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Discs gone with the wind?

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Abstract. We examine the conditions imposed on the vertical structure of accretion discs for them to exist. We find that a wind powered by a thermal instability exists in all discs with certain opacity laws. We find that for discs dominated by bremsstrahlung radiation, an inner hole develops below a certain accretion rate, and that for very low accretion rates, discs may not form at all. The consequences of our results are discussed in relation to black hole soft x-ray transient (BHSXT) accretion discs, and accretion discs in other systems.

Key words: accretion, accretion disks – stars: binaries: close – stars: circumstellar matter – stars: white dwarfs – X-rays: stars

1. Introduction

Accretion discs play an important role in astrophysics. Whenever material with some angular momentum falls towards a central object a disc is generally formed.

The basic philosophy behind the disc accretion model is the following: The matter first forms a disk shaped annular structure around the accreting object at the circularization radii appropriate to the specific angular momentum of the incoming matter. The thickness of the disc H is very small relative to the radius ϖ. The material in the disc rotates at near Keplerian velocities and hence there is shear between two adjacent layers. Viscosity, the exact nature of which is still under discussion, operates in discs, transports angular momentum outwards, and leads to a $small$ flux of matter in the radial direction. The viscosity generates heat which is assumed to be radiated in the perpendicular direction to the surface of the disc (the z direction). So while matter flows slowly inwards in the radial direction, energy dissipated by the inflowing matter is radiated in the perpendicular direction. Clearly, the structure of the disc depends crucially on how efficiently the energy dissipated by viscous heating can escape from the disc. If the vertical structure is such that it cannot cope with the dissipated energy, then it will heat up. If heating does lead to a configuration where the energy loss (cooling) is equal the energy gain (heating), a steady state disc can form. However, if no such equality exists, the matter will expand (or contract), and lead to the formation of a wind (or collapse). In the classical disc problem, the rate at which the disc must carry mass, namely $\dot{m}$, and the specific angular momentum are fixed by the outside. The disc must then adjust to the required structure which is consistent with these constraints, if such a structure can be found.

The disc is basically a 2D structure and all thermodynamic quantities vary with the cylindrical radial coordinate ϖ and z. Obviously, modeling in 2D is much more complicated than in 1D, and hence a frequently encountered approach to modeling the disc behavior and structure is to integrate the disc equations in the vertical z direction and ignore completely the disc energy loss mechanism along with the complicated radiative transfer in the vertical direction. One hopes that 'the disc will converge to the vertical structure needed to carry out the dissipated energy and radiate it out'. The purpose of this paper is to examine if such a state exists, namely: if for any accretion rate and opacity law, is there a corresponding vertical structure?

Viscous energy dissipation removes angular momentum and generates heat. Under certain conditions it is possible that the heat generated in the process cannot be removed by the radiation and as a consequence the gas does not cool and cannot be accreted. The gas then moves away from the z=0 plane and no disc structure forms. Such a possibility can be seen from the following example. Consider a constant temperature disc, namely $T(z) = Const$ but not $T(t)$. Next assume the standard assumption that the heating is proportional to the gas pressure P while the cooling is via free-free emission which is proportional to P^2. Then clearly, for sufficiently low pressures always heating will win and the gas will heat until one of the above assumptions breaks down!

The structure of discs in the vertical direction must allow for the viscous energy released in each volume element. If the disc is in thermal equilibrium, the energy balance dictates an $algebraic$ constraint which the hydrostatic equilibrium equation must satisfy. It is this condition which affects so dramatically the structure of the disc in the vertical direction. Shaviv & Wehrse (1986,SW) integrated the vertical structure equations including the radiation field and encountered numerical problems in convergence. They solved the problem by assuming an artificial cutoff in the viscous energy dissipation and in this way were able to obtain disc models and continuum spectra of classical

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nova disc with high accretion rate. Later Adam et al. (1987) showed how the problem encountered by SW is actually a thermal instability which leads to the formation of corona above the disc.

Czerny & King 1989a,b expanded the idea of SW (1986) and considered optically thin layer above an optically thick disc. They confirmed the results of SW and showed that viscous dissipation in the outer atmospheres of unilluminated accretion disc produces hot corona and thermally driven winds. Czerny & King 1989a also obtained a critical accretion rate above which a full hydrostatic equilibrium does not exist. The limiting accretion rate corresponds to $2.6 \times 10^{-3} \alpha \dot{m}_{Edd}$ where $\dot{m}_{Edd}$ is the critical accretion rate which leads to the Eddington luminosity and α is the viscosity constant. Here we treat the optically thick regime and show that the flow starts already in the optically thick region.

Dumont et al. (1991) (here after DCKL) discussed the structure of optically thin accretion disc. In particular, they investigated in detail the effects NLTE has on the vertically integrated structure equations. DCKL obtained the energy balance of the disc from equating the cooling to the required energy dissipation for a given accretion rate. The basic equations used by DCKL are the vertical height averaged equations, namely they assumed a height H given by:

$$\frac{H}{R} = c_s \left(\frac{R^3}{GM}\right)^{1/2}, \tag{1}$$

where c_s is the speed of sound and a surface density which is given by

$$\nu\Sigma = \frac{\dot{m}}{3\pi}\left[1 - \left(\frac{R}{R_*}\right)^{1/2}\right], \tag{2}$$

where $\dot{m}$ is the total accretion rate and the kinematic viscosity ν is given by

$$\nu = \alpha c_s H. \tag{3}$$

Thus, the first difference between the present treatment and that of DCKL is that DCKL considered averaged equations while here we use the local equation. The second difference is in the energy balance. Given the cooling expressed in terms of the vertical averaged equations, DCKL equate the averaged cooling to the required energy loss for a given accretion rate and thus obtain that:

$$n^2\Gamma = \frac{3}{16\pi}\frac{GM\dot{m}}{HR^3}\left[1 - \left(\frac{R}{R_*}\right)^{1/2}\right], \tag{4}$$

where $n = \Sigma/(2Hm_H$ is the vertically averaged number density of hydrogen nuclei and $n^2\Gamma$ the vertically averaged cooling rate. The advantage of this approach is that in this way the details of the dissipation mechanism and in particular the assumption expressed in Eq. 3 disappears from the equations. In short, DCKL investigated the radial structure while we here explore the vertical structure. The present work is a further investigation of the instability in the thermal balance of the vertical structure of the disc for low $\dot{m}$ systems. We show that it determines the fate of the entire disc structure.

2. The vertical structure

We assume that

– The disc is geometrically thin $H/\varpi << 1$ and all gradients in the vertical direction are much greater than the corresponding ones in the radial direction.
– The radiative transfer can be treated in the grey approximation and only the vertical directions are considered. ('two stream approximation')
– LTE prevails throughout the entire disc. Deviations from LTE will be discussed in a subsequent paper. We believe that LTE has an effect on the details but not on the general picture discussed here.

The equation for hydrostatic equilibrium in the vertical direction is

$$\frac{dP}{dz} = -g(z)\rho + \chi F/c, \tag{5}$$

where P, T, ρ and μ are the gas pressure, temperature, density and molecular weight respectively and z the vertical coordinate, $g(z) \approx g_0 z$ is the gravitational component in the z direction. The extinction coefficient is given by $\chi = \kappa + \sigma$ where κ and σ are the volume absorption and scattering coefficients respectively. The total radiative flux is F.

The condition of thermal equilibrium is:

$$\frac{dF}{dz} = \epsilon_{vis}, \tag{6}$$

where ϵ_{vis} is the viscous energy dissipation per unit volume.

We assume here that the entire energy dissipated by the viscous forces is carried away by radiation i.e. we ignore at this moment convection and any other non radiative energy transport mechanisms. The radiative transfer equation is:

$$\frac{dF}{dz} = \kappa(B - J), \tag{7}$$

and

$$\frac{dJ}{dz} = \chi F, \tag{8}$$

where B is the wavelength integrated Planck function and J is the wavelength integrated mean intensity. The thermal equilibrium condition can also be written as

$$\kappa(B - J) = \epsilon_{vis} = qP, \tag{9}$$

where $q = (3/2)\alpha\Omega_K$ in the notation of SW and we use the α model approximation for the viscous energy dissipation.

The boundary conditions are:

$$F(0) = 0, J(z_0) = 2F(z_0), P(z_0) = P_0.$$

The above equations, together with the boundary conditions form a closed set, and can be solved for the disc structure for given $F(z_0)$ and $P(z_0)$.

3. The thermal equilibrium condition

Consider the thermal equilibrium equation (9) at a point z_0 above the disc where the optical depth from infinity to z_0 is sufficiently small. Note the appearance of J in the equation. If

$$\int_{z_0}^{\infty} \kappa dz << 1 \tag{10}$$

and there is no significant positive source function gradient, J may be neglected in condition (9). In general, the value of J depends on the structure of the disc and the radiation inside it. We may write that $J = \sigma T_j^4$ where T_j is related to the effective temperature in the disc, which is given by the local mass transfer rate in the disc $\dot{m}(\varpi)$ and σ is the Stefan Boltzmann constant.

In this paper, we consider the important regime of low $\dot{m}$ and restrict the discussion mainly to regimes where absorption is due to hydrogen free-free. This is regime (b) in the SS models where gas pressure dominates over radiation pressure, and is expected to encompass most of the inner regions of the discs in compact star binaries and in AGN in the optically thin, namely the low $\dot{m}$ regime.

We approximate the absorption coefficient by a power law

$$\rho\kappa = AP^m T^n. \tag{11}$$

In the present numerical calculations we use free free opacity i.e. $m = 2$ and $n = -11/2$. The inclusion of more accurate expressions for the free free opacity as well as other sources of opacity and consideration of Compton cooling, is left for future investigation. We note however, that Compton cooling is not expected to be a significant coolant in the body of the disc in the regime that we presently consider. Here, we contend ourselves with an approximation that serves to demonstrate the fundamental physical effect that we wish to highlight in this paper.

For the adopted power law opacity, the thermal equilibrium condition (9) leads to

$$AP^2 T^{-11/2}\sigma(T^4 - T_j^4) = \frac{3}{2}\alpha\sqrt{\frac{GM}{R^3}}P. \tag{12}$$

This relation is shown in Fig. 1 for two cases. An accretion disc around a black hole of $10^8 M_\odot$ at $R = 10^{13}$ cm (10 Schwarzschild radii) for accretion rates of 10^{-2},10^{-4},10^{-6},10^{-8}, 10^{-10} $M_\odot y^{-1}$ and around a black hole soft x-ray transient (BHSXT of mass $10M_\odot$ at 10^9 cm ($\approx$ 300 Schwarzschild radii) for 10^{-7},10^{-9},10^{-11},$10^{-13}$and 10^{-15} $M_\odot y^{-1}$.

Along the sharply declining vertical portions of these curves, the mean intensity J plays an important role in determining the thermal structure, whereas the rising parts of the curves are dominated by optically thin free-free cooling. We note that for every J or equivalently $\dot{m}$, there is a minimal pressure below which there is no solution, and above which there are two solutions for a given pressure. In fact, such a minimum exists for all opacity laws with $m > 1$ and $n < -4$. A mass element which finds itself on the left hand part of these curves could be in thermal and hydrostatic equilibrium, whereas if it is on the rising part

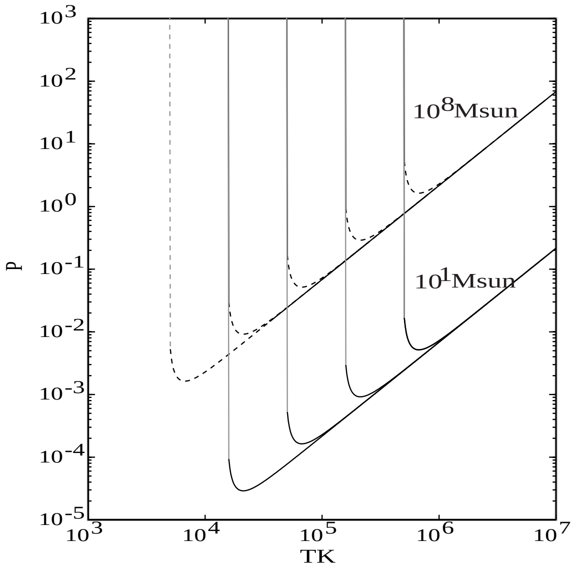

Fig. 1. The pressure – temperature relationship corresponding to thermal equilibrium in the upper atmosphere of a disc around a $10M_\odot$ black hole soft x-ray transient at 10^9 cm and a $10^8 M_\odot$ AGN disc at 10^{13} cm. The different curves correspond to mass transfer rates of 10^{-2}, 10^{-4}, 10^{-6}, 10^{-8}, 10^{-10} $M_\odot\, y^{-1}$ for the AGN disc and 10^{-7}, 10^{-9}, 10^{-11}, 10^{-13} and 10^{-15} $M_\odot\, y^{-1}$ for the BHSXT disc.

(on the right of the minimum), no hydrostatic state is possible with pressure rising outward. We note that for the adopted absorption the structure line of a classical disc should have mass elements which lie along the thermally stable sharply declining parts of these curves. We have used the above results to represent schematically an example of a possible disc structure in thermal and hydrostatic equilibrium. The example is shown in Fig. 2, where α_1 corresponds to the symmetry plane while β_1 corresponds to the surface of the disc, $z = z_0$ where eventually the vertical structure is truncated.

For a physical model, the point z_0 must be close to where the pressure tends to zero (in the sense that the matter above z_0 is sufficiently optically thin and produces only a negligible luminosity), or more realistically at some prescribed low value (say $P = 10^{-5} dyn/cm^2$) of the order of the environment pressure in the ISM (Wehrse, Wickramasinghe & Shaviv 1998). The shape of the cooling curve, however, precludes such a solution in general, as can be seen from Fig. 1. In this case the structure is better described in Fig. 2 by point β_2 which does not lie on any thermal equilibrium line but is in a region where the heating dominates over the cooling. The point β_2 cannot lie on the rising part of the curve since this configuration is thermally unstable. So far we have seen that there is a critical pressure $P_c = P_{min}$, and corresponding to this pressure we can define a critical temperature $T_c = T_{min}$. The quantities P_c and T_c will play an important role in the subsequent discussion. It is obvious that the temperature is always higher than T_j!

The minimum values of pressure and temperature for mass transfer rates of and for the AGN and black hole transient cases are shown in Fig. 4.

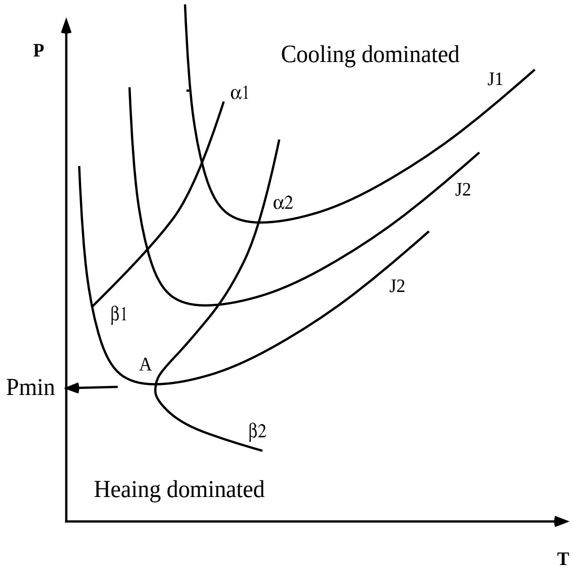

Fig. 2. The P-T plane, the thermal equilibrium lines $J(n = 1, 2 \ldots)$ and the structure lines in two cases. At each point in the disc there is a corresponding P-T curve with the appropriate J. The thermal equilibrium lines correspond to $J1 > J2 > J3$, where $J = \sigma T_j^4$. The heavy line marked $\alpha_1 - \beta_1$ denotes the structure line of a disc for which all points lie on a corresponding P-T equilibrium line. α_1 is the $z = 0$ plane and β_1 the outermost point. In the second structure line $\alpha_2 - \beta_2$, which is more realistic, point β_2 does not lie any longer on an energy equilibrium line.

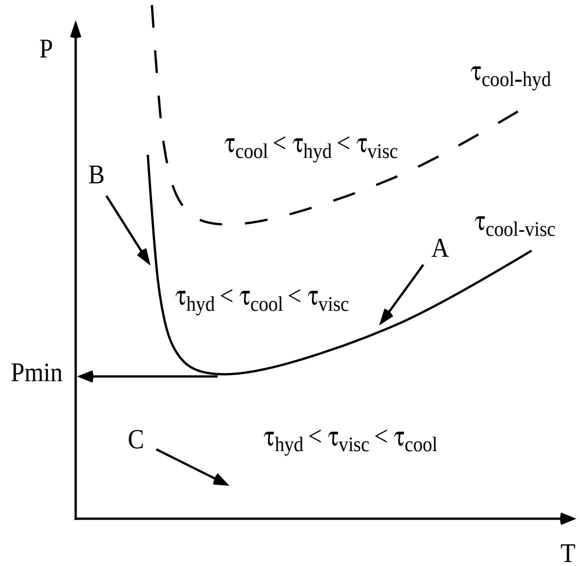

Fig. 3. The relevant time scales in the $P - T$ diagram. Since $\alpha \leq 1$ the curve $\tau_{cool-hyd}$ is above the curve $\tau_{cool-visc}$

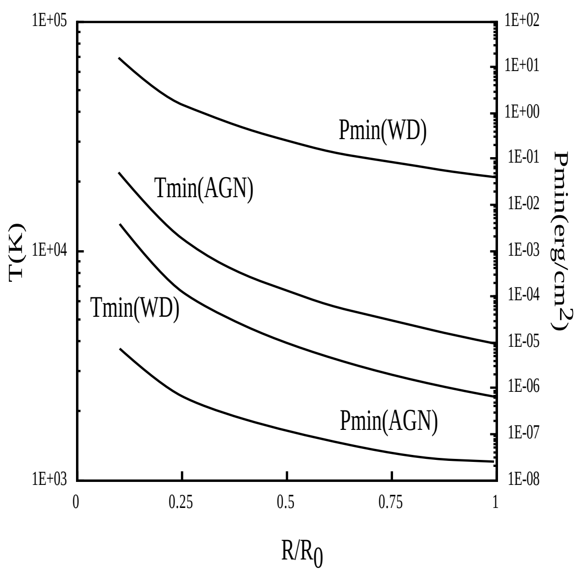

Fig. 4. The values of P_{min} and T_{min} as a function of the radial coordinate for a typical accreting WD with $1M_\odot$ and a BH with $10^9 M_\odot$. The units of the radial distance are $\varpi_{0,WD} = 10^{10} cm$ and $\varpi_{0,BH} = 10^{17} cm$.

The next question is what price the disc pays for not obeying the thermal equilibrium relationship. Said differently, the part $A - \beta_2$ is clearly *not*also in hydrostatic equilibrium. So what happens to it?

4. Understanding the P-T plane

To understand the behavior of the solution in the P-T plane described in Figs. (1, 3) we define the following time scales: The viscous heating time scale τ_{heat} defined by:

$$\tau_{heat} = \frac{\rho c_s^2}{\Gamma}, \tag{13}$$

where c_s is the adiabatic speed of sound and $\Gamma = \nu \rho v_\phi^2 \varpi^2$ is the local viscous heating per unit volume. The kinematic viscosity is given by $\nu = \alpha c_s H$. The cooling time is given by:

$$\tau_{cool} = \frac{\rho c_s^2}{\kappa(B - J)}. \tag{14}$$

Finally the hydrodynamic time scale is

$$\tau_{hyd} = \frac{H}{c_s}. \tag{15}$$

It is easy to see that

$$\tau_{visc} = \frac{1}{\alpha} \tau_{hyd} \tag{16}$$

and

$$\tau_{hyd} = \frac{\varpi^{3/2}}{\sqrt{GM}}. \tag{17}$$

Thus τ_{vis} and τ_{hyd} are constants in this approximation.

Let us draw now the following curves in the P-T plane:

$$\tau_{heat-cool}: \quad \tau_{heat} = \tau_{cool}, \tag{18}$$

$$\tau_{heat-hyd}: \quad \tau_{heat} = \tau_{hyd}, \tag{19}$$

and

$$\tau_{cool-hyd}: \quad \tau_{cool} = \tau_{hyd}. \tag{20}$$

It is instructive to examine the behavior of a mass element in the P-T plane from the point of view of the ratio between the various time scales. For this reason we redraw Fig. 1 and incorporate in it the relevant time scales. The curves marked with Jn are the thermal balance condition for the relevant J.

Consider a mass element at point A in Fig. 3. Since the hydrodynamic time is the shortest, the element wants to expand. However, the cooling time is shorter than the viscous heating time, and hence the element will move to lower pressures and lower temperatures. The motion will continue until the element will reach the thermal equilibrium line (lines marked by Jn in Fig. 2) appropriate for its conditions. The movement of an element at point B is in the opposite direction, and it will reach the thermal equilibrium line from the other side. This behavior can exist as long as $P(A) > P_{min}$ and $P(B) > P_{min}$. However, viscous heating of an element at C will cause it to heat and expand in such a way that it will never reach a thermal equilibrium line. We see also that the equilibrium line to the right of the minimal pressure is also unstable and for the same reason. An element below it does not approach the line, but moves away from it.

At this point we can return and examine the disc structure which penetrates into the region of point C. The part of the structure from point A (cf. Fig. 3) to point β_2 cannot be in hydrostatic equilibrium and will heat up and expand. Thus, this part of the disc represents a slowly expanding wind moving away from the $z = 0$ plane. Whenever the structure point representing the pressure decreases below P_{min} (on the thermal equilibrium line corresponding to the proper J) an expanding wind appears. *The expanding wind is an integral part of any disc with a proper opacity law.* The degree of mass loss through the wind depends how high is the point α_2 above P_{min}. For very optically thick discs, we expect $\dot{m}_{wind}$ to be very much less than the total external mass input rate into the disc.

The minimal pressure P_{min} is a function of ϖ and $\dot{m}$. Since *all* discs must lose mass via a wind at some level, $\dot{m}$ cannot be constant along ϖ. So $\dot{m} = \dot{m}(\varpi)$. The total mass loss rate from the disc in the form of a wind is $\dot{m}_{wind} = \dot{m}(\varpi_{out}) - \dot{m}(\varpi_{obj})$ where ϖ_{out} is the outer radius of the disc and ϖ_{obj} is the inner one, or to a good approximation the radius of the accreting object. For sufficiently high $\dot{m}(\varpi_{out})$ we have

$$\dot{m}(\varpi_{out}) - \dot{m}(\varpi_{obj}) << \dot{m}(\varpi_{out}). \tag{21}$$

The disc models calculated by SW were calculated in this limit.

When $\dot{m}(\varpi_{out})$ decreases the above condition is not satisfied any more and the change of $\dot{m}$ throughout the disc must be taken into account. At a certain accretion rate which we denote by $\dot{m}_h$, $\dot{m}(\varpi)$ will vanish for some $\varpi_h > \varpi_{obj}$. In this case the disc structure disappears before reaching the central object and the disc has a hole inside.

We can use our calculations to roughly estimate the critical mass transfer rate below which we expect a hole to develop at some radius ϖ_h as follows: We assume for this estimate an isothermal atmosphere and the free-free opacity law $\kappa = 4.0 \times 10^8 P^2 T^{-11/2}$ (Allen 1976). For a thermally driven wind, the specific mass loss rate $\dot{m}^{sp}_{wind}$ will be given approximately by the product of the density ρ_c and the adiabatic sound speed c_s (cf. Czerny & King, 1989a,b) evaluated at the structure point where hydrostatic equilibrium breaks down for the first time, namely at P_c and T_c. Hence:

$$\dot{m}^{sp}_{wind} = \rho_c c_s = \rho_c \sqrt{\frac{R_g T_c}{\mu}} = P_c / \sqrt{\frac{R_g T_c}{\mu}}. \tag{22}$$

P_c and T_c can be obtained from a complete solution to our equations, or approximately as follows. Using the condition Eq. 12, we find that the minimum in the pressure-temperature relationship occurs at

$$T_c^4 = \frac{n}{n+4} T_j^4, \tag{23}$$

or, for sufficiently high point in the disc, where T_j is to a good approximation equal to T_e. Hence we can write that:

$$T_c^4 = \frac{n}{n+4} \frac{3GM\dot{m}}{8\pi\sigma\varpi^3} f, \tag{24}$$

where f depends on the nature of the inner boundary condition. Assuming $f \approx 1$, which is valid for a point far from the inner boundary

$$T_c = 7.1 \times 10^6 (\frac{n}{n+4})^{1/4} \dot{m}_{16}^{1/4} (\frac{10M/M_\odot}{\varpi_9^3})^{1/4}. \tag{25}$$

The pressure at the critical point is given by

$$\frac{P_c^{m-1}}{T_c^{-n-4}} = 5.2 \times 10^{-5} (\frac{\alpha}{0.1})(\frac{-3n}{8})(\frac{4 \times 10^8}{A}) (\frac{10M/M_\odot}{\varpi_9^3})^{1/2}. \tag{26}$$

For the particular case of free-free opacity we obtain the following expressions.

$$T_c = 8.5 \times 10^4 \dot{m}_{16}^{1/4} (\frac{10M/M_\odot}{\varpi_9^3})^{1/4} \tag{27}$$

$$\frac{P_c}{T_c^{3/2}} = 1.07 \times 10^{-5} (\frac{\alpha}{0.1})(\frac{4 \times 10^8}{A})(\frac{10M/M_\odot}{\varpi_9^3})^{1/2} \tag{28}$$

and hence

$$\dot{m}^{sp}_{wind} = 1.0 \times 10^{-4} (\frac{\alpha}{0.1})(\frac{4 \times 10^8}{A}) (\frac{10M/M_\odot}{\varpi_9^3})^{3/4} \dot{m}_{16}^{1/4}. \tag{29}$$

The total wind mass loss rate outside a given radius ϖ is obtained by integrating the above expression, and is given approximately by

$$\dot{m}_{wind} = 5 \times 10^{15} (\frac{\alpha}{0.1})(\frac{4 \times 10^8}{A}) (10M/M_\odot)^{3/4} (\varpi_9)^{-1/4} \dot{m}_{16}^{1/4} g/s, \tag{30}$$

where we have assumed that the outer radius of the disc is much larger than the radius at which the inner boundary conditions should be imposed. (Note that $f < 1$ so the above estimate is an upper limit for $\dot{m}_{wind}$.)

The specific mass loss rate (and of course also the integrated mass loss rate) increases as the inner radius decreases. For certain conditions it is possible that the wind mass loss rate will equal the mass transfer rate into the disc before the stellar surface is reached. Under these circumstances, a hole may develop in the inner regions of the disc. The radius of this hole can be obtained from the condition $\dot{m} = \dot{m}_{wind}$, and is given by

$$\varpi_{h,9} = 6.0 \times 10^{-2} (\frac{10M/M_{\odot}}{\dot{m}_{16}})^3 (\frac{\alpha}{0.1})^4 (\frac{4 \times 10^8}{A})^4. \quad (31)$$

The above estimates indicate that for a typical BHSXT of mass $10M_{\odot}$, a hole of radius of approximately 10^9 cm will develop for a mass transfer rate of $10^{15} gs^{-1}$. We note the strong dependence of the radius of the hole on the uncertain viscosity parameter α. Nevertheless it is interesting to note that the numerical value given above is close to the value deduced for the hole in the disc in BHSXT GRO1655-40 (Hameury et al. 1997).

We can make similar estimates for the white dwarf and the AGN black hole cases, but here more careful attention will need to be paid to the nature and details of the opacity. The estimates based on free-free opacity are likely to be appropriate only for restricted regions of CV and AGN discs.

The above results imply that in general, for very low $\dot{m}$, there is a regime where no static solutions will exist, and the matter will expand as soon as it tries to lose angular momentum. Matter that tries to accrete at extremely low accretion rates is blown away from the central object by the very process that is supposed to remove the angular momentum.

The extreme sensitivity of the radius of the hole to the value of α may in principle be used to infer a good estimate of α when such a hole and it radius are observed.

How far will the material go? From observation of Fig. 2 one can conclude that the material will expand in the P-T plane for ever. However, the matter loses angular momentum via the dissipation that heats it. At the same time, as z increases the effective gravitational acceleration increases for a while and then decreases. It is not clear a priori which one wins. If the loss of angular momentum is the larger, then as the matter move up it moves inward. When the matter will reach radial motion there will be no more viscous heating and the thermal instability will die away. On the other hand, if the loss of angular momentum is minimal the matter will move outward. It is instructive to consider the energy balance. The gravitational energy released by the matter as it moves a distance ΔR is:

$$E_{grav} \Delta R = \frac{GM}{2R^2} \Delta R \quad (32)$$

while the energy dissipated in this ring is

$$D(R) = \frac{3GM}{8\pi R^3} \left[1 - \frac{R_*}{R} \right] 2\Pi \Delta R. \quad (33)$$

A necessary but not sufficient condition for a wind is that the dissipation inside the ring $R - R + \Delta R$ be greater than the gravitational energy of the matter at this point, in other words:

$$(1\text{–}3 \frac{R_*}{R}^{1/2}) \geq 0, \quad (34)$$

a condition which is satisfied only for $R > 9R_*$. Hence, for sufficiently large radii, there is a possibility for a strict mass loss from the system via a wind. The extra energy needed to lift the matter to infinity comes from the dissipation energy at lower radii which was convected outward. But at lower radii there is apparently not enough energy and the consequence is a hot corona and eventually radial inflow into the accretor.

Our estimates of the wind expansion ignores many effects like radiation pressure, radiation pressure in lines etc. It seems therefore that the present estimate is a lower limit to the actual phenomenon.

5. Additional consequences and conclusions

Our discussion here was carried out under the assumption that the radial advection of energy is negligible and the material rotates with Keplerian velocities. Could it be that as the thermal balance ceases to exist the disc will choose to develop a flow in the radial direction leading to a non negligible radial advection term rather than develop a wind? Our analysis indicates that this is impossible. Actually, advection *aggravates* the instability that we have discussed because when advection is assumed to occur, the temperature in the disc is higher and approaches the virial temperature (Narayan & Yi 1994; Shaviv et al. 1998). When discs with an advection term and without it are compared, one finds that the disc with advection is much more bloated in the z direction violating the 1D approximation (Wehrse et al. 1998) and therefore more prone to developing outflows in the z direction. The claim that radial advection stabilizes the thermal instability and produces another stable branch (Narayan & Yi 1994), is based on a 1-D analysis which ignores the thermal instability and its implications for the vertical direction.

It is possible that in the binary context, for certain ranges in the mass transfer rate from the companion, one may encounter a situation where the disc structure oscillates between two different states, one with a hole in the centre, and another with a more complete disc structure extending to the central object. Such structures appear to be common in the CV and the black hole x-ray transient contexts. We consider a situation where the mass transfer rate from the companion is in the range which cause the outer regions of the disc to cycle between two stable branches, driving the disc into high and low mass transfer states, as in the disc instability model (for a review cf. Cannizzo 1996). In the low state (quiescence) the mass transfer rate through the central regions of the disc is low, and the disc may develop a hole in the central regions for the reasons described previously. Accretion from a hot coronal wind will continue during this phase. As mass transfer from the companion continues, the surface density in the outer regions will build up to the critical surface density which triggers thermal instability, and causes the outer regions of the disc to transform to the second stable branch. A heat front will propagate through the disc, which now has a hole in the centre, and transform the entire disc into the hot high viscosity (outburst) state. During the hot phase, the viscosity is higher, and the mass transfer rate through the disc increases. The hole will then fill up and the disc will extend all the way to the star. X-rays from the inner disc will change the boundary condition on the surface of the disc, and may result in enhanced coronal winds from the disc during this phase.

We have argued that considerations of vertical structure lead one to conclude that a vertical outflow is an avoidable consequence. For high $\dot{m}$ discs, thermally driven winds will play a minor role in the energetics and structure of the disc, although additional effects, such as radiation pressure driven outflows may come into play in certain regimes. However, as $\dot{m}$ decreases, thermal winds may begin to play the dominant role, and below a certain mass transfer rate, a hole may develop in the centers of discs. In extreme cases, the bulk of the disc may be lost in the wind. We have shown that this could occur if free-free opacity dominates over a significant portion of the disc, as may be the case for black hole soft x-ray transients. Thermal winds will also be increasingly important in low $\dot{m}$ CV and AGN black hole discs where other opacities dominate, but more detailed calculations are required to estimate the critical mass transfer rate below which holes develop.

The following are likely consequences of our model:

1. Thermally driven winds are a natural consequence of the shearing motion that occurs in discs, and disc viscosity, and should be common place in astrophysical systems which exhibit discs.
2. Due to the prevalence of wind mass loss from discs, the disc temperature profile is expected to be flatter than predicted by the Shakura & Sunyaev (1973) thin disc models, particularly in systems with low $\dot{m}$.
3. Low $\dot{m}$ discs should exhibit centrally evacuated regions and lack the high temperature thermal component usually attributed to the central regions of discs.
4. Direct accretion via a disc is not expected to occur in low $\dot{m}$ systems, so that a classical boundary layer around the equator will not form. The accretion in classical nova is apparently spherically symmetric.
5. The wind is expected to heat to temperatures of the order of at most 0.4 of the virial temperatures during its expansion phase, and be detected at x-ray energies.

The above characteristics may explain the following observations:

1. The overwhelming evidence for holes in the centers of accretion discs around white dwarfs in dwarf novae during quiescence (Warner 1995), and in the centers of black hole x-ray transients such as GRO1655-40 during quiescence, from the measured UV and Soft x-ray delays (Hameury et al. 1997).
2. The evidence for a hard x-ray spectrum without an accompanying soft x-ray black body component in the spectra of some black hole transients during quiescence (Tanaka & Lewin 1995).
3. The evidence for ‘discs’ which are underluminous in comparison to deduced mass input rates, around black holes in binaries (Narayan et al. 1996), and in the centers of some AGN and our galactic centre (Narayan et al. 1995).

Finally, we note that other models have been developed to explain central holes in CV discs. We note in particular the coronal siphon model of Meyer & Meyer-Hoffmeister (1994). These authors assume the existence of a corona and calculate the evaporation of the disc by means of an electron conduction flux from the corona back to the disc. In the present case the thermal instability generates the expanding hot layers which eventually become an expanding corona. We expect that thermal winds will be more effective in creating central holes in cases where viscous dissipation ($\propto \Omega_K^2$) plays a relatively more important role in determining the vertical structure of discs.

6. Critique of the basic assumptions

One of the fundamental assumptions in the present analysis is the assumption of LTE. However, this assumption may not be valid for optically thin discs. A crucial feature of the above analysis and consequences is the fact that the cooling fails to overcome the heating. From thermodynamics we can generally claim that NLTE cooling will be less powerful than LTE cooling, as indeed found by DCKL that the LTE overestimates the actual cooling. Thus, we expect that NLTE will *aggravate* the instability and enhance the heating.

A similar argument holds for very cool discs when the actual opacity decreases. Again we encounter a situation where the cooling decreases while the heating prevails.

We assumed the existence of a thermal balance. Clearly, as soon as the balance is violated the gas starts to expand and the momentum equation in the vertical direction contains two additional terms, namely $v_r \partial v_z/\partial r$ and $v_z \partial v_z/\partial z$. Strictly speaking the 1D approximation is not valid anymore. Thus the entire considerations developed here must be reconfirmed using 2D approach.

The actual calculations and figures were carried out for free-free opacity and naturally the question arise what happens under other conditions. The situation does not change in an essential way when the opacity law changes (at least as long as $m \neq 1$). The reason is the appearance of J in the energy balance. If we write $J = \sigma T_{eff}^4$ then obviously, as the temperature decreases towards T_{eff} the pressure rises almost irrespective of n leading to a situation where T increases as P decreases. The fact that P does not increases, namely that a given P does not have two solutions for T implies that P continues to decrease while T increases preserving the hydrostatic equilibrium. In this case there is always hydrostatic equilibrium though the temperature rises as the pressure decreases. Hence no breakdown of the hydrostatic solution and no wind. However, here we recall the model of Meyer & Meyer-Hoffmeister (1994) in which electron conduction from a hot corona into the photosphere leads to evaporation of the disc. In a more general case in which the opacity changes along the radial distance, there can be a wind in one zone and no wind in the others.

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