The spectra of accretion discs in active galactic nuclei

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Accepted 1992 March 5. Received 1992 February 20; in original form 1991 December 18

SUMMARY
We present fully self-consistent calculations of the spectra emitted by accretion discs in active galactic nuclei. We treat the inner, radiation-pressure-dominated regions of standard $\alpha$-discs. Incoherent Compton scattering and emission/absorption of radiation by hydrogen and helium are included in computing the vertical temperature structure within the disc and the emergent spectrum. Comptonization modifies the EUV/soft X-ray continuum and the He ii spectral features. In particular, an important fraction of the emergent luminosity appears in the form of a steep soft X-ray spectrum. We find that the observed UV flux and soft X-ray excess of Mkn 841 are consistent with a face-on accretion disc with $\alpha = 0.1$, surrounding a black hole with $M = 2 \times 10^7 M_\odot$ and $L = 0.28 L_{\text{edd}}$.

Key words: accretion, accretion discs – radiation mechanisms: Compton and inverse Compton – radiative transfer – galaxies: active – galaxies: nuclei – X-rays: galaxies.

1 INTRODUCTION
Active galactic nuclei (AGN) are commonly thought to be powered by accretion onto a massive black hole (see Rees 1984 and references therein). The angular momentum of the infalling material is expected to cause an accretion disc to form around the black hole and produce much of the luminosity in a quasi-thermal bump in the EUV part of the spectrum. Observational evidence for this emission has been accumulating over the past few years from the ‘big blue bumps’ seen in the UV (e.g., Shields 1978; Malkan & Sargent 1982) and the spectral excesses seen in the soft X-ray wavebands (e.g., Arnaud et al. 1985; Wilkes & Elvis 1987; Turner & Pounds 1989; Urry et al. 1989; Kruper, Urry & Canizares 1990; Brinkmann 1992). Direct observation of much of the EUV emission is not possible for low-redshift objects due to photoelectric absorption in our Galaxy. Most distant AGN, such as quasars, are not bright enough to have been well measured in the UV. Other evidence of accretion discs is found in the reflected hard X-ray continuum and iron emission line recently observed in AGN (Pounds et al. 1990; Matsuoka et al. 1990; Nandra et al. 1991) and the shapes of Hα emission lines (Chen, Halpern & Filippenko 1989), as well as in some interpretations of the broad-line region (e.g., Collin-Souffrin & Dumont 1990).

The available data show that the EUV bump is likely to dominate the luminosity of AGN. In view of this, it is particularly important to understand its spectrum and to interpret what it may tell us about the mechanism for energy release and the behaviour of accreting black holes. The simplest approach for calculating an accretion disc spectrum is to assume that each point on the disc radiates like a blackbody at the local effective temperature, $T_{\text{eff}}$. This approach is completely inadequate, however, for calculating the high-energy tail of the emitted spectrum. At soft X-ray frequencies, free–free and bound–free opacities are low and photons can diffuse out from large Thomson depths below the disc surface. These Comptonized photons are represented by a temperature $T_{\text{rad}} > T_{\text{eff}}$, so that thermal emission by the disc is a more important source of soft X-rays than expected from the blackbody approximation.

Shakura & Sunyaev (1973), Czerny & Elvis (1987) and Wandel & Petrosian (1988) have used analytic fitting formulae to discuss the effects of Comptonization on the spectrum emitted by the accretion disc. Most detailed calculations, however, have treated incoherent Compton scattering only approximately (Sun & Malkan 1987; Litchfield, King & Brooker 1989; Laor & Netzer 1989). Here we present calculations of the spectra emitted by AGN accretion discs that are fully self-consistent in the EUV/soft X-ray spectral regime.

2 RADIATION-DOMINATED ACCRETION DISCS
We assume the basic structure for spatially thin accretion discs given by the standard theory of Shakura & Sunyaev (1973). The inner regions of the disc are dominated by
radiation pressure,

\[ P \approx P_{\text{rad}} = \frac{u}{3}, \]  

and Thomson opacity,

\[ \kappa \approx \kappa_T = (0.2 \text{ cm}^2 \text{ g}^{-1})(1 + X), \]

where \( u \) is the radiation energy density and \( X \) is the mass fraction of hydrogen. Shakura & Sunyaev found the bulk properties of the accretion disc under these conditions. At radial distance \( R \) from a central black hole of mass \( M \) accreting matter at rate \( M \), the flux emerging from the surface of the disc is given by

\[ F_0 = \frac{3GM\dot{M}}{8\pi R^2} \left[ 1 - \left(\frac{3R_s}{R}\right)^{1/2}\right], \]

where \( R_s = 2GM/c^2 \) is the Schwarzschild radius. For a spatially thin disc, the equation of hydrostatic equilibrium is

\[ \frac{dP}{d\rho} = \frac{1}{z} \frac{du}{dz} = -\frac{GM\rho}{R^3} z, \]

where \( \rho \) is the gas density and \( z \) is the vertical distance from the mid-plane of the disc. Assuming vertical energy transport via radiative diffusion, the flux at height \( z \) is

\[ F(z) = -\frac{c}{3\rho \kappa_T} \frac{du}{dz} = \frac{GM\rho}{\kappa_T R^3} z, \]

where the second equality follows from equation (4). [Shakura, Sunyaev & Zillitinevich (1978) have shown that convection transports no more than \( \sim 30 \) per cent of the vertical energy flux.] Equations (3) and (5) imply that the surface of the disc is at height

\[ H = \frac{3M_T}{8\pi c} \left[ 1 - \left(\frac{3R_s}{R}\right)^{1/2}\right] \]

above the mid-plane.

Assuming an energy generation rate proportional to the gas density (\( dF/dz \propto \rho \)), equation (5) implies that the density is uniform in the vertical direction:

\[ \rho(z) = \rho_0. \]

Finally, consideration of the angular momentum of the infalling gas gives the relation

\[ \frac{M}{2\pi} \left(\frac{GM}{R^3}\right)^{1/2} \left[ 1 - \left(\frac{3R_s}{R}\right)^{1/2}\right] = 2\int_0^H T_{\Phi} \, dz, \]

where \( T_{\Phi} \) is the tangential stress between adjacent radial layers of the disc. If it is assumed that a simple ‘\( \alpha \) law’,

\[ T_{\Phi} = \alpha P, \]

holds (at least on average) in equation (8), the equations form a closed set and can be solved for the value of \( \rho_0 \). Choosing dimensionless parameters (different from those of Shakura & Sunyaev 1973)

\[ M_0 = \frac{M}{10^8 M_\odot}, \quad \dot{M}_0 = \frac{\dot{M}}{1 M_\odot \text{ yr}^{-1}} \quad \text{and} \quad r = \frac{R}{R_s}, \]

the properties of the disc are summarized by

\[ F_0 = 3.88 \times 10^{13} \frac{\dot{M}_0(1 - \sqrt{3}/r)}{M_0^{3/2} r^3} \text{ erg cm}^{-2} \text{ s}^{-1}, \]

\[ \rho_0 = 1.65 \times 10^{-13} \frac{M_0 r^{3/2}}{\alpha(1 + X)^3 M_0^{1/2}} \text{ g cm}^{-3}, \]

and

\[ h = \frac{H}{R} = (1.70) \frac{M_0(1 + X)(1 - \sqrt{3}/r)}{M_0}. \]

At a given radius, we treat the disc as a uniform slab of gas with density, half-thickness and dynamic heating rate (\( \Gamma_{\text{dyn}} = F_0/\dot{H} \)) given by equations (11) to (13). A composition \( X = 0.75, \ Y = 0.25 \) and \( Z = 0 \) is assumed. The canonical value \( \alpha = 0.1 \) is used.

A full non-LTE calculation of the radiative transfer and the vertical temperature (and excitation) structure is performed for this dynamically heated slab. The radiative transfer is computed using the Fokker–Planck approximation of Kompaneets (1957) to treat the effects of incoherent Compton scattering, and the spatial diffusion approximation to treat the vertical transport:

\[ \frac{\partial n}{\partial t} = \frac{N_e \sigma_T}{m_e c} \frac{1}{\varepsilon} \frac{\partial}{\partial \varepsilon} \left[ \varepsilon^4 \left( n + n^2 + kT \frac{\partial n}{\partial \varepsilon} \right) \right] + \frac{\partial}{\partial \varepsilon} \left[ \frac{c}{3\rho \kappa_T} \frac{\partial n}{\partial \varepsilon} \right] + i_j h^3 c^3 \frac{1}{8\pi^2} \frac{1}{1 - c \rho \kappa_T n}, \]

where \( n \) is the photon occupation number, \( \varepsilon \) is the photon energy, \( N_e \) is the free-electron number density, \( T \) is the gas temperature, \( i_j \) is the spectral emissivity, \( \kappa_T \) is the total opacity and \( \kappa_T \) is the opacity due to free-free and bound-free absorption (see Ross, Weaver & McCray 1978). This approach gives an especially accurate treatment of the radiation at high photon energies (where the free-free opacity is low) that diffuses out from deep layers of the disc. Equation (14) is solved in finite-difference form over a rectangular grid of space and energy points using the variable forward-differencing scheme of Chang & Cooper (1970) to treat the Fokker–Planck operator (first term on the right-hand side).

The inner boundary condition employed for this radiative transfer equation depends on whether or not the disc is effectively thick at high photon energies. At depths greater than the thermalization length, local thermodynamic equilibrium (LTE) is expected to hold (see Rybicki & Lightman 1979). If LTE were to hold throughout the disc, substitution of the condition \( n = n_T^{\text{LTE}} \) into equation (5) would yield (Laor & Netzer 1989)

\[ T_{\text{LTE}} = T_{\text{eff}} \left[ 1 + \frac{3}{4} \frac{T_{\text{LTE}}}{2T_0} \right]^{1/4}, \]

where \( T_T \) is the Thomson depth below the surface and \( T_0 \) is the total Thomson depth at the mid-plane of the disc. To find the thermalization depth, we assume \( T = T_{\text{LTE}} \) and evaluate
the effective optical depth due to free–free absorption,
\[ \tau_{\text{ff}}(e) = \frac{3 \tau_{\text{ff}}(e)}{\tau_i + \tau_{\text{ff}}(e)} \],
(16)
for high-energy photons. As a reference energy, we conservatively choose an energy 10 times greater than the expected Wien energy for the emergent spectrum (see Section 3). The gas is treated down to the point where \( \tau_i = 1 \) and the radiation there is assumed to be LTE at the temperature given by equation (15). However, if \( \tau_i < 1 \) holds all the way to the mid-plane of the disc, the entire half-thickness of the disc is treated and the inner boundary condition consists of setting the spectral flux equal to zero \( (F_\nu = 0) \) at the mid-plane. Notice that equation (15) is only used to set the inner boundary condition for the radiative transfer calculation; the actual gas temperatures are then found self-consistently (see below). The outer boundary condition for the emergent spectral flux has been described by Foster, Ross & Fabian (1986).

Although hydrogen and helium are kept highly ionized by the strong radiation fields within these discs, their bound–free absorption and line emission still have important effects on the emergent spectrum. \( \text{H} \) and \( \text{He}^\ast \) are each treated as a three-level atom (with complete \( \ell \)-mixing) plus continuum. Recombination to \( \text{He} \) is negligible. The complete equations of statistical equilibrium for each level (e.g., see Krolik & McKee 1978) are solved to find the fractional level populations. The reduction of spontaneous line emission by resonance trapping is treated using the escape/destruction probabilities derived by Avrett & Hummer (1965) and Hummer (1968).

The temperature of the gas is heated by the complete equation for thermal equilibrium. All heating and cooling rates due to Compton scattering, free–free processes and atomic processes are included (e.g., see Osterbrock 1974), along with the dynamic heating rate \( (T_{\text{dyn}}) \). The non-LTE solution is achieved when the radiation field, the gas temperature and the atomic level populations are completely self-consistent at each depth within the disc.

We treat the uniform slab of gas as ending abruptly at \( z = H \). As Shakura & Sunyaev (1973) have pointed out, the gas density must decline steeply for \( z > H \), so that a thin \(( \tau \ll 1 \) ) ‘atmosphere’ lies atop the thick \(( \tau \sim 500 \) ) slab of gas. Because of the escape of high-energy photons from the hotter gas below, the ionization parameter (ratio of ionizing radiative flux to gas density) is high in the outer layers of the slab. (For example, radiative heating is found to dominate over viscous heating in our non-LTE calculations for the outer layers of the slab.) With declining density, the ionization parameter will remain high throughout the atmosphere and the gas there will be highly ionized like the gas immediately below it. As a result, there should be little difference between the spectral features calculated with such an atmosphere in place and those computed assuming that the ‘photosphere’ lies within the uniform-density slab of gas instead.

3 DISC SPECTRA

We consider first a black hole with \( M_b = 0.01 \) and \( M = 0.01 \), corresponding to \( L = \frac{1}{2} L \) where \( L \) is the luminosity in the Eddington limit. The emergent spectrum at the radial position \( r = 7 \) along the accretion disc is shown in Fig. 1(a). This is highly representative of the spectra from the inner regions of the disc since the quantity \( F_\nu R^2 \), which measures the contribution to the total luminosity, peaks at \( r = 7 \). The emergent spectrum only agrees with the blackbody approximation at low frequencies. At higher frequencies, the spectrum decreases to a ‘modified blackbody’ (see Rybicki & Lightman 1979), with the \( \text{H} \) Lyman edge and the \( \text{He} \) \( \beta \) Lyman \( \alpha \) and \( \beta \) emission lines evident. The line emission is greatly reduced by the effects of resonance trapping; Compton scattering of line photons out of the narrow resonant line core is largely responsible for allowing some \( \text{He} \) line photons to escape. In the soft X-ray band \(( \nu \geq 4 \times 10^{16} \text{ Hz} \) ), the emergent spectrum exceeds the effective blackbody spectrum and, because of Comptonization, closely resembles a Wien law.

The predicted Wien law shown in Fig. 1(a) was found as follows. On the assumption that the gas temperature was resonably approximated by the LTE value \( T = T_{\text{LTE}} \) given

![Figures](https://academic.oup.com/mnras/article-abstract/258/1/189/967234)

**Figure 1.** Spectra for \( M = 10^8 \) \( M_b \) and \( L/L_{\text{Edd}} = \frac{1}{2} \). (a) Spectral flux at \( R = 7 R_g \). Solid curve: computed emergent flux; dashed curve: effective blackbody spectrum; dash–dot curve: predicted Wien law. (b) Spectral luminosity for entire disc. Lower curves show contributions to total luminosity due to individual annuli; dotted curves indicate blackbody approximations used for outermost annuli. Upper solid curve gives total emitted luminosity; dashed curve shows total expected luminosity using the blackbody approximation.

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by equation (15), the temperature was found at the depth where the Compton y-parameter was unity

\[ y = \frac{4kT}{m_e c^2} \tau_y^2 = 1. \]  

(17)

As high-energy photons diffuse outward from deeper, hotter layers of the disc, incoherent Compton scattering continually shifts the spectrum towards a Wien law at the local temperature. The temperature satisfying equation (17) should represent the lowest temperature for which this Comptonization is complete. Thus a Wien spectrum at this temperature containing the entire emergent flux was assumed. For this model, the Wien temperature is \( 2.4 T_{\text{eff}} \). The soft X-ray (\( \varepsilon \geq 150 \text{ eV} \)) portion of the actual emergent spectrum carries 66 per cent of the total flux and can be seen to agree well with this predicted Wien law.

Fig. 1(b) shows the total spectral luminosity emitted by one face of the accretion disc. The region \( 3 \leq r \leq 300 \) has been divided into 20 annular rings, each emitting approximately 5 per cent of the total luminosity, and the typical emergent spectrum for each ring has been found. Multiplication of these spectral fluxes by the areas of the annuli has given the individual contributions to the spectral luminosity shown by the lower curves in Fig. 1(b). For the outermost annuli with negligible contribution to the spectrum at the He \( \alpha \) Lyman features, the local blackbody approximation has been employed. Since the H \( \alpha \) Lyman features are dominated by emission/absorption at large radii, we have not attempted to treat them accurately. The total spectral luminosity (summed over the 20 annuli) is given by the upper solid curve in Fig. 1(b). No Doppler broadening has been applied, so the total spectrum represents what would be observed for a face-on disc. The frequency at which the total spectrum flattens from a Rayleigh-Jeans law depends on the outer limit chosen for the accretion disc; therefore the spectrum for \( \nu \leq 10^{15} \text{ Hz} \) is only approximate. In the EUV, the spectrum falls off approximately as \( \nu^{-0.5} \). A hump (due to blended Wien laws) appears in the soft X-ray regime, where the emergent spectrum greatly exceeds what would be predicted by the blackbody approximation (dashed curve in the figure). Of the total luminosity, 38 per cent appears as soft X-rays for this model.

Using the least-squares method, we have tried fitting commonly used formulæ to the emergent soft X-ray spectrum. This blend of Wien laws is poorly fitted by an exponential or thermal bremsstrahlung spectrum, but the soft X-ray spectrum above 150 eV is well fitted by a single blackbody with \( kT = 54.0 \text{ eV} \). (By contrast, the spectrum using the blackbody approximation for the entire disc would be fitted by a much softer single blackbody with \( kT = 25.6 \text{ eV} \).) Immediately below 150 eV, however, the true spectrum is extremely flat and is not well fitted by any commonly used formula.

Similar results for the same central mass but half the accretion rate (\( M_\odot = 0.005 \)) are shown in Fig. 2. At \( r = 7 \) the He \( \alpha \) Lyman edge is found to be in emission, as radiative recombination to the ground state dominates over photo-absorption in this case. [Husfeld et al. (1984) have discussed He \( \alpha \) emission edges in the context of hot white dwarf atmospheres.] Also, the soft X-ray spectrum does not as closely match the predicted Wien law. This is because the gas now becomes effectively thick (\( \tau_y = 1 \)) at about the same depth where \( y = 1 \), so that Comptonization is less complete and occurs at a somewhat lower temperature than predicted. In the spectrum for the entire disc, the soft X-ray hump is not so pronounced, although it still contains 26 per cent of the total emergent energy. A single blackbody fit to the spectrum above 150 eV has \( kT = 48.2 \text{ eV} \).

Now consider the spectrum for a higher central mass, \( M_\odot = 0.1 \). Fig. 3(a) shows the total disc spectrum for \( M_\odot = 0.1 \) (\( L = \frac{1}{5} L_{\text{edd}} \)). The He \( \alpha \) spectral features are more pronounced and the soft X-ray hump is reduced to 17 per cent of the emergent luminosity. The spectrum above 150 eV resembles a single blackbody with \( kT = 36.2 \text{ eV} \). Fig. 3(b) shows the effect of cutting the accretion rate in half. Here the soft X-ray spectrum resembles a 32.9-eV blackbody and contains 8 per cent of the emitted energy.

Finally, consider the spectrum for \( M_\odot = 1.0 \). Fig. 4(a) shows the total spectrum for \( L = \frac{1}{5} L_{\text{edd}} \). Although a strong He \( \alpha \) Lyman edge now appears, this edge would be much deeper without incoherent Compton scattering acting to fill it in. To illustrate this effect, Fig. 5 compares the actual emergent spectrum at \( r = 7 \) to the result obtained when

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Compton scattering is treated as coherent in the calculation (except that a single Compton scattering is still allowed to remove resonantly trapped line photons from the narrow line cores). Since Compton heating/cooling are absent, the coherent-scattering model has a somewhat different temperature/excitation structure than the true model, contributing to some of the differences seen. The important point is that the coherent-scattering model exhibits a much deeper He II Lyman edge without Compton down-scattering acting to fill it in (see Ross et al. 1978 for further examples).

In Fig. 4(a), the soft X-ray tail of the emergent spectrum accounts for little more than 3 per cent of the total luminosity and resembles a 23.5-eV blackbody spectrum. Fig. 4(b) shows the disc spectrum for \( L = \frac{1}{2} L_{\text{edd}} \). The soft X-ray tail only contains 0.6 per cent of the emitted energy and resembles a 22.3-eV blackbody spectrum.

There is one important effect that metals, which we have neglected in our calculations, can have on the soft X-ray spectrum. Carbon can produce a strong K-edge in the emergent spectrum. The location of this K-edge in the spectrum will depend on the dominant ion of carbon. Since the L-shell ionization potential of C IV is similar to the

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**Figure 3.** Spectral luminosities for entire discs with \( M = 10^7 M_\odot \).
(a) Accretion rate corresponding to \( L/L_{\text{edd}} = 1 \). (b) Accretion rate corresponding to \( L/L_{\text{edd}} = \frac{1}{2} \). Curves have the same meaning as in Fig. 1(b).

**Figure 4.** Spectral luminosities for entire discs with \( M = 10^8 M_\odot \).
(a) Accretion rate corresponding to \( L/L_{\text{edd}} = 1 \). (b) Accretion rate corresponding to \( L/L_{\text{edd}} = \frac{1}{2} \). Curves have the same meaning as in Fig. 1(b).

**Figure 5.** Spectral flux at \( R = 7R_g \) for \( M = 10^8 M_\odot \) and \( L/L_{\text{edd}} = \frac{1}{2} \). Solid and dotted curves show emergent flux when Compton scattering is treated as incoherent and coherent, respectively, in the calculations.
ionization potential of He II, the fractional abundances of C IV and He II will be similar. Helium is highly ionized in our models, with the fraction of He II typically \( \leq 10^{-7} \). Given the lower total abundance of carbon, \( N_C = 4.5 \times 10^{-3} \ N_{\text{He}} \) (Morrison & McCammon 1983), C IV (and less ionized species) will have a negligible effect on the radiation. The observed K-edge will be at \( \epsilon = 392 \text{ eV (} \nu = 10^{17} \text{ Hz) due to C IV or even} \epsilon = 490 \text{ eV due to C V} \) (rather than the C I K-edge at \( \epsilon = 284 \text{ eV} \)). Since our disc spectra for central masses \( M \geq 10^8 \ M_\odot \) have very low spectral luminosities above \( 10^{17} \text{ Hz} \), the introduction of a C IV K-edge will have no effect on any of our conclusions. Only for the high-\( M \), low-\( M \) case shown in Fig. 1 is the possibility of carbon K-edge absorption likely to be important. The result there will depend on the precise nature of the decrease in density of the material close to the surface of the disc. Caution should consequently be applied when using our model calculations for low-mass, high accretion rate discs.

Summarizing our results, fractions of the emergent flux in the soft X-ray band (\( \epsilon \geq 150 \text{ eV} \)) at \( r = 7 \) are shown as functions of the central mass for three different values of \( L/L_{\text{edd}} \) in Fig. 6(a). The soft X-ray fractions for the entire disc are 0.5–0.6 times these values. The higher the value of \( L/L_{\text{edd}} \) (i.e., \( \dot{M}/M \)) and the lower the central mass, the greater the fraction of energy emitted as soft X-rays. Thus soft X-rays constitute an important fraction of the emergent flux for the wedge-shaped region shown in Fig. 6(b). Here the solid curve shows the value of \( L/L_{\text{edd}} \) for which 10 per cent of the flux at \( r = 7 \) is in the soft X-ray band. The dashed curve at \( L/L_{\text{edd}} = 0.4 \) is the locus of models with \( \tau_e = 1 \) at \( \epsilon = 3kT_{\text{mid}} \), where \( T_{\text{mid}} \) is the expected (LTE) temperature at the mid-plane of the disc. At luminosities higher than this, the gas becomes an inefficient, optically thin bremsstrahlung emitter and temperature inversions occur within the disc. Besides, at \( L/L_{\text{edd}} = 0.4 \) the thin-disc approximation is clearly beginning to fail, since \( h/r = 0.18 \) at \( r = 7 \). The triangles in Fig. 6(b) indicate the models presented in this section.

### 4 APPLICATION TO Mkn 841

For purposes of illustration, we apply our results to observations of the Seyfert 1 galaxy Mkn 841. Arnaud et al. (1985) have presented EXOSAT observations of Mkn 841 made in 1984 July. They found that the 2–20 keV spectrum was well fitted by

\[
f_{\nu} = 9.3 \times 10^{-4} \left( \frac{\nu}{10^{15} \text{ Hz}} \right)^{-0.59} \text{ mJy.} \tag{18}
\]

They also found a total 0.15–2 keV flux of \( 6.7 \times 10^{-11} \text{ erg cm}^{-2} \text{ s}^{-1} \), whereas the extrapolated hard X-ray spectrum only implies a soft X-ray flux of \( 1.1 \times 10^{-11} \text{ erg cm}^{-2} \text{ s}^{-1} \). We assume, as they did, that the excess soft X-ray flux of \( 5.6 \times 10^{-11} \text{ erg cm}^{-2} \text{ s}^{-1} \) is due to emission by the accretion disc.

Furthermore, Elvis et al. (1986) have presented IUE observations of Mkn 841 made in 1983 February. Their UV continuum spectrum (corrected for interstellar reddening but not including their greyshift to match earlier optical data) is shown by the solid curve in Fig. 7. For \( \nu \geq 1.5 \times 10^{15} \text{ Hz} \), the observed spectral flux is \( f_{\nu} = 4 \text{ mJy.} \)

Table 1 shows the fluxes for the models presented in Section 3 scaled to the distance of Mkn 841. Here we have assumed face-on discs at the distance implied by the redshift \( z = 0.0366 \). We use \( H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1} \) as an appropriate (geometric) mean of the currently accepted limits on the value for the Hubble constant. The third column in Table 1 gives the total soft X-ray flux above 150 eV, while the fourth column gives the UV spectral flux at \( 1.5 \times 10^{15} \text{ Hz} \). For \( L/L_{\text{edd}} = \frac{1}{3} \), the total soft X-ray flux rises monotonically with increasing \( M \), even though the fraction of the luminosity emitted in soft X-rays decreases. However, the same is not true for \( L/L_{\text{edd}} = \frac{1}{6} \), where the soft X-ray flux is seen to turn down at the highest value of \( M \). As a result, models with \( L/L_{\text{edd}} = \frac{1}{6} \) cannot produce a soft X-ray flux as large as

![Figure 6](https://academic.oup.com/mnras/article-abstract/258/1/189/967234/fig6)

**Figure 6.** Properties of emergent spectrum at \( R = 7R_\odot \). (a) Fraction of total flux in soft X-rays (\( \epsilon \geq 150 \text{ eV} \)) as a function of central mass for \( L/L_{\text{edd}} = \frac{1}{3} \) (solid curve), \( \frac{1}{6} \) (dashed curve) and \( \frac{1}{10} \) (dotted curve). (b) Limits to thin accretion disc spectra with substantial soft X-ray fractions. Solid curve: locus of models emitting 10 per cent of total flux in soft X-rays at \( R = 7R_\odot \). Below this curve, less than −5 per cent of the total disc luminosity appears in soft X-rays. Dashed curve: locus of models with \( \tau_e = 1 \) at \( \epsilon = 3kT_{\text{mid}} \), where \( T_{\text{mid}} \) is the expected temperature at the mid-plane of the disc. Above this curve, discs are effectively thin (leading to temperature inversions) and geometrically thick. Triangles indicate the models presented in Section 3.
observations were separated by more than a year and the brightness of the source may have varied during that time. Also, we have ignored any UV emission associated with the production of the hard X-rays. What we have shown is that an accretion disc is consistent with the UV-soft X-ray spectrum of Mkn 841; we have not, of course, found a unique solution.

Nevertheless, our results are also consistent with the Einstein observations of Mkn 841 performed by Urry et al. (1989). Using a two-component model consisting of a power law plus a soft blackbody, they found that the ‘ultra-soft’ excess of Mkn 841 was fitted by a 28.8-eV blackbody. For our model with $M=2 \times 10^7 M_\odot$ and $L=0.28 \ L_{\text{edd}}$, the spectrum above 150 eV is well fitted by a blackbody with $kT=30.8$ eV. It is also interesting to note that Laor (1990), using very different assumptions about the disc structure, the rotation of the black hole and treating electron scattering as coherent, fitted the Mkn 841 data with an accretion disc corresponding to $M=4 \times 10^7 M_\odot$ and $L=0.24 \ L_{\text{edd}}$.

We have made no attempt to treat the inclination of the disc. An inclination $\mu = \cos \theta < 1$, in addition to implying Doppler broadening of the spectral features due to the Keplerian motion of the emitting gas, reduces the observed flux by a factor $\mu$. As long as $\mu$ is not too small, this would simply mean choosing a slightly larger value for $M$ and a slightly different value for $L/L_{\text{edd}}$. It is clear, however, that the soft X-ray excess of Mkn 841 can be explained by an accretion disc with a sub-Eddington luminosity.

### 5 DISCUSSION

The inner regions of accretion discs in AGN are found to be strong emitters of soft X-rays, contrary to expectations from the naive blackbody disc approximation. The closer the luminosity to the Eddington limit and the lower the mass of the central black hole, the greater the fraction of the emission in soft X-rays.

Our calculations are for ‘bare’ accretion discs around central masses of $10^6-3 \times 10^8 M_\odot$ and $0.1 \leq L/L_{\text{edd}} \leq 0.4$, covering bolometric luminosities of $1 \times 10^{43-46}$ erg s$^{-1}$. A hot corona above the disc could augment the soft X-rays further by inverse Compton scattering of UV photons (e.g., see Czerny & Elvis 1987; Zdziarski & Coppi 1991). Even without such a corona, however, a bare accretion disc can be expected to produce a steep soft X-ray component in the emitted spectrum like that observed for Mkn 841. If bare accretion discs are found to provide an adequate description to the soft X-ray spectra of AGN, then this would be strong evidence that the harder X-ray emission is due to relativistic electrons which do not cool to the extent that a substantial population of thermal electrons accumulates (i.e., the Lorentz factor of the electrons always exceeds $\sim 2$).

We have so far only carried out numerical calculations for the case in which $\alpha=0.1$. An increase of $\alpha$, but presumably keeping it less than unity, means that the density in the disc is lower at fixed values of $M_h$, $M_0$ and $r$. This shifts the allowed soft X-ray emitting region in Fig. 6(b) to lower values of $L/L_{\text{edd}}$ and higher values of $M_0$. Lower densities also mean that the computed temperature profile does not necessarily decrease monotonically with height in the disc, but may increase at the outside. (No such temperature inversions are found in the $\alpha=0.1$ models presented here). Such effects do
not change the gross properties of the spectrum, but can have a significant influence on the near-UV emission and the H I Lyman emission/absorption. Reducing $a$ increases the density in the disk and so shifts the allowed region in Fig. 6(b) in the opposite sense. The enhanced radiation pressure due to He II absorption might also reduce the gas density in discs around massive objects and/or lead to a wind from the disc. We shall investigate the extent of these effects in future work.

Note that we have not included general relativistic effects in our work, except to truncate the disc at $r = 3$, following Shakura & Sunyaev (1973). In particular, we expect more soft X-ray emission from accretion discs surrounding Kerr black holes, since they extend closer to the central object (Novikov & Thorne 1973; Laor & Netzer 1989). Gravitational lensing of light rays also has an effect on the observed spectrum (e.g., Sun & Malkan 1989).

Complex He II emission and absorption is found in our model spectra of discs around massive black holes ($M_\ast \sim 1$). The Hubble Space Telescope may be able to detect such features in bright, high-redshift quasars. A less direct handle on the spectrum in this band in many more objects should be obtainable from its ionizing influence (particularly emission-line ratios) on the surrounding broad-line region (Fabian et al. 1986; Binette, Courvoisier & Robinson 1988).

Several arguments against simple accretion discs in AGN have been put forward recently. These are: (i) no H I Lyman edge has yet been observed when models predict a substantial feature (Antonucci, Kinney & Ford 1989; Czerny & Pompianischi 1990); (ii) detailed studies of the response of emission lines of different ionization levels in the broad-line region in NGC 5548 and other AGN do not show the time delays expected if the ionizing continuum originates in an accretion disc (Peterson et al. 1991; Krolik et al. 1991); and (iii) the observed levels, or limits, on continuum polarization are less than expected from an emission region dominated by electron scattering (Antonucci 1988).

Our work does not refute these arguments, but we make the following comments. On point (i), our calculations are least accurate for the Lyman features of hydrogen (our aim was to model the EUV spectrum), so by themselves are not a good guide at those wavelengths, which carry a small fraction of the bolometric luminosity of the disc. On point (ii) we find that the spectra extend over a broader wavelength range than for a blackbody disc. Therefore, photons that lead to efficient ionization ($13.6 \lesssim \varepsilon \lesssim 50$ eV) originate from a narrower range of disc radii than in a simple blackbody disc. On point (iii), Laor, Netzer & Piran (1990) have shown that a combination of opacity and relativistic effects suppresses the degree of polarization to a level consistent with observations.

Potentially of greater importance are the effects of hard X-ray irradiation of the surface of the disc, density inhomogeneities and differences in the value of $a$. The observations of the fluorescent iron line and reflected continuum in X-ray spectra of AGN (Pounds et al. 1990; Matsuoka et al. 1990; Nandra et al. 1991) suggest that the surface of the disc is exposed to an intense hard X-ray flux. This means that the surface layers should be hotter than we have calculated. The energy of the fluorescent iron line (6.4 keV) requires that the gas density is higher than that of an $a = 0.1$ disc (George, Fabian & Ross 1991), perhaps involving so-called $\beta$-discs in which the tangential stress in the disc is proportional to the gas pressure rather than the total pressure (Sakimoto & Coroniti 1981) as assumed in equation (9). We intend to calculate the spectra of irradiated discs in future work.

ACKNOWLEDGMENTS

We thank Mark Lepper for assistance with the FORTRAN programming. Part of this work was performed while RRR was on a Faculty Fellowship from the College of the Holy Cross. ACF acknowledges financial support by the Royal Society.

REFERENCES


Spectra of AGN accretion discs


