Thermal Conduction and Thermal Instability in the Transition Layer between an Accretion Disc and a Corona in AGN

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ABSTRACT

We study the vertical structure of the transition layer between an accretion disc and corona in the context of the existence of two-phase medium in thermally unstable regions. The disc is illuminated by hard X-ray radiation and satisfies the condition of hydrostatic equilibrium. We take into account the energy exchange between hot corona ($\sim 10^8$ K) and cool disc ($\sim 10^4$ K) through the radiative processes and due to thermal conduction. We make local stability analysis of the case with conductivity and we conclude that thermal conduction does not suppress thermal instability. In spite of continuous temperature profile $T(\tau)$ there are regions with strong temperature gradient where spontaneous perturbations can lead to cloud condensation in the transition layer. We determine the minimum size $\lambda_{TC}$ of such a perturbation.

Key words: accretion, accretion discs – galaxies: Seyfert – atomic processes, conduction, instabilities.

1 INTRODUCTION

The X-ray emission from active galactic nuclei (AGN) has been observed for more than thirty years. Although a lot of research were done on this topic in the case of radio-quiet AGN we still don’t know neither the mechanism leading to formation of the hot X-ray emitting plasma nor the geometry of the X-ray source (for a review, see Mushotzky, Done & Pounds 1993, Czerny 1994).

Recent X-ray data show that a typical X-ray spectrum of a Seyfert 1 galaxy consists of a power law component with energy index $\sim 0.8-1.0$, high energy cutoff above $\sim 200$ keV, and a reflection component characterized by low-energy cut-off due to the absorption by cold matter, high-energy cut-off due to the Klein-Nishina effect and the presence of the Kα line originating in a cold medium (e.g. Pounds et al. 1990, Zdziarski et al. 1995).

Underlying power-low corresponds to emission from hot region and this radiation is most probably thermal (annihilation line for par production wasn’t observed in any spectrum, Maciolek-Niedziewski et al. 1995). Second component describes reflection of the radiation of hot plasma by cold matter. Approximately half of the primary emission is observed unprocessed and half of it is reflected. Fast variability of the line component observed in a number of sources (Yaqoob et al. 1996) shows that the cold medium is located cospatially with the hot one although some contribution from the reflection by a distant dusty/molecular torus (Krolik et al. 1994, Iwasawa et al. 1997) also might be present.

An attractive scenario of the coexistence of the hot and cold medium is an optically thin corona above an optically thick accretion disc. The existence of such a corona was suggested by Liang & Price (1977).

Main question about coronae refers to the manner how the energy dissipation proceeds in order to heat up the plasma to an (electron) temperature $\sim 10^9$ K. Usually the fraction of energy dissipated in coronae $f$ is a free parameter of a model (Haard & Marschi 1991, 1993, Kusunose & Mineshige 1994, Svensson & Zdziarski 1994). Another approach is to assume that the flux in the corona is generated by viscosity, similarly like in accretion disc. Parameterization by viscosity coefficient $\alpha$ allows to predict the ratio of the energy generated in the disc to that generated in corona. Nevertheless in this models full disc/corona vertical structure is averaged by two hot and cold layers, and the place of transition is either adopted arbitrarily or based on crude approximations (Nakamura & Osaki 1993, Życki et al. 1995, Witt et al. 1997).

It is shown (Krolik, Mecke & Tarter 1981 hereafter KMT) that the illumination of the cold matter ($\sim 10^4$ K) by hard X-rays leads to spontaneous stratification into hot and cold part due to the thermal instability. The two layers are in pressure equilibrium and the discontinuity in the tem-
perature and density reflects the discontinuity in the cooling mechanism: hot layer is cooled by Compton scattering while cold layer is cooled by atomic processes. This phenomenon underlies the spontaneous division of the flow into disc and a corona and therefore it is important to study it in some detail.

The nature of this phenomenon is well understood due to simple analytical studies (KMT 1981, Begelman, McKee & Shields 1983) based on local determination of the temperature. Later studies which included the radiative transfer in the irradiated slab were not suitable for the purpose to analyze the problem since either the instability showed up in these codes as numerical problems (Raymond 1993) or the iteration had to be started below the instability zone (Sincell & Krolik 1997). Therefore we follow the basic semi-analytical approach to the description of the effects of irradiation but we enrich the physics involved in the process.

The vertical structure of such a transition between the Compton cooled and atomic cooled medium in the context of accretion disc/corona system was studied in our previous paper (Różańska & Czerny 1996 hereafter RC96). The model was parameterized by the hard X-ray flux illuminating the transition layer from above and by the soft UV flux (thermal radiation of the cold disc) from below. We took into account the hydrostatic equilibrium and we assumed purely radiative energy exchange between the corona and the disc.

The result of that paper suggested that as soon as there was enough of the hard X-ray emission liberated e.g. in the upper layer of the corona there was a well defined transition layer between the disc and the corona. In this zone two-phase medium may form, i.e. cool clouds can coexist with hot plasma under the constant pressure.

In the present paper we discuss the transition layer considering also the effect of thermal conduction. For the first time this effect was studied in relation to dissipative corona by Maciołek-Niedźwiecki, Krolik & Zdziarski (1997). They studied the influence of the heat flux transport on the structure of the corona and the transition between the hottest layers ($10^8$ K) and Compton heated disc atmosphere ($10^7$ K). Their approach did not allow to determine properly the basis of the corona as they did not include atomic processes and were unable to reach cool disc layers cooled predominantly by lines and bound-free transitions. Our paper is therefore complementary to theirs as we describe more deep parts of the transition starting from Compton heated layers and reaching down the cool disc ($10^4$ K).

We show that the thermal conduction allows to find very sharp but nevertheless continuous solution for density and temperature profile throughout a transition region, without the need for the two-phase medium. However, we study the local stability of such a solution and we show that this solution is not thermally stable which supports the view of the existence of the two-phase medium in the transition zone, as suggested by RC96.

2 ASSUMPTIONS OF THE VERTICAL STRUCTURE

We consider the optically thick and geometrically thin accretion disc illuminated on the top by hard radiation flux. The origin of this hard radiation in radio quiet objects is still unknown and usually two models are under debate. High energy photons can rise through the Comptonization: (1) on thermal electrons in optically thin corona above the disc or in the innermost part of the optically thin disc (2) on non-thermal electrons in the corona or at the basis of the jet-like (albeit not highly collimated) flow.

We do not specify which interpretation is the correct one. Instead, we treat the produced hard X-ray flux as external at a given radius. Therefore, if the hard X-ray emission is produced in the corona it means that we study the boundary layers between the upper, very hot coronal zone and cool disc. On the other hand, if the X-ray emission is generated somewhere in the innermost part of the accretion flow and illuminates the disc at larger radii we actually calculate the entire corona with the transition zone. To simplify the terminology we just call the studied layer ‘the transition zone’ in both cases.

The incident flux of hard radiation $F_{\text{hard}}$ illuminating the upper surface of the transition layer is a free parameter in our calculations. It is only that part of whole radiation flux generated in X-ray source, which is directed towards the disc. We take into account the decrease of this flux by ‘absorption on the spot’ which means full extinction of flux through real absorption and through scattering:

$$F_{\text{hard}}(\tau) = F_{\text{hard}} \exp(-\tau),$$

where $\tau$ is the optical depth measured from the top of the slab downwards to the cold disc. We don’t take into account extinction of flux after one scatter, so the matter on considered optical depth does not become the source of radiation. Such an assumption is valid as long as absorption is the dominant source of opacity. Recent papers on X-ray spectrum reflected from accretion disc suggest that in case of illumination almost 90% of incident luminosity is deposited within the cold disc, while the remaining 10% is reflected (George & Fabian 1991, Matt, Perola & Piro 1991, Haardt & Maraschi 1993). There are strong observational evidences that absorption dominates in the illuminated matter (Magdziarz & Zdziarski 1995). So most of the X-ray energy is converted into soft radiation and reemitted according to the formula:

$$dF_{\text{soft}1}(\tau) = F_{\text{hard}}(1-a)d\tau$$

We assume the albedo of hard radiation from the disc atmosphere to be constant with optical depth and equal $a = 0.15$.

Another source of energy is the gravitational energy of the accreting matter itself liberated via the viscous dissipation (Czerny & King 1989a,b). The soft radiation flux generated in this way increases upward throughout the disc according to the formula:

$$dF_{\text{soft}2}(\tau) = \frac{3}{2} \frac{\Omega \alpha}{\rho \kappa_{\text{tot}}} P d\tau,$$

where $\Omega$ is the angular velocity of keplerian motion, $\alpha$ is the viscosity parameter, as introduced by Shakura & Sunyaev (1973), $P$ means the gas pressure, and $\kappa_{\text{tot}}$ is the total opacity coefficient.

Total soft flux emitted by the disc is a sum of two components described above:

$$dF_{\text{soft}}(\tau) = dF_{\text{soft}1}(\tau) + dF_{\text{soft}2}(\tau)$$

The spectrum of the soft radiation is that of a black body with temperature $T_{\text{soft}}$ given by:

$$T_{\text{soft}} = \frac{1}{\kappa_{\text{tot}}} \left( \frac{c}{4 \pi G} \right)^{1/4} \left( \frac{P}{\rho} \right)^{1/4}$$

where $c$ is the speed of light, $G$ is the gravitational constant, $P$ is the pressure, $\rho$ is the density, and $\kappa_{\text{tot}}$ is the total opacity coefficient, which is given by the sum of the absorption coefficients of the gas and dust.

$$\kappa_{\text{tot}} = \kappa_{\text{gas}} + \kappa_{\text{dust}}$$

The gas opacity is given by:

$$\kappa_{\text{gas}} = \pi \rho \sigma T_{\text{coll}}$$

where $\sigma$ is the Stefan-Boltzmann constant, $T_{\text{coll}}$ is the gas temperature, and $\rho$ is the gas density.

The dust opacity is given by:

$$\kappa_{\text{dust}} = \pi \rho_{\text{dust}} a_{\text{dust}}$$

where $\rho_{\text{dust}}$ is the dust density and $a_{\text{dust}}$ is the dust grain size.

The accretion disc is considered to be optically thick, i.e. the optical depth is large enough so that the radiation from the disc is not completely absorbed. This assumption is supported by the fact that the accretion disc is opaque to X-rays, which indicates that the radiation is strongly absorbed by the disc material.
body with the temperature $T_{bb}$ whilst hard X-ray emission is assumed to be a power law with the energy index $\alpha_E = -0.9$ (here $F_\nu \sim \nu^{\alpha_E}$). We neglect the change of the spectral shape in any of the two components with the optical depth.

Strong illumination by X-rays heats up the outer parts of the disc and the very hot, optically thin slab is created above the disc. Such a corona generally reduces the radiation pressure, so we assume for the transition zone the following equation of state:

$$P = \frac{k}{\mu m_H} \rho T, \quad (5)$$

with the value of molecular weight $\mu = 0.5$ for cosmic chemical composition, i.e. we neglect the radiation pressure throughout the zone.

The vertical distribution of the pressure and the density $\rho$ is defined by the condition of the hydrostatic equilibrium in vertical direction:

$$\frac{1}{\rho} \frac{dP}{dz} = -\Omega^2 (H_d - z) + \frac{F_{soft}}{c} \kappa_{tot} - \frac{F_{hard}}{c} \kappa_{tot}. \quad (6)$$

$H_d$ is the half of the disc thickness and the coordinate $z$ is measured from the disc half-thickness, $H_d$, downwards.

The temperature of a slab is determined by balancing radiative heating and cooling and heat transport via the thermal conduction from the neighbor layers. Time-independent energy balance equation in vertical direction becomes:

$$-\rho \mathcal{L}(\rho, T) = \frac{dq}{dz}, \quad (7)$$

where $\mathcal{L}$ is generalized loss-heat function defined as energy losses minus energy gain via the radiative processes and $q$ is the conductive heat flux.

The solution for vertical disc/corona structure depends on what kind of physical processes are taken into account. In the hot corona Comptonization and bremsstrahlung plays main role, but if we want to find the structure of the transition layer it is obvious that going closer to the cool, dense disc also atomic absorption and emission becomes important. Furthermore, the accreting matter heats itself by viscous forces. Also thermal conduction in vertical direction gives substantial contribution to the heating/cooling since the transition layer is characterized by strong temperature gradient. All physical processes which are included in our energy balance equation are described below in section 3.1.

3 ENERGY BALANCE - EQUILIBRIUM

3.1 Radiative processes

Assuming only radiative energy exchange between the layers, the condition of thermal equilibrium is satisfied by:

$$\mathcal{L}(\rho, T) = 0. \quad (8)$$

Considering following processes: atomic cooling (spontaneous emission, bremsstrahlung), atomic heating (absorption in lines and continuum), Compton heating and cooling and viscous heating, the energy equation becomes:

$$\rho \Lambda(T) - F_{hard} \gamma(T) - \kappa_{es} F_{tot} \frac{4k}{m_c c^2} (T_{1c} - T) -$$

$$\frac{3}{2} \Omega^2 \frac{k}{\mu m_H} \rho T = 0, \quad (9)$$

where $\Lambda(T)$ and $\gamma(T)$ are cooling and heating functions respectively. We use the same $\Lambda(T)$ and $\gamma(T)$ functions determined by CLOUDY code as in our previous paper (RC96) which contains full discussion how to obtain these relations. The first function represents the total cooling rate via the bremsstrahlung, bound-bound and bound-free emission in erg cm$^{-3}$ s$^{-1}$ g$^{-2}$ units, the second one describes the absorption in lines and continuum in the cm$^{-2}$ g$^{-1}$ units.

For radiative processes the energy balance equation is algebraic and we can expect more than one solution for temperature. In fact, we concluded in RC96 that the vertical structure of the transition zone is not continuous, if hydrostatic equilibrium is imposed. For two narrow ranges in optical depth the matter is described by three different values of temperature and one of them correspond to thermally unstable solution. Other two stable solutions cannot be matched continuously. The transition therefore might be discontinuous and located arbitrarily within three-solution zone (as suggested by Koo & Kallman 1994) or, as we argue, the state of gas is determined by a two-phase equilibrium. We concluded there was a narrow slab of matter between disc and corona where cool clouds could coexist with hot plasma under the constant pressure.

To describe the nature of the two-phase medium it is convenient to use the ionization parameter $\Xi$ defined after KMT (1981) as:

$$\Xi \equiv \frac{F_{ion}}{nk_B T_c}, \quad (10)$$

where $F_{ion}/c$ is the pressure of the ionizing radiation, $c$ is the velocity of light, $n$ is the number density of particles [cm$^{-3}$] and $k_B$ - Boltzman’s constant.

For very large $\Xi$ only Compton processes are important and the temperature approaches an asymptotic value (i.e. inverse Compton temperature) independent of $\Xi$. With decreasing ionization parameter bremsstrahlung gradually becomes more important, leading to the temperature decrease. $\Xi^*_c$ is the minimum value of ionization parameter when high temperature equilibrium is possible. For lower $\Xi$ bremsstrahlung cooling overwhelms photoionization heating.

Similar situation is observed for small $\Xi$. The temperature first increases with increasing ionization parameter (because of increasing of ionization level) until ionization parameter reaches critical value $\Xi^*_r$. At this point photoionization overwhelms the line cooling and low temperature equilibrium becomes unstable.

For narrow range of ionization parameter $\Xi^*_c < \Xi < \Xi^*_r$, cold, dense matter can exist in pressure equilibrium with hot, less dense gas. Both branches are connected by a third one with the negative slope which corresponds to thermally unstable solutions.

The ionization state of the irradiated gas in thermal equilibrium is described as a relation $T(\Xi)$ which depends on: ionizing flux and the shape of cooling and heating curves $\Lambda(T)$ and $\gamma(T)$. In our situation hard X-ray radiation is predominantly responsible for the ionization of the gas, and for the eventual presence of the multi-phase region so we accept $F_{ion} = F_{hard}$. 

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3.2 Thermal conduction

Trying to solve disc/corona transition we should expect large temperature gradient so free electrons can efficiently transport the heat from upper layers to the lower ones. Considering such a situation we solve the equation (1) where $\rho c$ contain all radiative processes described above.

Usually the conductive heat transport is treated in two limits, depending on the effective mean free path of electrons in the medium. The classical thermal conductivity is based on the assumption that the mean free path is short in comparison with the temperature scale height $T/|\nabla T|$. For a plasma of cosmic abundance the conductivity is (Draine & Guelianni 1984):

$$\kappa = 5.6 \times 10^{-7} \phi_e T^{5/2},$$

where $\phi_e$ is the factor corresponding to reduction in the mean free path due to magnetic fields and turbulences (it is taken to be $\phi_e = 1$ for equal ion and electron temperature).

The heat flux in classical case is expressed by the diffusion approximation:

$$q_{\text{class}} = -\kappa \frac{dT}{dz}. \quad (12)$$

When the mean free path is comparable to or greater than the temperature scale height the heat flux becomes 'saturated' and nonlocal theory of heat conduction is required. In this situation the heat flux depends on the electron distribution function and is limited by two constraints.

First is that divergence of the current must vanish, and second, that electrons streaming through the ions should be stable against various plasma instabilities (like ion acoustic instability). To estimate of the flux we rewrite after Cowie and McKee (1977):

$$q_{\text{sat}} = 5\phi_e \rho c^3 = 5\phi_e c \rho,$$

where $c^2 = p/\rho$ is the isothermal sound speed, $\phi_e$ is the dimensionless uncertainty parameter. For Maxwellian distribution of electrons $\phi_e = 1.1$, but experimental evidences show that usually $\phi_e \sim 0.3$ (Max et al. 1980) and we will use this value in our calculations.

For the purpose of numerical computations, it is convenient to have the expression which gives smooth transition from classical diffusive to the saturated transport. So we adopt (after Balbus & McKee 1982) very useful formulae defining effective heat flux as:

$$q = -\kappa \frac{dT}{dz} \quad (14)$$

where $\kappa$ is the ratio of the classical to the saturated heat flux at a given temperature:

$$\kappa = \frac{q_{\text{class}}}{q_{\text{sat}}}. \quad (15)$$

The energy balance is the differential equation of second order in this case and we expect the temperature to be the monotonic function of the optical depth. This means the thermal conduction can in principle suppress two-phase medium in the transition zone between disc and corona. The temperature profile $T(\tau)$ could be smooth or very sharp. Smooth one secures stability of the vertical structure solution, but the sharp one should be studied more carefully.

Perhaps small perturbations of thermodynamic parameters can lead to thermal instability and cloud condensation.

4 LOCAL THERMAL INSTABILITY

4.1 General considerations

The thermal stability of the matter interacting with radiation field was discussed by Field (1965). He assumed a perturbation of thermodynamic parameters in uniform medium which in equilibrium state attains $L(\rho_0, T_0) = 0$. He neglected any energy transport and dynamic flows, so his isobaric criterion for thermal instability was the following:

$$\left(\frac{\partial L}{\partial T}\right)_p < 0,$$

where $\left(\partial L/\partial T\right)_p$ is the isobaric evolution of the loss-heat function from the equilibrium.

Balbus(1986) made more general picture since he took into account the dynamical perturbations. In this case the loss-heat function of the medium can be a function of density, temperature and possibly space and time. Thus the isobaric instability criterion becomes:

$$\left(\frac{\partial \ln L}{\partial \ln T}\right)_p < 1. \quad (17)$$

We analyze a local thermal instability taking into account the energy transport. Our medium in this case is characterized by temperature $T_0(z)$ and density $\rho_0(z)$ profile. We neglect any dynamic flows and we assume that $T_0(z)$ is attained by balance between radiative processes and thermal conduction. Our unperturbed medium is in stationary state denoted by us as $U(\rho_0, T_0)$ and it satisfies the condition:

$$\rho_0 U(\rho_0, T_0) = \rho_0 L(\rho_0, T_0) + \frac{d\rho_0}{dz} = 0. \quad (18)$$

When total cooling exceeds total heating (via radiative processes and thermal conduction) then $U > 0$. We perturb two independent thermodynamic variables (density and temperature) under the condition of constant pressure. We are interested only in isobaric perturbations because they can lead to clouds condensation. The specific entropy of the material will change by an amount $\delta S$. This change evolves in time as:

$$\frac{d}{dt} \delta S = \frac{d}{dt} \delta S = \frac{d}{dt} \left( \frac{dQ}{dT} \right) = -\frac{d}{dz} \left( \frac{U}{T} \right). \quad (19)$$

If $\delta S > 0$ then its time derivative should be less than zero to return perturbed entropy to the background one (stability). For $\frac{d}{dz} \delta S$ with positive sign the positive perturbation will grow up from the background entropy (instability). Therefore the instability criterion follows:

$$\left(\frac{\partial (U/T)}{\partial S}\right)_p < 0. \quad (20)$$

Taking into account that in isobaric perturbation $TdS = C_p dT$ and $U = 0$ for stationary state, our criterion reduces to:

$$\left(\frac{\partial U}{\partial T}\right)_p < 0. \quad (21)$$
4.2 The perturbation growth rate

In order to find the possibility of the existence of unstable regions we perform the local thermal stability analysis. It means that we assume our medium to be homogeneous locally. This assumption is done, because our problem is local. We cannot find full disc structure from the bottom to the disc surface for the reason that below our transition layer there is another thermally unstable zone as a result of radiation pressure domination (Pringle, Rees & Pacholczyk 1973, Shakura & Sunyaev 1976). Global analysis does not allow us to separate this two effects.

At each distance $z$ from the disc surface $z$ we are looking for perturbations growing up in time For this purpose we consider the perturbation of two independent parameters: density and temperature, in the form of a wave in space and time:

$$T = T_0 + T_1 \exp(ikz + i\omega t), \quad \rho = \rho_0 + \rho_1 \exp(ikz + i\omega t),$$

where $k$ is a wavenumber and $\omega$ is frequency. All quantities with index 0 characterize unperturbed state.

Field (1965) considered such a perturbation of thermodynamic parameters in homogeneous medium and he derived the critical wavenumber above which the thermal conduction between perturbed element and surroundings suppresses the thermal (radiative) instability:

$$k_c = \left[ -\frac{\rho_0}{\kappa} \left( \frac{\partial \mathcal{L}}{\partial T} \right)_p \right]^{1/2}.$$

The inverse of critical wavenumber gives the characteristic length scale over which perturbation grows up in a result of thermal instability. Let us to denote this length scale of linear perturbation cannot be longer than:

$$\frac{1}{\lambda^*}.$$

A $\lambda^*$ differs from usually used Field length which is defined as:

$$\lambda_F = \left[ \frac{T \kappa}{\rho \mathcal{L}_M} \right]^{1/2},$$

where $\mathcal{L}_M \equiv \max(\text{cooling}, \text{heating}) > 0$.

Physically it means that if perturbed element has a radius $r < \lambda^*$, then such a cloud will always evaporate due to thermal conduction. For perturbation with $r > \lambda^*$ cloud can condensate depending on radiative conditions in the surrounding medium (McKee & Begelman 1990).

For our non-uniform medium with energy transport we still assume only isobaric perturbations $\partial P = 0$ and we neglect dynamical flows (so eulerian and lagrangean perturbations are equal). The matter is characterized by time dependent energy balance equation:

$$\frac{3}{2} \frac{dP}{dt} - \frac{5}{2} \frac{P}{\rho} \frac{d\rho}{dt} + \rho \mathcal{L}(\rho, T) + \frac{d\varrho}{dz} = 0.$$  

For thermal instability analysis we use the heat flux in classical diffusion approximation (equation 25.3).

The first two terms of equation (25) are $\rho T (dS/dt)$ calculated for perfect gas.

Considering only linear deviations from the equilibrium we derive the relation between $\omega$ and $k$ putting perturbations of density and temperature to equations of energy balance and state. The growth rate for each value of $z$ is:

$$\omega = \frac{\kappa_0 k^2 + \rho_0 \left( \frac{\partial \mathcal{L}}{\partial T} \right)_p}{5\rho_0 k_B/m_H}. $$

Here $\omega$ is always real either negative, or positive.

We can write temperature deviation as follows:

$$\delta T = T_1 \exp(\text{Re}(i\omega)t) \exp(ikz).$$

If the real part of $i\omega$ has positive sign then the perturbation will tend to increase, for negative sign perturbation will be suppressed. Therefore we can find the critical wavenumber denoted here by $k_{TC}$ which bounds this two cases. From condition $\text{Re}(i\omega) = 0$, $k_{TC}$ is given by:

$$k_{TC} = \left[ -\frac{\rho_0}{\kappa} \left( \frac{\partial \mathcal{L}}{\partial T} \right)_p \right]^{1/2}.$$  

This relation reduces to the one of Field (1965) in the case of homogeneous medium when $\mathcal{L}_0(\rho_0, T_0) = 0$. For perfect gas we have:

$$\left( \frac{\partial \mathcal{L}}{\partial T} \right)_p = \frac{\mathcal{L}_0}{T_0} - \rho_0 \left( \frac{\partial \mathcal{L}}{\partial \rho} \right)_T.$$  

The characteristic scale length above which the thermal instability develops can be written as $\lambda_{TC} = 1/k_{TC}$. So we expect the perturbed matter with size $r > \lambda_{TC}$ would condense in clouds with different density and temperature than the surrounding medium. We guess that such a perturbation can form a two-phase region in the non-uniform matter.

Local linear stability analysis is valid until $\exp(\text{Re}(i\omega)t)$ is small enough to neglect second order terms. It gives strong constraint on time over which this analysis applies. The time scale of linear perturbation cannot be longer than:

$$t_{max} = \frac{1}{\text{Re}(i\omega)}.$$  

Computation of further evolution would require global stability analysis. This in turn, would be possible only after the unperturbed full vertical structure is computed. Such complete models are not available at present since either surface layers or deep layers are ignored in the course of computing (e.g. Sincell & Krolik 1997, Rózsa et al. 1998). The results of local stability analysis cannot therefore describe the final state of the transition zone. However they may show the existence of the thermal instability zone and strongly indicate the possibility of the formation of the two-phase medium.

5 RESULTS

5.1 Parameters and Boundary conditions

To solve the structure of the transition layer we integrate differential equations presented above as functions of $\tau$ starting from the arbitrary (but low enough) values of the optical depth and density as: $\tau_0 = 2 \times 10^{-5}$, $\rho_0 = 10^{-14} \text{gcm}^{-3}$.

In a pure radiative case the temperature is determined by solving an algebraic equation (41) for each optical depth. In the case with thermal conduction we integrate second order differential equation (44) and the initial temperature on the top is computed from equation (41). Second boundary condition is set as the thermal flux being zero on the surface $(d\mathcal{F}_\mathcal{T}/d\tau = 0)$.

Soft flux generated via viscous dissipation $F_{\text{diss},0}$, $F_{\text{hard},0}$ and X-ray radiation spectrum at the disc surface are free parameters in our model. We consider in detail a case of a
disc emitting soft flux $F_{\text{disc0}} = 5 \times 10^{13} \text{erg cm}^{-2} \text{s}^{-1}$ which corresponds to the black body temperature $3.16 \times 10^{4}$ K. Such a flux is expected from a disc around massive (mass of $10^8 M_\odot$) black hole at $10 R_{\text{sch,hr}}$ if the accretion rate within a disc itself is about $2.8 \times 10^{24}$ g/s (i.e. about 0.01 in Eddington units).

We assumed X-ray radiation spectrum with energy cutoff at 200keV, as usually observed in Seyfert galaxies (Zdziarski et al. 1995) and with spectral indexes: $\alpha_E = -0.987$ for $h\nu < 200$keV and $\alpha_E = -2.1$ for $h\nu > 200$ keV. Results are presented for different values of hard radiation flux: $F_{\text{hard0}} = 5 \times 10^{14}$ and $5 \times 10^{15} \text{erg cm}^{-2} \text{s}^{-1}$.

In the case without thermal conduction the thickness of the accretion disc was adopted as a minimum possible value $H_{\text{dmin}}$ for which the disc is optically thick (RC96). For $H_d < H_{\text{dmin}}$ disc becomes transparent, whereas for $H_d > H_{\text{dmin}}$ disc is no more geometrically thin and hot atmosphere becomes very narrow in comparison with cold matter. Therefore we adopt the disc thickness $H_{\text{dmin}}$ in our approach.

Other quantities adopted in our calculations are typical for standard accretion disc model. We take the viscosity parameter equal $\alpha = 0.1$, the Keplerian angular velocity is $\Omega = 10^{-5} \text{s}^{-1}$ (this value corresponds to Keplerian motion at radius $r = 5.1 \times 10^{14} \text{cm}$ around black hole at the mass of $10^8 M_\odot$), opacity coefficient for electron scattering is $\kappa_\text{es} = 0.34 \text{ cm}^2 \text{g}^{-1}$ and Kramer’s coefficient $\kappa_0 = 3.8 \times 10^{21} \text{ cm}^{-2} \text{ g}^{-1}$. Total opacity coefficient used in relation 3 is computed from the standard relation $\kappa_{\text{tot}} = \kappa_\text{es} + \kappa_0 T^{-3.5}$.

The results are presented as the functions of optical depth $\tau$ or of the distance from the disc surface $z$. Both parameters are connected by relation:

$$dz = \frac{dr}{\kappa_{\text{tot}}(T) \rho}$$

(30)

5.2 The vertical structure of the transition layer

The computations show that in our model heat flux is carried according to classical description, because $\sigma \ll 1$ (see eq. [6]).

Cool, optically thick disc is almost isothermal so, the thermal conduction between the layers is negligible. Similar situation is observed in hot corona, but some heat flux can appear there until corona has really high temperature (classical conductivity is proportional to the $T^{5/3}$). Our numerical results show that heat flux generated in corona is about six orders of magnitude smaller than the one in transition layer.

Fig.1 presents the temperature profiles $T(\tau)$ (I) and stability curves $T(\Xi)$ (II) of the disc illuminated by: a) $F_{\text{hard0}} = 5 \times 10^{14} \text{ erg cm}^{-2} \text{s}^{-1}$ and b) $F_{\text{hard0}} = 5 \times 10^{15} \text{ erg cm}^{-2} \text{s}^{-1}$. Computations for the case when thermal conduction is taken into account (solid line) are compared to the ones without thermal conduction (points). Near the surface, the disc atmosphere for both cases is heated up to the same temperature for given $F_{\text{hard0}}$. The value of this temperature is close to the Inverse Compton temperature which is determined by the spectral shape of incident X-ray radiation. In our case $T_{\text{IC}} = 8.19 \times 10^{8}$ K on the top of the discussed layer. The zone with hot plasma is thicker for larger $F_{\text{hard0}}$. When the density increases with the optical depth the importance of the bremsstrahlung cooling increases and the temperature starts to fall down. For the case with conductivity and for lower $F_{\text{hard0}}$ this point appears at lower optical depth than for pure radiative case.

It is shown in previous paper (Figure 5 of RC96) that our method is applicable for $\tau > 1$ and we do not have to consider the radiative transport equation more carefully. Going further inside the disc temperature reaches the intermediate stable region at $T \sim 8 \times 10^{8}$ K. Since photoionization is the predominant process in the thin transition layer between hot $T \sim 10^8$ K and cold $T \sim 10^4$ K matter, we can expect that heavy elements abundance influences the shape of instability curve. For some configuration of elemental abundance even two intermediate stable regions can exist (Hess et al.1997). In our model one intermediate equilibrium region forms and it is wider for larger values of X-ray radiation flux.

The most interesting feature of the presented profiles is that for the pure radiative case there are two narrow ranges of $\tau$ when matter can reach three different values of temperature and one of them is unstable $dT/d\Xi < 0$. For the case with conductivity the profiles are very sharp but continuous opening a possibility that the thermal conduction suppresses thermal instability in the transition layer between disc and corona.

Finally, the temperature drops to the value of $2.8 \times 10^4$ characteristic for the cool disc. Bigger flux of X-ray radiation illuminating the surface leads to higher optical depth reached by cold matter. Our computations for the case with conductivity are stopped at the point where the thermal conduction flux vanishes at low temperature equilibrium.

Figures named by II (right hand side) show ionization curves and we can easily see that thermal conduction does not change the ionization state of matter in regions of ther-

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Figure 2. Vertical distribution of I) pressure, II) density and III) optical depth at two different values of incident hard radiation flux: left hand side a) $F_{\text{hard,0}} = 5 \times 10^{14}$ erg cm$^{-2}$ s$^{-1}$, right hand side b) $F_{\text{hard,0}} = 5 \times 10^{15}$ erg cm$^{-2}$ s$^{-1}$. Solid wide line represents the case with conductivity, thin dotted line shows the pure radiative case.

5.3 Thermal instability criteria for the case with conductivity

Now we would like to focus on the case with thermal conduction, because it seems to be more realistic. It is obvious that in medium with strong temperature gradient conductive heat transport plays a big role and influences the structure of region. In our model there are two very narrow (first $dT = 1 \times 10^{-4}$ and second $dT = 1.6 \times 10^{-6}$) zones with sharp temperature and density profiles in the transition layer between an accretion disc and a corona. Thermal heat flux rises to $3 \times 10^8$ in these places.

To discuss the stability of such a solution we use local theory developed in Section 4. As was mentioned in above this theory is valid for short period of time comparable with $1/Re(\omega)$. Otherwise we should treat our problem more globally which is more complicated and is not presented in this paper.

Another assumption of applicability of local theory is that perturbation of considered parameter $x$ is smaller than characteristic scale length $x/|\nabla x|$ on which this parameter changes. Therefore we are looking for any perturbations bigger than $\lambda_{TC}$, but smaller than $T/|\nabla T|$, because they can lead matter to condensate into clouds. Treating the problem locally we determine the dependence of $\lambda_{TC}$ and the temperature scale length (named here $H_T$) on the optical depth.

For a wide range of $\tau$ the solution is stable, i.e. $k_{TC}^2 < 0$ (equation (2)). Only for those two narrow zones with extremely strong temperature gradient, wavenumber becomes positive and we can find size of perturbations that lead to thermal instability. Fig. 3. presents variation of $\lambda_{TC}$ and $H_T$ versus $\tau$ for first (I) and second (II) zone, and for two different values of hard radiation flux $F_{\text{hard,0}} = 5 \times 10^{14}$ a) and $F_{\text{hard,0}} = 5 \times 10^{15}$ b). Because possible unstable regions are very thin, we rescale drawings by subtracting the optical depth of the beginning of zones $\tau_0$ from $\tau$. Values of $\tau_0$ are written in the lower left corner for each case. Note that the optical thickness $dr$ of zones does not depend on X-ray radiation flux. Only optical depth $\tau$ and distance of the zones from the disc surface $z$ is bigger for higher $F_{\text{hard,0}}$

Characteristic temperature scale length is always one order of magnitude bigger than $\lambda_{TC}$ for the first zone in case aI and more than half of order of magnitude in case bI. So the perturbations with size $\lambda_{TC}$ can lead to cloud condensation. The half order of magnitude difference between $\lambda_{TC}$ and $H_T$ is observed for the second zone, but not from the beginning where there is no place even for minimal perturbations (aII, bII). The higher hard radiation flux the smaller perturbations are required to produce thermal instability in the transition layer.

We expect our clouds to be under influence of gravity and particularly radiation pressure. The life time of clouds is very short and probably they don’t have enough time to grow up to the much bigger size than the minimal size of perturbations. So the most probable size of clouds is in order of $\lambda_{TC}$ and usually is taken into account in calculations of cloudy regions (Kurpiewski et al. 1997, Krolik 1998). As we can see in unstable zones there is no space for a big clouds. But this is still enough for several clouds in size comparable with $\lambda_{TC}$ and because they are most important the results give good evidence for instability existence in the transition layer.

6 DISCUSSION

In this paper we described the influence of thermal conduction on the vertical structure of the transition layer between a hot corona and a cool accretion disc. We found that heat transport via conductivity does not suppresses thermal instability which occurs in the case when energy exchange
through the disc/corona boundary is only due to radiation processes.

The temperature profile for the case with conductivity is continuous but very sharp particularly in the place where thermal instability was found for pure radiative case. Local stability analysis of stationary state with thermal conduction leads to the conclusion that in zones with extremely sharp temperature gradient, the spontaneous perturbations can cause cloud condensation. The estimated size of possible clouds is \( \sim \lambda \) and is much smaller than the size of thermally unstable zone in pure radiative case which occurs \( \sim 10^{16} \text{ cm} \) for both cases of \( F_{\text{hard0}} \).

Therefore we expect the structure from the corona surface to the disc midplane as follows: first the outermost zone, hot corona cooled by inverse Compton emission (Begelman, McKe & Shields 1983, Begelman & Mckee 1983, Ostricker, Mckee & Klein 1991), second, the transition layer with possibility of the development of the two-phase medium discussed in this paper, and third, main body of the disc (Shakura & Sunnayev 1973). We argue that upper parts of the transition consists of cool clouds at column density \( N_H > 10^{22} \) embedded in hotter medium, but closer to the disc the situation changes continuously in the opposite one when hot blobs at \( N_H \sim 10^{19} \) are deep in cooler matter.

Using local stability analysis we are not able to find the evolution of two-phase zones. We can only find the possibility of clouds existence. We cannot to integrate full vertical disc structure, because it is very difficult to describe all physical processes, particularly in the disc interior (three body recombination). Also on the distance from the Black Hole equal \( 10^5 \text{ cm} \) disc is radiation pressure dominated and another thermal instability plays role. Global analysis does not allow us to separate these two effects. (In the disc interior convection may be important as well.) In addition we found that this is not Sturm-Liouville problem, because we can find only finite number of eigenvalues (this is against main assumption of Sturm-Liouville problem). In the future we should integrate equations numerically without any assumptions about character of perturbations and about differential operator.

These results should be treated as preliminary and further research is necessary first to improve the physical content of the model and second, to explore the range of parameters. In the future work the transfer of hard radiation through the disc matter should be consider more carefully especially in the aim of better determination of absorption and reflection of X-rays depending on optical depth.

The position of possible unstable regions depends on the shape of atomic cooling and heating rates. We tested the accuracy of the adopted functions in our previous paper. Recent studies on photoionization codes show that abundance of iron and oxygen ions (with smaller contribution of magnesium, silicon and sulphur) plays crucial role in the nature of instability. Varying of this abundance strongly influences the heating and cooling behavior and the suppression of thermal instability might be achieved through small changes (Hess et al.1997).

Our results are important in the aim of determine strictly the position of cloudy boundary layer. According to the present results, possible unstable regions in stationary state occur for \( \Xi = 0.28 \) (first zone) and \( \Xi = -0.19 \) (second zone). Improved models should be used in the future to determine the dependence of these values on accretion disc parameters (including radius), X-ray radiation spectrum and chemical composition of the gas. The global stability analysis should be done by solving second order differential equation and finding global modes of thermal instability.

We do not consider dissipative corona above transition layer. The heat transport via conductivity from such hot \( \sim 10^6 \text{ K} \) region can change our picture. We expect to do it in the future.

Unstable zones might be very dynamical and it is difficult to study the time evolution of such a region. Clouds can evaporate or condense or to be accelerated by radiative pressure. The first attempt to compute the spectrum from cloudy transition layer was done by Kurpiewski et al. (1997). They found that broad emission lines observed in quasars spectra can come from clouds which form continuously at the basis of the corona because of thermal instability.

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