssim 1000 \text{ \AA}$ and decrease toward approximately unity at longer wavelengths (Figure 5). Overall, disk sizes are predicted to be more anomalous in the ultraviolet than in the optical, and response-function-weighted radii are predicted to be more discrepant from their corresponding blackbody values than flux-weighted radii, by up to a factor of two and a half even at the density of a standard thin disk (Figure 6). However, we have not accounted for specific features in the bound-free opacity such as the temperature-dependent iron opacity bump (see Jiang et al. 2016 and references therein) or the Balmer and Paschen edges. In particular, if the disk atmospheres contain sufficient H I, disk sizes and time lags at the wavelengths of H I edges may differ noticeably from our predictions. It remains to be seen whether or not such an effect can explain the discrepant U-band lags noted in F16.

If scattering atmospheres are important in AGN disks, dependences of the scattering atmosphere properties on AGN parameters such as $M$ and $\dot{M}$ may help explain why SDSS...
quasars do not show the expected trends of continuum color or emission line properties with such parameters (Bonning et al. 2007, 2013).

The models discussed here do not by themselves match all AGN observations. They produce SEDs that peak at shorter wavelengths than observed in many AGN. While the models can show a non-zero lag between X-ray and ultraviolet emission (e.g., Figure 1 for log(\(\mu_{\text{eff}}/\rho\)) = -4), as seen in NGC 2617 (Shappee et al. 2014) and NGC 4151 (Edelson et al. 2017), the slow increase in the lag from ultraviolet to optical wavelengths in those objects is not reproduced. A disk that transitions at small radii to a torus emitting at extreme ultraviolet wavelengths (e.g., Edelson et al. 2017; Gardner & Done 2017) may alleviate these problems.

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Software: MATLAB.

Appendix

Density and Temperature in Gas Pressure-dominated Regions of an Accretion Disk

We desire equations for \(\rho\) and \(T\) in regions of an accretion disk where \(P = P_{\text{gas}}\) and \(\kappa = \kappa_{\text{es}} + \kappa_{\text{refl}}\). For comparison with SS73, we use cgs units and neglect general relativistic corrections.

We will require \(D(r)\), the rate at which mechanical energy is transformed into heat energy per unit surface area over both sides of the disk. Following Frank et al. (2002), we write

\[
D(r) = \frac{3G\rho M}{8\pi r^3} \left(1 - \frac{r_m}{r}\right) = \frac{3G\rho M}{8\pi r^3} f(r).
\]

where \(M\) is the black hole mass, \(\dot{M}\) is the mass accretion rate, and \(f(r)\) involves the innermost disk radius \(r_m\).

To find \(\rho\), we begin with the radiative energy transport equation \(\dot{\epsilon}_r T^4 = \kappa \Sigma F\) (NT73 Equation (5.8.1c)) where \(T\) is the interior disk temperature, \(\kappa = \kappa_{\text{es}} + \kappa_{\text{refl}}\) is the opacity, \(\Sigma\) is the mass surface density, and \(F(r)\) is the total surface flux emitted at \(r\). We take \(\kappa_{\text{refl}} = \kappa_{\text{refl}}^{\text{cm}} b T^{-7/2}\) so that \(\kappa_{\text{refl}}^{\text{cm}}\) has units of \(\text{cm}^{-2} \text{g}^{-1} \text{cm}^{-3} \text{K}^{7/2}\). We assume that energy released is radiated locally, so that \(F(r) = D(r)\). Substituting in for \(\kappa\), \(\Sigma = 2\rho h\), and \(F(r) = D(r)\),

\[
\dot{\epsilon}_r T^4 = (\kappa_{\text{es}} + \kappa_{\text{refl}}^{\text{cm}} b T^{-7/2}) \times 2\rho h \times \frac{3G\rho M}{8\pi r^3} f(r).
\]

The above is a quadratic equation for \(\rho\). We substitute for the disk half-thickness \(h = (P r^3/\rho G M)^{1/2}\) (NT73 Equation (5.8.1a)) rewritten using \(P = \rho k_B T/\mu_p\) as \(h = (k_B T r^3/\mu_p G M)^{1/2}\), where \(\mu_p\) is the mean mass per particle in grams. The quadratic equation and its solution are:

\[
\begin{align*}
\rho a q^2 + b q + c q &= 0 \text{ with } a_q = \kappa_{\text{refl}}^{\text{cm}} T^{-7/2}; b_q = \kappa_{\text{es}}; \\
c_q &= -4\pi\sigma_b T^7/2\mu_p^{1/2} r^{3/2} / \left(3 (k_B G M)^{1/2} \dot{M} f(r)\right)
\end{align*}
\]

To find \(T\), we use two expressions for the vertically integrated shear stress, \(W = 2\hbar\alpha P = 2\hbar\alpha p k_B T/\mu_p\) (NT73 Equations (5.7.5d) and (5.8.1b)) and \(W = \left[f(r) M (GM/r^3)^{1/2}\right]/2\pi\) (NT73 Equation (5.6.14a)), to find a second equation relating \(\rho\) and \(T\):

\[
\rho = \frac{f(r) M (GM/r^3)}{4\pi\alpha (k_B T/\mu_p)^{3/2}}.
\]

Equating Equations (15) and (16) and solving for \(T\), we find:

\[
T = \left[\frac{\kappa_{\text{refl}}^{\text{cm}} f(r) M (GM/r^3)}{2\pi\kappa_{\text{es}} (k_B T/\mu_p)^{3/2}} \times \left(-1 + \frac{1 + 3\kappa_{\text{refl}}^{\text{cm}} (k_B T/\mu_p)^{1/2}}{3\kappa_{\text{es}} (GM/r^3)^{1/2} \dot{M} f(r)}\right)^{-1}\right]^{1/5}.
\]

We have verified that these equations have the correct numerical coefficients and dependencies on \(r\), \(M\), and \(\alpha\) in the limits \(\kappa_{\text{refl}}^{\text{cm}} \ll \kappa_{\text{es}}\) and \(\kappa_{\text{refl}}^{\text{cm}} \gg \kappa_{\text{es}}\). The limiting solutions are approached slowly; e.g., the disk \(T\) is higher at large radii (SS73 region c) than predicted by the limiting solution because \(T\) has a steeper radial dependence in region b than in region c, resulting in the limiting \(T\) in region c being approached from above. In disk regions where \(P = P_{\text{rad}} + P_{\text{gas}}\) and \(\kappa = \kappa_{\text{es}}\), we searched for analytic solutions for \(\rho\) and \(T\) but found only implicit equations.

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