ACCRETION DISKS IN SYMBIOTIC STARS.

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ABSTRACT.

We give an overview over the theory of geometrically thin α-accretion disks; further we introduce the two different proposed mechanisms that can cause outbursts of accretion disks; and finally we compare the results of these models applied to symbiotic stars (=SS).

1. INTRODUCTION.

Accretion disks were *invented* almost a quarter of a millennium ago: In 1755 Immanuel Kant explained the origin of the solar system by a model that we today would call an accretion disk model. The main principles behind accretion disks, namely the inward transport of matter with outward transport of angular momentum, were formulated by von Weizsäcker (1943) and, especially, Lüst (1952). It took another two decades until accretion disk physics made the next major step forward: in 1973 Shakura and Sunyaev introduced the so-called α-accretion disks. We shall follow the line of this model and its approximations in presenting the basic physics of accretion disks in Sect. 2.

While this ansatz was successful in describing the overall behaviour of accretion disks, especially in dwarf nova systems, there remained the problem of the outbursts or active phases of these systems. As early as 1972 and 1974, Bath, and Osaki, resp., proposed two different models. While in Bath's description the outburst takes place in the envelope of the companion star (an instability there leads to a strong increase in the mass transfer rate; what we observe is this material being processed through the disk), in Osaki's picture an essentially constant mass transfer feeds an unstable accretion disk (matter first stored in the outer regions of the disk, is released by a – then unknown – instability and processed through the disk). Meyer and Meyer-Hofmeister (1981) found an instability with the properties required by Osaki's model. In Sect. 3 the different possible outburst mechanisms are discussed.

The application of accretion disks to and their outburst behaviour in SS was analyzed by several authors (Bath, 1977; Bath and Pringle, 1982; Plavec, 1982; Kenyon and Webbink, 1984; Kenyon, 1985, 1986; Duschl, 1985a, 1986a, 1986b; among others). We are discussing SS with a characteristic behaviour like CI Cyg, Z And, or AR Pav. In Sect. 4 we shall describe the present situation and compare the different models for SS. An independent indication that accretion disks exist in SS comes from the observations of for instance bipolar outflow (Solf, 1984).

2. ACCRETION DISKS ...

In this Section we give a concise overview over accretion disk physics; for a more detailed treat-

J. Mikolajewska et al. (eds.), The Symbiotic Phenomenon, 137–148.

https://doi.org/10.1017/S025292110010329X Published online by Cambridge University Press

ment, and for original references we refer to Frank *et al.* (1985).

We treat the accretion disk in a cylindrical coordinate system $\{s, z, \varphi\}$, and introduce a radial distance $r = \sqrt{s^2 + z^2}$. The structure and evolution is described by the continuity equation, the Navier-Stokes-equations, and equations for the energy transport. In the following we first present these basic equations and introduce the relevant quantities (2.1.); then we discuss the commonly used approximations (*thin α-disk theory*) (2.2.); in the next paragraph (2.3.) we shall describe the set of resulting equations that are actually solved; and finally the results for stationary disks are presented (2.4.). If not stated otherwise, we use cgs-units.

2.1. Basic equations.

2.1.1. The continuity equation.

$$\frac{d\rho}{dt} + \rho \cdot \underline{\nabla}\underline{v} = 0 , \tag{1}$$

where ρ is the matter density, t the time, and $\underline{v}$ the velocity vector.

2.1.2. Navier-Stokes-equation.

$$\frac{d\underline{v}}{dt} = -\underline{\nabla}\Phi - \frac{1}{\rho} \cdot \left(\underline{\nabla}P + \underline{\nabla}\tilde{\Theta}\right) , \tag{2a}$$

P being the pressure, and $\tilde{\Theta}$ the viscous stress tensor (see e.g. Landau and Lifshitz, 1959). The gravitational potential Φ is that of a point mass, M:

$$\Phi = -\frac{G \cdot M}{r} . \tag{2b}$$

We assume the gas to be an ideal one (P_g: gas pressure), and take into account radiation (P_r). With the temperature T, and the molecular weight μ we get ($\Re$ is the gas constant.)

$$P = P_g + P_r = \frac{\rho \cdot \Re \cdot T}{\mu} + \frac{a}{3} \cdot T^4 . \tag{2c}$$

2.1.3. The Energy Transport Equation.

$$\frac{d\varepsilon}{dt} = -\frac{P}{\rho} \cdot \underline{\nabla}\underline{v} - \frac{1}{\rho} \cdot \left(\left(\tilde{\Theta}\underline{\nabla}\right)\underline{v} - \underline{\nabla}\underline{q}\right) . \tag{3a}$$

ε is the specific internal energy, and $\underline{q}$ the heat flux vector. After Baker and Kippenhahn (1962) we take ε to be

$$\varepsilon = \frac{3}{2} \cdot \frac{P}{\rho} \cdot (2 - \beta) , \tag{3b}$$

with β being the gas pressure in units of the total pressure:

$$\beta = \frac{P_g}{P} . \tag{3c}$$

The heat flux , $\underline{q}$, consists of a radiative ($\underline{q}_r$) and a convective ($\underline{q}_c$) part:

$$\underline{q} = \underline{q}_r + \underline{q}_c . \tag{3d}$$

For the radiative part we use the diffusion approximation, and for the convective contribution the mixing length formalism as described by Kippenhahn *et al.* (1967).

2.2. Approximations.

2.2.1. Symmetry.

We assume the accretion disk to be symmetric in azimuth, i.e. in the φ-direction, and in the z-direction.

2.2.2. Mass of the disk.

The mass of the accretion disk shall be negligible compared to the central mass (M); only this justifies a) taking (2b) for the potential, and b) neglecting selfgravitation in the vertical direction.

2.2.3. Velocities in the disk.

Velocities in the vertical, i.e. z-direction are neglected; which is equivalent to assuming that equlibrium in this direction is always reached on time scales shorter than the ones we are interested in.

The influence of pressure on the horizontal disk structure shall be very small compared to that of gravity:

$$\frac{\partial P}{\partial s} \ll \rho \cdot \frac{\partial \Phi}{\partial s} ; \tag{4a}$$

this means $v_s \ll v_\varphi$. We further assume that the velocities do not vary in the z-direction, i.e.

$$\frac{\partial v_s}{\partial z} = \frac{\partial v_\varphi}{\partial z} = 0 . \tag{4b}$$

2.2.4. Surface density, vertical scale height.

The time dependent evolution of the disk will be evaluated only for the equations integrated in vertical direction, and a surface density, Σ, is introduced:

$$\Sigma = \int_{-\infty}^{\infty} \rho \, dz . \tag{5a}$$

Futher a vertical scale height, h, is defined, where $\bar{\rho}$ is the density averaged in z-direction:

$$h = \frac{\Sigma}{2 \cdot \bar{\rho}} . \tag{5b}$$

2.2.5. Time scales.

Lightman (1974) has shown that in geometrically thin accretion disks the longest time scale is that over which variations due to viscosity occur. Only processes which run on this time scale are taken into account; this means that all other processes are regarded as reaching equilibrium instantaneously. During outbursts also processes on shorter time scales are included, if important.

2.3. Resulting Equations.

Applying the approximations described in Paragraph 2.2. to the equations introduced in Paragraph 2.1. gives the set of equations to describe the stationary structure and time dependent evolution of α-accretion disks.

2.3.1. Stationary accretion disks.

For a (constant) mass flow rate, $\dot{M}$, one gets from the continuity and Navier-Stokes-equations (subscripts s, z, and φ denote the components of the vectors/tensors in the respective directions; in brackets we give the results for an even simpler approximation where integrations and differentiations with respect to z are replaced by multiplications with, and divisions by, h, resp.):

$$\begin{aligned}
\dot{M} &= -2\cdot\pi\cdot s\cdot v_s\cdot\int_{-\infty}^{\infty}\rho\,dz \quad (= -2\cdot\pi\cdot s\cdot\Sigma\cdot v_s)\ ,\\
v_\varphi &= \sqrt{\frac{G\cdot M}{s}}\ ,\\
\dot{M}\cdot s\cdot v_\varphi &= 2\cdot\pi\cdot s^2\cdot\int_{-\infty}^{\infty}\Theta_{s\varphi}\,dz+\dot{I}\quad\left(\approx 2\cdot\pi\cdot s^2\cdot\Theta_{s\varphi}\cdot(2\cdot h)+\dot{I}\right)\ ,\\
\frac{\partial P}{\partial z} &= -\frac{G\cdot M}{s^2}\cdot\frac{z}{s}\cdot\rho\quad\left(\Rightarrow P\approx\frac{\Sigma\cdot h\cdot\Omega^2}{2}\right)\ .
\end{aligned} \tag{6a}$$

$\dot{I}$ is the net flux of angular momentum; $\dot{M}$ and $\dot{I}$ are determined by boundary conditions. In all models we give $\dot{M}$ as the mass inflow rate at the disk's outer boundary, and assume no net flux of angular momentum, i.e. $\dot{I}=0$.

In geometrically thin accretion disks heat flows predominantly in the vertical direction, i.e.

$$\underline{q}\approx q_z\cdot\underline{e}_z\ , \tag{6b}$$

where $\underline{e}_z$ is a vector in the z-direction of unity length. From Eqs. (3) we get:

$$\begin{aligned}
\frac{\partial q_z}{\partial z} &= -\Theta_{s\varphi}\cdot s\cdot\frac{d\Omega}{ds},\\
\frac{\partial T}{\partial z} &= \begin{cases} -\frac{3\cdot\kappa\cdot\rho}{4\cdot a\cdot c\cdot T^3}\cdot q_z\ , & \text{for radiative energy transport,}\\ \frac{G\cdot M}{s^2}\cdot\frac{z}{s}\cdot\rho\cdot\frac{T}{P}\cdot\nabla_{conv}\ , & \text{for convective energy transport,}\end{cases}
\end{aligned} \tag{6c}$$

where ∇_{conv} is the convective gradient, κ the opacity, and $\Omega = v_\varphi/s$ the angular velocity.

2.3.2. Time dependent evolution of accretion disks.

The evolution of the disk is described – under these approximations – by a diffusion-type equation:

$$\frac{\partial\Sigma}{dt}-\frac{3}{s}\cdot\frac{\partial}{\partial s}\left(\sqrt{s}\cdot\frac{\partial}{\partial s}\left(\sqrt{s}\cdot f\right)\right)=0\ , \tag{7a}$$

with the normalized stress tensor element integrated in vertical direction,

$$f=\frac{2}{3\cdot\Omega}\cdot\int_{-\infty}^{\infty}\Theta_{s\varphi}\,dz\ . \tag{7b}$$

2.4. Stationary accretion disks.

2.4.1. Material functions.

In order to solve Eqs. (6) we have to define a set of material functions; for our models we use:

$$\begin{aligned}
\Theta_{s\varphi} &= \alpha\cdot P\ ,\\
\alpha &= \min\left(0.05,\ \ 5\cdot\left(\frac{h}{s}\right)^{1.5}\right)\ .
\end{aligned} \tag{8}$$

We choose α following Meyer and Meyer-Hofmeister (1983), but introduce also a satuaration value (Duschl, 1986a). The exact values in the α-prescription were chosen to give the best fit to the observed time scales.

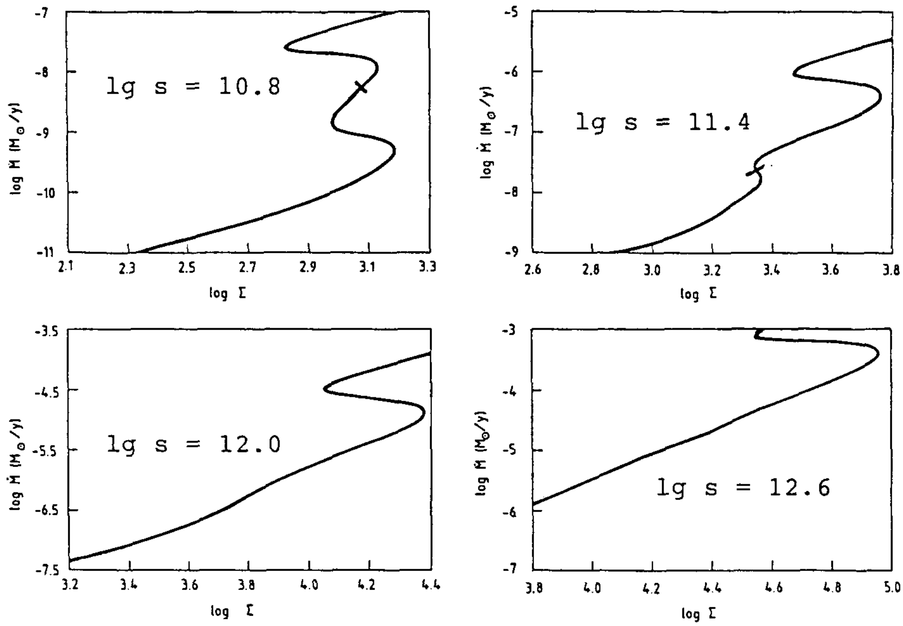

Figure 1: *The mass transfer rate as function of the surface density for different distances to the accretor.*

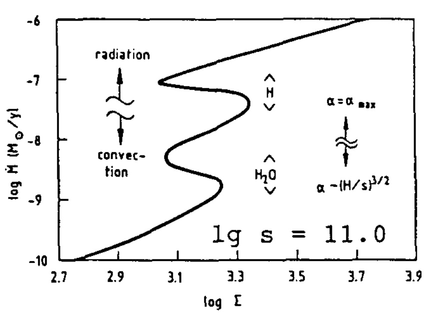

Figure 2: *The different domains of energy transport, viscosity parameter, and main contributors to the opacity, for a typical case ("H" indicates the region of ionization/recombination of hydrogen, "H_2O" the domain of disintegration/formation of water).*

We take κ from Cox and Stewart (1969) for $\log T \geq 4$, and Alexander (1975) for $\log T \leq 3.8$; in between we interpolate between the two sets of tables; the chemical composition is taken to be $X = 0.739, Y = 0.240$. The specific heat is calculated as described by Kippenhahn *et al.* (1967) for the regime of the Cox and Stewart-opacities, and taken from tables by Sharp (1981) for the lower temperatures. For the molecular weight we follow Kippenhahn *et al.* (1967), and Alexander (1975), resp.. The central accreting mass is assumed to be one solar mass.

Figs. 1 show the surface densities for different radii and mass transfer rates. Tickmarks on the $\log \Sigma - \log \dot{M}$-curve mark the points where the saturation value of α is reached; for higher mass transfer rates α is constant; at the two largest radii (10^{12}, and $10^{12.6}$cm) the entire curve shown has a constant viscosity parameter.

In Fig. 2 the different domains are indicated for a characteristic radius, 10^{11} cm. The saturation of the viscosity was already discussed above. The two distinct *S*-shaped features in the curve result from strong changes of the opacity within a comparatively small temperature range: at about $\dot{M} = 10^{-7.2}$ M$\odot$ */yr* it is due to the ionization/recombination of hydrogen, while at around $10^{-8.6}$ M$\odot$ */yr* it is due to the formation/disintegration of the water molecule. These S-shaped features indicate the regions of the disk instabilities.

The resulting accretion disks a) are geometrically thin, and b) have negligible mass compared to M, so that the approximations introduced above are justified. In contrast to the results for dwarf novae the accretion disks in SS are always optically thick.

3. ... AND THEIR OUTBURSTS ...

There are basically two different reasons why accretion disks may show outbursts: Either the mass inflow into the disk changes strongly, or – at an essentially constant inflow rate – there exists an intrinsic instability in the disk.

Integrating the energy equation (6c) in the vertical direction and introducing an effective temperature, T_{eff}, one obtains locally a relation between T_{eff} and $\dot{M}$:

$$\sigma \cdot T_{eff}^4 = \frac{3}{8 \cdot \pi} \cdot \frac{G \cdot M \cdot \dot{M}}{s^3} , \tag{9}$$

i.e. the higher the local mass flux, the higher is the corresponding temperature. This shows that an outburst has to be associated with an increased mass flow – at least somewhere in the disk.

The remaining, but crucial question is which physical reason stands behind the increase in $\dot{M}$.

3.1. The mass transfer instability (=MTI).

Bath proposed in 1972 that actually the instability is situated in the atmosphere of the companion; because of such an instability the mass overflow to the accretion disk may vary quite strongly, and cause thus the outbursts. Although there seem to exist some problems (Gilliland, 1985) the question whether such an instability exists or not is not settled as yet (Bath, 1975; Edwards, 1985; among others).

In the following we shall not deal with the question how such an increase of $\dot{M}$ can originate, but take for the time being the optimistic point of view that some suitable mechanism exists.

Eq. (7a) allows us to define an accretion time scale, τ_{accr}, that describes the time within which variations of the mass flow at some radius – because of whatever reason – reach the inner edge of the disk due to viscous transport:

$$\tau_{accr} = \frac{s \cdot v_\varphi}{\alpha \cdot c_s^2} , \tag{10a}$$

where we used the fact that according to Eqs. (6a) (integrated in the z-direction)

$$\frac{h}{s} = \frac{c_s}{v_\varphi} . \tag{10b}$$

Comparing the fastest changes in the observed light curves with τ_{accr}, one obtains a lower limit for the viscosity parameter α, as variations on time scales shorter than τ_{accr} cannot be caused by a viscous process that involves large parts of the disk. Further the definition of α gives an upper limit for it as one expects turbulence to be sub-sonic (assuming that the length scale of the turbulence is of the order of the smallest length scale in the disk, the vertical scale height h):

$$\alpha \not> 1 . \tag{10c}$$

Bath and Pringle (1982) find that the MTI model needs $\alpha \not< 0.2$ to be capable of reproducing the light curves of SS. For smaller α, τ_{accr} becomes so large that not even a step-like change of $\dot{M}$ at the outer edge suffices to reproduce fast modulations of the luminosity.

The results of model calculations for SS in the framework of the MTI will be presented and discussed in Sect. 4.

3.2. The disk instability (=DI).

Meyer and Meyer-Hofmeister (1981) were the first to calculate the stationary vertical disk structure with a realistic description of the energy transport a n d the opacity. They realized that an instability as proposed by Osaki (1974) may exist since in some regions Lightman's (1974) stability criterion

$$\frac{\partial \dot{M}}{\partial \Sigma} \geq 0 \tag{11}$$

is violated. (The exact form of this condition depends on the choice of $\dot{I}$ (see Eq. (6a)); for $\dot{I}$ as chosen here, Eq. (11) is exact, for other $\dot{I}$ the principles of this reasoning do not change.)

The calculations presented in Figs. 1 and 2 show that for a certain range of $\dot{M}$ and s condition (11) is violated. The resulting behaviour of the disk is shown in Fig. 3.

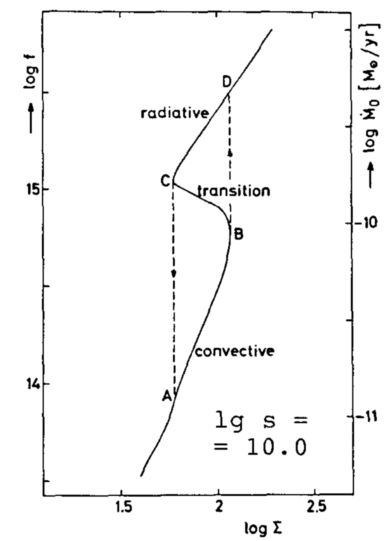

Figure 3: *The limit cycle ABCD in an unstable disk (the Figure is taken from Meyer and Meyer-Hofmeister's (1981) models for dwarf novae; in addition to mass flow rate and surface density, it also shows the viscosity integral (Eq. (7b)).*

If at some radius the stationary solution lies on an unstable branch, we shall instead have a limit cycle type evolution of the disk. This will influence neighbouring regions thus leading to a *cooling* or *heating* front moving through the disk. "Cooling (heating) front" means that the transition from the stable branch with high (low) local mass flow, i.e. effective temperature, to the other stable branch is initiated. Such a front can move through the disk until it reaches zones where for the involved range of $\dot{M}$ only stable solutions exist. Meyer (1984) has discussed their behaviour in detail; other authors used slightly different physical and numerical approximations to describe this phenomenon (see e.g. Faulkner *et al.*, 1983; Mineshige and Osaki, 1983, 1985; Papaloizou *et al.*, 1983; and others).

Such fronts move through the disk with a velocity that is of the order of the sound velocity of the front's final state times the viscosity parameter α. Thus one can define a time scale, $\tau_{f,a}$, that describes within which timespan a front can change the energy output of an accretion disk markedly (a stands for c or h, which again stand for cooling, or heating front, resp.):

$$\tau_{f,a} = \frac{s}{\alpha \cdot c_s (T_a)} \cdot \tag{12a}$$

The (relatively) thinner an accretion disk, the shorter is τ_f compared to τ_{accr}:

$$\tau_f = \frac{h}{s} \cdot \tau_{accr} \; ; \tag{12b}$$

the type of the front determines h.

So in the DI model one can reproduce variations on shorter time scales than is possible with the MTI model, i.e. has a larger allowed range of values of α, although $\dot{M}$ is kept constant. On the other hand the rise and the decay time are related to each other (Eq. (12a)). With τ_{f+} being the rise, and τ_{f-} the decay time scale, we get:

$$\frac{\tau_{f+}}{\tau_{f-}} = \frac{\alpha_c \cdot \sqrt{T_c}}{\alpha_h \cdot \sqrt{T_h}} . \tag{12c}$$

Here we introduced the possibilty that the hot and the cool state may have different values of the viscosity parameter as is indicated from dwarf nova models (e.g. Meyer and Meyer-Hofmeister, 1983; Mineshige and Osaki, 1983; Smak, 1984). A further restriction arises from the timescale within which outbursts follow each other. One expects the disk to be refilled after an outburst within the accretion time scale of the cool state, so the typical repetition time scale of outbursts, τ_{out}, is

$$\tau_{out} = \frac{s \cdot v_\varphi}{\alpha_c \cdot c_{s,c}^2}. \tag{12d}$$

All these arguments are only of the *order of magnitude* type but are shown to be reasonably accurate by numerical models that will be discussed later.

There are obviously stronger restrictions on the parameters in the DI model than there are in the MTI model. The reason is that in the latter one is free to choose any suitable form for the mass overflow rate as one cannot give a selfconsistent physical model for the actual form as yet. In contrast to that, in the DI model these restrictions come from the physical processes that determine the outburst; this means that the DI model has a higher degree of selfconsistency compared to the MTI model. In this sense the higher degree of freedom in the latter ansatz actually results from the partial lack of a physical description, rather than being model inherent.

4. ... IN SYMBIOTIC STARS.

In the following we shall discuss the two outburst models and their application to SS; for the details we refer to the papers by Bath and Pringle (1982; MTI), and Duschl (1985a, 1986b; DI).

4.1. The accretor: main sequence star vs. white dwarf.

The first question is whether one expects a main sequence star, or a white dwarf to be the accreting object. Integrating Eq. (9) one gets for a stationary accretion disk a luminosity, L, of

$$L = \frac{G \cdot M \cdot \dot{M}}{s_*} , \tag{13}$$

where s_* is the radius of the accreting star. One can deduce from this for both cases some average accretion rates; in the case of a white dwarf these would be smaller by a factor of the order 10^2. According to Eq. (10a) we have $\tau_{accr} \propto T_{eff}^{-1}$, estimating the relation between temperature and mass inflow at the (here relevant) outer radius of the disk (that does not depend on the central object) we find that $T_{eff} \propto \dot{M}^{1/4}$. This would be equivalent to a higher accretion time scale for a white dwarf accretor. We give as an example the following numbers (s_0: outer radius of the disk): $M = M_\odot$; $s_* = R_\odot$; $s_0 = 10^{12}\ cm$; $L = 10^{36}\ erg/s$, and $\alpha = 1$, and get an accretion time scale of the order one year, while for a white dwarf we would reach already several years, which is definitely too long. As we have natural upper limits for α and M (white dwarf), and observed values for L, the only way out would be a much smaller s_0. From the observed orbital periods of SS, and from the relative sizes of accretion disks in dwarf novae (e.g. Smak, 1982), we regard that as very unlikely. This argument only uses τ_{accr} and thus applies strictly only to the MTI. But similar arguments are true for τ_{out}. So both models point strongly towards a main sequence accretor.

This is in agreement with the analysis of spectra of SS (e.g. Mikołajewska and Mikołajewski, 1983; Kenyon and Webbink, 1984; Kenyon, 1985, 1986; for a more detailed compilation see Kenyon, 1986).

4.2. The MTI ansatz.

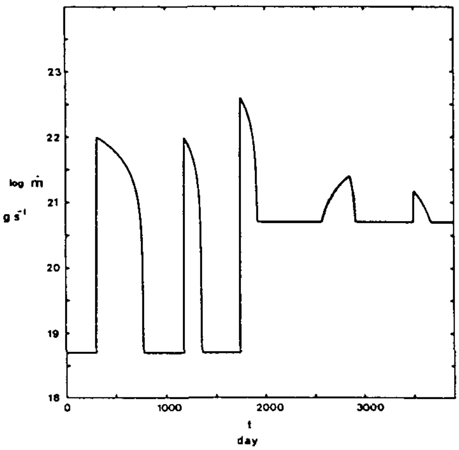

Figure 4: *The evolution of the mass transfer rate for CI Cyg in the model calculations by Bath and Pringle (1982); the Figure is taken from their paper.*

Bath (1977), and Bath and Pringle (1982) investigated how far the MTI model is capable of reproducing the lightcurves of SS. They chose *CI Cyg* as their example. In Figure 4 we show the evolution of the mass transfer rate they needed for their best fit of the lightcurve. Their model dimensions were the following: M=M$\odot$; $s_* = R_\odot$; $s_0 = 8.5 \cdot 10^{12}$ cm; $\alpha = 1$. They find good agreement with the observed lightcurves. But there remains a problem with the evolution of the mass transfer rate which this model needs to reproduce the lightcurves: a) The minimum mass overflow rate between outbursts has to vary itself by more than order of magnitude within few orbital periods; this could be attributed to long term variations in the companion's atmosphere, so that it does not seem to be a serious problem, although it is a remarkable feature in the evolution of the transfer rate; b) The structure of the bursts of $\dot{M}$ itself varies; while most bursts show a very sharp increase over several orders of magnitude, followed by a much shallower decline, at least for some cases the opposite configuration is necessary; this points towards two different types of instabilities being needed in the companion star's envelope (which remains to be explained).

The basic assumption of the disk theory presented here is that the disks are geometrically thin; as stated by Bath and Pringle, during the maxima of the outbursts this becomes a poor approximation as the disk's thickness reaches values of a third of the distance to the accretor, i.e. the neglected quantities (which are of relative order $(h/s)^2$) are less than an order of magnitude smaller then those taken into account, so one is at the limit of the applicability of the theory. This latter statement is also true for the DI ansatz.

4.3. The DI ansatz.

Duschl (1985a, 1986a, 1986b) and Mineshige (1986) have applied the DI model to SS. Since in this type of model the degrees of freedom are far fewer one cannot reasonably try to reproduce an individual system as this would need *fine tuning* of all parameters to an exactness that is far beyond what one can reach observationally; since one on the other hand does not expect a unique parameter set to fit all the observations, it is also not suitable as a part of observational diagnostics. One therefore attempts to model an average system and analyze which parameter range is allowed; these results are then compared with observations.

We also assume a one solar mass main sequence star as accretor; here we take the α-prescription as defined in Eq. (8). The outer radii are varied from $10^{11.6}$ to $10^{12.8}$ cm, the – constant – mass inflow rate between 10^{-6} and 10^{-3} M$\odot$ $/yr$. What one wants to reproduce are lightcurves with typical variations of the order of several month to years; the increase in luminosity is sharper than the decrease; one observes only minima that last for about the same timespan as – or less than –

the outbursts (in contrast to dwarf novae where the minimum state is far longer than the outburst). One can mark a region in the $s_0 - \dot{M}$-diagram where the parameter combinations lie that are capable of producing light curves like the ones observed in SS (Fig. 5). The limitations are the following: A: (dashed line) The disks are no longer geometrically thin, i.e. the approximations break down (this might very well be a technical rather than a physical limit); B: τ_{out} becomes too long as the disks become too large; C: τ_{out} becomes too long as $\dot{M}$ ($\Rightarrow T_{eff} \Rightarrow c_s$) becomes too small; D: The maximum brightness becomes so small that it is no longer consistent with observations; E: There are no bursts at all as the entire disk is stable.

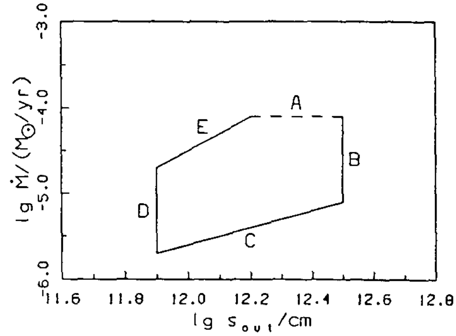

Figure 5: *The domain where the models for SS according to the DI model can be found.*

One might wonder how changes of this α-prescription would change these results. While the value of α determines the overall timescale of the evolution, it has almost no influence on the relation between τ_{out} and τ_f as h/s is only weakly dependent on α. So a change in the viscosity parameter would not leave unaffected the good absolute and relative agreement of time scales as well as the luminosities, that is achieved in the models. Changes in α would either change the frequency of outbursts or the luminosity; either are undesired effects thus making the possible range for an α-law rather small; this means that the results shown in Fig. 5 are not strongly dependent on α.

4.4. Comparison to dwarf nova models.

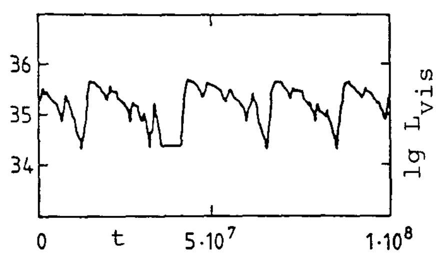

Figure 6: *A typical light curve for a SS in the DI model; the parameters for this model are: $\dot{M} = 10^{21}$ g/s, and lg $s_0 = 10^{12}$ cm.*

The MTI as well as the DI model were both originally developed to model dwarf nova outbursts (e.g. Bath and Pringle, 1981; Meyer and Meyer-Hofmeister, 1981). While the behaviour of accretion disks in the MTI model is essentially the same as in the case of dwarf novae – only scaled to the dimensions of SS – in the DI model new types of behaviour of the fronts show up. In dwarf novae one has usually only either a heating or a cooling front moving through the disk, or – in the minimum state – no fronts at all; in the models for SS one almost always finds fronts moving through the disk; often there are several fronts of both types moving through the disk; and these fronts can interact. Because of the different timescales (see Eq. (12c)) two types of interaction are possible: a) A slower moving cooling front is *overtaken* by a faster heating front, and the two fronts cancel each other out; b) Two fronts of the same type move towards each other, also cancelling themselves out. These different interactions are the reason for the typically not very smooth light curves as shown in Fig. 6. Duschl (1986b) gave a color figure showing the evolution of the fronts in an accretion disk of a SS; this may be compared to the case of dwarf novae (Duschl, 1985b; there the color bar was reversed during the printing process!)

4.5. Comparison of the two outburst models.

Both models are capable of reproducing the observed light curves reasonably well. The DI model has a higher degree of physical selfconsistency. In the field of dwarf novae where the comparison of both models with observations is in a much more detailed state than in that of SS, there are strong indications that the DI model is in better agreement with the observations although this question is far from being settled.

In both models one has T_{eff} during the maxima of the outbursts of several 10^4 K; in addition one expects a boundary layer between the disk and the star where the characteristic temperatures during the outbursts reach values of the same order, i.e. sensible temperatures for SS.

5. SUMMARY.

There exist many – mutually incompatible – models for SS, and hopefully nobody will expect one model to be applicable to all SS. For a subgroup of the *type-B*-SS accretion disks seem to be the main reason for the behaviour we observe. Members of this subgroup are stars like CI Cyg, Z And, and AR Pav.

ACKNOWLEDGEMENTS.

I thank Prof. R. Kippenhahn, and Drs. E. and F. Meyer for all I have learnt from them about accretion disks in many, very valuable discussions; I also thank Drs. H.-C. Thomas and S. White who read this manuscript and helped by their suggestions to improve it, and Drs. G.T. Bath, E. and F. Meyer, and J.E. Pringle, for allowing me to use figures from their papers in this review.

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