Accretion Disks in Transient Systems

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Abstract. I review recent advances in our understanding of accretion disks in transient systems — the dwarf novae and the soft X-ray transients. The primary theme will be the ongoing development of theory in response to the observations. The accretion disk limit cycle model appears to provide a unifying framework within which we may begin to understand what is seen in different types of interacting binary stars, and also to constrain parameters which enter into the theory.

1. Introduction

The last 15 years have witnessed major advances in our understanding of the steady state and time dependent nature of accretion disks in dwarf novae (DNe) and soft X-ray transients (SXTs). These systems consist of semi-detached binaries in which a Roche lobe filling K or M star transfers matter into an accretion disk surrounding a compact object. In the DNe the compact object is a white dwarf (WD), whereas in the SXTs it is a neutron star (NS) or black hole (BH). The advances discussed in this work have come about largely because of the development of the accretion disk limit cycle mechanism as a model to account for outbursts in DNe and SXTs. The outbursts one sees in these systems typically last for days to weeks and recur on intervals of weeks to years. Such short time scales, coupled with the relative brightnesses of some of the nearer systems, make these objects ideal laboratories for studying the physics of accretion disks.

In this review I will touch briefly on six aspects of the limit cycle instability — with an emphasis on the connection between observation and theory. These areas are:

(i) the steady state physics underlying the instability,
(ii) the general time dependent evolution of disks subject to the limit cycle instability,
(iii) the “problem” of the delay of the UV flux with respect to the optical during the onset of a fast-rise outburst,
(iv) Warner’s (1987) \( M_V(\text{peak}) \) versus \( P_{\text{orbital}} \) relation for DNe at maximum light, and its relevance to the \( M_V \) versus \( (L_X/L_{\text{Edd}})^{1/2}P_{\text{orbital}}^{2/3} \) relation for SXTs found by van Paradijs & McClintock (1994),

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(v) the implication of observed exponential decays of outbursts seen in dwarf novae and SXTs for the Shakura & Sunyaev (1973) viscosity parameter $\alpha$, and

(vi) the effect of evaporation of material from the inner disk on the long term light curves of the SXTs.

The work I will discuss makes the fundamental assumption that our steady state and time dependent modeling has direct relevance for the outbursts that we observe in DNe and SXTs, and that we can therefore work backwards from the observations to place limits on the magnitude of the viscosity parameter $\alpha$. In the early stages of this line of inquiry, most modelers were content with simply estimating how big $\alpha$ is. More recently we are trying to go one step further and place limits on the functional dependence of $\alpha$ on local physical variables. The work described by John Hawley represents the more fundamental effort in this regard as it attempts to calculate the viscous dissipation from first principles (Hawley, Gammie, & Balbus 1995, 1996). The ultimate goal is for workers in these two philosophically distinct camps to compare results and see how their findings about the viscosity — usually expressed in terms of $\alpha$ for historical reasons — agree or disagree. Perhaps in another five to 10 years workers will begin to be able to cross-compare findings between the two groups.

My personal strategy in carrying out the task of going from observation to theory is that I model the brightest and best studied objects. This means that for DNe, I take SS Cygni, and for SXTs, I take A0620-00. SS Cyg is the brightest DN in the sky. Its orbital period is $\sim 6.6$ hr (Joy 1956), and it shows outbursts of amplitude $\delta m_V \simeq 3.3$ which last for $\sim 10-20$ days and recur every $\sim 2$ months (Campbell 1934, Cannizzo & Mattei 1992). A0620-00 is a black hole candidate with $P_{\text{orbital}} \sim 7.75$ hr (McClintock & Remillard 1986) which had the brightest X-ray outburst seen in an SXT (Nova Mon 1975). It is the only SXT which lies within 1 kpc of Earth (Chen, Shrader, & Livio 1997).

2. Background/Steady State Physics

Although one can find in the pre-1980’s literature certain statements which foresaw in some way the possible operation of a disk instability as the agent responsible for DN outbursts, one can trace the roots of modern day research on the limit cycle instability back to the Sixth North American Workshop on Cataclysmic Variable Stars and Low Mass X-ray Binaries which was held in the summer of 1981 in Santa Cruz. At this meeting Jim Pringle discussed in general terms the relation between the integrated viscosity $\nu \Sigma$ and the surface density $\Sigma$ which an accretion disk must possess in order to obtain limit cycle behavior. The parameter $\nu$ is the “kinematic viscosity coefficient”, $\nu = 2\alpha P/(3\Omega \rho)$, and $\alpha$ is the Shakura & Sunyaev (1973) parametrization of the $\phi-r$ component of the viscous stress tensor expressed in terms of the local pressure $P$. When $\nu \Sigma$ is plotted in terms of $\Sigma$, one sees an “S” shaped curve. In the audience at this meeting were several accretion disk modelers, notably Friedrich Meyer, Doug Lin, and John Faulkner. Meyer had just completed a paper with his wife Emmi Meyer-Hofmeister on the vertical structure of accretion disks which was submitted the same month as the meeting (and appeared early in 1982), but nowhere in their paper had they thought to plot their solutions of the disk structure as $\nu \Sigma$ versus $\Sigma$. Traditionally, one had always plotted all physical quantities versus
radius. Upon his return to Germany, Meyer produced a short paper with Meyer-Hofmeister using some of the solutions taken from their structure paper. It was submitted to Astronomy & Astrophysics in October 1981, and it appeared in December 1981. Figure 2 in this article shows a plot of $\dot{M}$ (which is equivalent to $\nu\Sigma$, and which also scales with $T_{\text{midplane}}$ or $T_{\text{effective}}$) versus $\Sigma$ computed from the vertical structure at one radius and for one $\alpha$ value (Meyer & Meyer-Hofmeister 1981, 1982). In this figure we see the hysteresis between $\dot{M}$ and $\Sigma$.

Shortly thereafter other workers, many of whom were not at the Santa Cruz meeting, became aware of the importance of calculating the vertical structure of accretion disks, and numerous other papers appeared showing the S-curves for rates of accretion giving disk temperatures near $\sim 10^4$ K, the dividing line between neutral and ionized hydrogen (cf. Cannizzo 1993a for a complete list of references). The physics controlling the vertical structure depends heavily on the opacity. For $T \gtrsim 10^4$ K, the opacity is large because of the ready availability of free electrons to contribute to bound-free and free-free processes. The opacity $\kappa \sim \rho^b T^b$ with $b \sim -2$ to $-3$. For $T \lesssim 10^4$ K the opacity falls steeply with decreasing temperature as electrons and protons recombine. For $6000$ K $\lesssim T \lesssim 8000$ K, one has $b \sim 10 - 20$, and for $T \lesssim 6000$ K the opacity is small. The form of $\kappa$ at low temperatures is complicated due to its dependence on a variety of molecules (Cannizzo & Wheeler 1984). It is the large change in both the amplitude and functional form of $\kappa(\rho, T)$ between neutral and ionized gas which ultimately leads to the shift in the $\dot{M} - \Sigma$ (or, equivalently, $T_{\text{eff}} - \Sigma$) relation between the lower and upper branches of the S-curve.

The specific shape and scalings of the local minima $\Sigma_{\text{min}}$ and maxima $\Sigma_{\text{max}}$ in $T_{\text{eff}}(\Sigma)$ are dependent somewhat on the assumed strength of convection, usually cast in terms of $l/H_P$, the ratio of the mixing length to pressure scale height (Mineshige & Osaki 1983, Pojmański 1986, Cannizzo 1992, Ludwig, Meyer-Hofmeister, & Ritter 1994). Broadly speaking, $\Sigma_{\text{max}}$ and $\Sigma_{\text{min}}$ increase with $l/H_P$ – this effect being more pronounced for smaller $\alpha$ values. All studies to date which integrate the steady state vertical structure equations assume standard stellar mixing length theory, which is questionable for accretion disks because of the strong shear. Nevertheless, it seems difficult to imagine that a more correct treatment of convection by researchers in the distant future could somehow eliminate the S-curve and produce a monotonic $T_{\text{eff}} - \Sigma$ relation, given that the opacity change is more fundamental than convection to the overall $T_{\text{eff}} - \Sigma$ morphology. Indeed, the omnipresence of the hysteresis in so many articles by independent groups gives us confidence in its reality.

A separate issue is that of viscous dissipation and angular momentum transport induced by convective motions within the accretion disk vertical structure. John Hawley argued convincingly that any such mechanism related to turbulence which couples linearly to the shear cannot transport angular momentum outward as required (cf. Balbus, Hawley, & Stone 1996). Many previous workers had treated turbulent heat transport and angular momentum transport in accretion disks collectively as “passive contaminants” which have similar physical properties, in spite of much work to suggest the contrary (for a discussion see Balbus & Hawley 1997). Calculations show only a weak inward transport of angular momentum if convection is the relevant agent (Ryu & Goodman 1992, Stone & Balbus 1996). Although a few early limit cycle papers tried to connect
α to some function involving the ratio of the convective speed to the sound speed (e.g., Smak 1982, Cannizzo & Cameron 1988, Duschl 1989), no one currently active in the field continues to advocate such a model.

3. Time Dependent Behavior/Outburst Rise Time versus Triggering Location within the Disk

Soon after the first vertical structure papers appeared, workers began to consider the global, time-dependent evolution of the disk in response to the S-curve relation which exists at every radius. The instability was found to operate in the following manner: During the time corresponding to the quiescence interval for DNe, matter arriving at the outer disk edge from the secondary star piles up at large radii, thereby augmenting Σ. When one plots Σ as functions of r and t, one sees that, in quiescence, Σ(r) is bounded by Σ_{min}(r) and Σ_{max}(r). These two physical quantities scale linearly with radius, so the mass enclosed within a cylinder of radius r scales steeply with radius, M(r) ∼ r^3. Thus most of the mass lies at large radii. As mass continues to accumulate, eventually we must have Σ(r_{trigger}) > Σ_{max}(r_{trigger}) at some radius r_{trigger}. The radius of onset of the instability r_{trigger} can lie anywhere in the disk. As gas within this annulus becomes unstable, it begins to heat quite vigorously, and its movement in (T_{eff}, Σ) or (T_{mid}, Σ) takes it from the lower to the upper stable branch of the S-curve, although often by a circuitous route (Mineshige & Osaki 1985, Fig. 12). The viscosity coefficient ν scales as αT_{midplane}, so the factor of ∼ 100 increase in T_{midplane} and thence ν gives rise to an isolated, high viscosity annulus (Lin et al. 1985) which attempts to spread with great urgency into the neighboring regions of low viscosity. This spreading leads to a subsequent depletion of Σ in the immediate vicinity of r_{trigger}, and the propagation of two sharp spikes in Σ, one moving to smaller radii and one to larger radii. Within these spikes the relative contrast in Σ(r) can be as much as ∼ 10^2 − 10^3 against the background level. (One wonders whether some of our basic assumptions are breaking down within the spikes.) The spikes represent heating fronts which rush into the regions of the disk lying at both smaller and larger radii from r_{trigger}. Upon traversing a given radial location, they initiate a heating from the lower to upper S-curve branch. Ultimately, most of the disk transforms to the high state. It is interesting that regions for which Σ_{quiescence}(r) < Σ_{max}(r) before the arrival of the heating front are able to make the cold→hot transition quite readily, due to the large value of Σ within the heating front. The value of Σ(r) within the outwardly progressing heating front falls as the outer disk edge is approached, however, and at some point the condition Σ_{heating front}(r) > Σ_{max}(r) cannot be maintained. At this radius the spike in Σ associated with the heating front can no longer cause heating because it cannot augment the local, background Σ value to the required level. The spike stalls and collapses down, merging back into the local flow within the disk from whence it arose.

All of the activity just described accompanies the observed rise of an outburst in a DN or SXT. The Σ(r) profile which accompanied the quiescent state is radically transformed after the disk finishes making the low-to-high transition. The disk attempts to enforce the condition M(r) constant with radius as a minimum energy condition. It has been known for some time that, for steady
or near steady disks with opacity laws relevant for ionized gas, \( \Sigma(r) \sim r^{-3/4} \). This is nearly the inverse of the profile which existed in the disk during quiescence, therefore during the subsequent evolution one sees a juxtaposition of material from large to small radii. This activity greatly increases \( \Sigma \) near the inner edge and thereby increases the level of accretion onto the central object. Smak (1984=S84) and Cannizzo et al. (1986=CWP, see also Cannizzo 1996a) found that the rise time for outbursts depends on whether the outbursts begin at small or large radii in the disk – so-called type B or type A outbursts in Smak’s terminology, or “inside-out” versus “outside-in” outbursts in the language of CWP. The reason for the difference in rise times versus outburst triggering location can be understood by looking at the surface density evolution. For type B outbursts, the surface density at small radii is small initially, as is the total surface area of the hot portion of the disk \( A_{\text{hot}} \). The surface density near the central object sets the level of accretion onto the central object (Cannizzo 1996b), and therefore the EUV and soft X-ray fluxes are small. The optical flux \( f_V \) is set by \( A_{\text{hot}} \) (Cannizzo 1996b), and therefore \( f_V \) is also small initially for the type B outbursts. As the outward moving heating front travels to larger radii, both \( A_{\text{hot}} \) and \( \Sigma(r_{\text{inner}}) \) increase, thereby gradually raising the optical and EUV fluxes.

The behavior for type A outbursts differs dramatically. For large \( r_{\text{trigger}} \) values, the initial value of \( A_{\text{hot}} \) just after triggering is a substantial fraction of the total disk, which results in a rapid increase of the optical flux to its near-maximum value. Also, since most of the mass is stored at large radii, this means that the total amount of material in the inwardly moving heating spike is also quite large. When the inward traveling heating spike arrives at the inner edge, it contains so much mass that the EUV flux rises dramatically to its maximum value. One finds, for SS Cyg parameters for instance, that the heating front arrives at the inner edge about one day after the optical flux rises. Evidence to support this picture is provided by EUVE observations of SS Cyg showing an outburst for which the EUV flux increases by \( \sim 10^3 \) in a few hours about a day after the rise in visual flux (Mauche 1996). Mauche also presents observations of a slow rise outburst in SS Cyg which is presumably due to an outburst triggered at a small radius in the disk.

After the phase of triggering and heating of the disk has ended and the \( \Sigma(r) \) profile has shifted to \( \sim r^{-3/4} \), there are two possibilities for the subsequent evolution. If the surface density in the outer disk \( \Sigma(r_{\text{outer}}) < \Sigma_{\text{min}}(r_{\text{outer}}) \), and if irradiation of the outer disk is not significant, then a wave of cooling begins to form at large radii and starts propagating inward, causing successive annuli of gas to make the high-low transition as it passes. Just as the heating front operates by enforcing a local augmentation of the surface density, so the cooling wave operates by enforcing a local depletion of the surface density. This situation comes about because, at the interface between the hot, inner disk and the cool, outer disk there exists a vigorous outflow of gas. This outflow is in fact \( \sim 3 \) times the rate of accretion onto the central object. As the matter passes through the transition front, it cools and its viscosity plummets. The matter comes almost to a complete halt (in terms of its radial velocity \( v_r \)) once it exits the front at large radii. By the time the front has traversed the disk completely to \( r_{\text{inner}} \), the entire disk has been transformed back into neutral gas. The rate of outflow across the front so greatly exceeds the rate of accretion onto the central object.
that only $\sim 2 - 3\%$ of the gas stored in quiescence is accreted onto the central object.

If the cooling front is not able to propagate, either because $\Sigma(r_{\text{outer}}) > \Sigma_{\text{min}}(r_{\text{outer}})$, or because irradiation is strong enough to keep the outer disk gas ionized, then the disk evolution is much slower. One must rely on mass loss from the inner disk caused by accretion of gas onto the central object to lower the surface density. After some time has elapsed and the disk mass has decreased, one has $\Sigma(r_{\text{outer}}) < \Sigma_{\text{min}}(r_{\text{outer}})$, and/or $T_{\text{irradiation}}(r_{\text{outer}}) < T_{\text{eff}}(\Sigma_{\text{min}}(r_{\text{outer}}))$ so that the cooling front can form and begin to propagate. The subsequent evolution is then as described in the previous paragraph. One sees evidence of both viscous and thermal decay in the outbursts of SS Cyg. The long outbursts with their flat topped maxima give evidence for viscous decay followed by the faster thermal decay, whereas the short outbursts which rise to maximum and then almost immediately decay back to quiescence show only the thermal decay. In Cannizzo’s (1993b=C93b) model for SS Cyg, about 3% of the stored disk mass is accreted onto the WD during a short outburst, whereas about 30% is accreted during a long outburst.

4. The “Problem” of the UV Delay

In the computations of S84 and CWP which showed evidence for fast and slow rise outbursts, one sees that the optical flux precedes the FUV flux by $\sim 2/3$ day in the fast rise outbursts. This is due to the fact that the formation of the enlarged radiating area of the disk accompanying the initial triggering precedes by $\sim 1$ day the arrival of the heating spike at the inner disk edge. A comparison in CWP of the spacing between the rise in $V$ and 1050 Å flux as seen by the Voyager spacecraft for a fast rise outburst in SS Cyg showed theory and observation to agree. Both S84 and CWP assumed Planckian flux distributions locally for the disk. Subsequent work by Pringle, Verbunt, & Wade (1986), and Cannizzo & Kenyon (1987) cast doubt on CWP’s findings. In these later works, Kurucz or stellar-like flux distributions were assumed, and the delays between rising flux at $V$ and FUV wavelengths decreased dramatically. The spectral distributions in the later works, however, also showed sharp absorption edges due to the Lyman and Balmer series. La Dous (1989) compiled a listing of published Balmer decrements, most of them measured in SS Cyg in outburst, and found them to be smaller than expected from theory. One can understand why the stellar-like flux distributions give such small FUV delays by comparing stellar and Planckian accretion disk flux distributions for a disk with a given $T_{\text{eff}}$ profile (Wade 1984). In the stellar-like distributions, the Lyman edge effectively cuts off the spectrum below $\sim 912$ Å and redistributes much of this flux into regions between $\sim 1000$ and $\sim 2000$ Å. Therefore, if one compares fluxes at, say, $\sim 1050$ Å and $\sim 5500$ Å as in CWP, one finds the $f_{\text{FUV}}/f_V$ ratio to be much larger for the stellar/Kurucz distributions than for the Planckian distributions (Wade 1984, Fig. 1). This contributes to an earlier rise in $f_{\text{FUV}}$ with respect to $f_V$ than would have been expected in the Planckian models. The earlier works which utilized the more primitive Planckian flux distributions may have characterized better the relevant distributions insomuch as the $f_{\text{FUV}}/f_V$ ratios were concerned, because one does not see strong absorption edges in the observed
flux distributions of DNe. For this reason, I do not believe there exists a “UV problem” as has been repeated so often in the literature. Until one has in place a reliable set of basis functions for the accretion disk flux distributions, it may be better to use the light curves computed from \( V \) and \( \dot{M}(r_{\text{inner}}) \) to compare with the observed \( V \) and EUV fluxes, such as those reported in Mauche (1996).

5. Warner’s \( M_V(\text{peak}) - P_{\text{orbital}} \) Relation for DNe and the van Paradijs & McClintock \( M_V - (L_X/L_{\text{Edd}})^{1/2}P_{\text{orbital}}^{2/3} \) Relation for SXTs

Warner (1987) found an empirical relation between the absolute visual magnitude of dwarf novae at maximum light \( M_V(\text{peak}) \) and orbital period \( P_{\text{orbital}} \). One can imagine how such a relation might be expected physically in the limit cycle model. As noted earlier, during the interval of quiescence when matter is being accumulated as it arrives from the secondary star, the \( \Sigma(r) \) distribution is bounded by \( \Sigma_{\text{min}}(r) \) and \( \Sigma_{\text{max}}(r) \). The upper bound \( \Sigma_{\text{max}} \) ensures that the disk mass stored in quiescence cannot exceed a “maximum mass” \( M_{\text{max}} = \int 2\pi r dr \Sigma_{\text{max}}(r) \). Now, since \( \Sigma_{\text{max}}(r) \propto r \), the maximum mass will be an increasing function of orbital period. C93b derived an expression for \( M_{\text{max}}(r_{\text{outer}}) \) by equating the disk mass stored in quiescence to that in the near steady state disk which exists just after the \( \Sigma(r) \) profile has changed from its quiescent shape to its outburst shape. The expression obtained by C93b was plotted by Warner (1995) with the data used in his earlier work, and there appeared to be consistency. More recently, Cannizzo (1998a) ran time dependent models for a range of values of \( r_{\text{outer}} \) (equivalent to a range in orbital period), to test the accuracy of C93b’s analytical estimate.

Does such a relation exist for outbursts in the SXTs? Van Paradijs & McClintock (1994=vPMc) looked at the dependence of \( M_V \) on the soft X-ray flux \( L_X \), expressed in Eddington units, and \( P_{\text{orbital}} \). They determined a correlation between \( M_V \) and \( (L_X/L_{\text{Edd}})^{1/2}P_{\text{orbital}}^{2/3} \). VPMc present a toy model which shows that this relation is to be expected if the optical flux is produced primarily by the irradiation of the outer disk by X-rays coming from accretion onto the central object. Indeed, by comparing the vertical scale on the \( M_V \) plots presented in Warner (1987) and vPMc, we immediately see that SXTs are \( 3 - 7 \) mag brighter at peak outburst than DNe. For a system at \( P_{\text{orbital}} \sim 8 \) hr representative of our adopted canonical SXT A0620-00, we may estimate how large the irradiation effect is. The difference between the two empirically determined relations at \( P_{\text{orbital}} \sim 8 \) hr is \( \Delta M_V \sim 3 \) mag. Some of this difference is due to the larger effective area of the disk around a \( \sim 7M_\odot \) BH versus a \( \sim 1M_\odot \) WD, at the same \( P_{\text{orbital}} \). For \( P_{\text{orbital}} \) constant, the orbital separation scales as \( M_1^{1/3} \) (if \( M_1 >> M_2 \)), so the radiating area of the disk which gives rise to the optical flux scales as \( \sim M_1^{2/3} \). So in going from a \( \sim 0.7 - 1M_\odot \) WD to a \( \sim 7M_\odot \) BH, we increase the disk area by \( \sim 7^{2/3} - 10^{2/3} \sim 4 - 5 \), which reduces \( \Delta M_V \) by \( \sim 2 \) mag. Thus \( \Delta M_V(\text{corrected}) \sim 1 \) mag. During the discussion, van Paradijs pointed out that this is only a rough estimate: there may also be a \( \Delta T_{\text{eff}} \) correction in going from DNe to BH SXTs. In any event, in the next section we will directly compare the observed \( M_V(\text{peak}) \) value determined from the 1975 outburst of
A0620-00 with that given by our time dependent model for A0620-00. This will provide a more direct estimate.

6. The Ubiquitous Exponential Decays in DNe and SXTs

Bailey (1975) looked at the decay light curves of DNe spanning a range in orbital period and determined the decays to be roughly exponential — the associated decay time constant scaling with orbital period. The “Bailey relation” is $t_e(V) \simeq 3 \text{ days} \left( \frac{P_{\text{orbital}}}{P_{\text{orbital}}(\text{SS Cyg})} \right)$. It is interesting to note that, for many of the brightest and best studied SXTs, including A0620-00, one also finds exponential decays. The associated $e^{-}$-folding time constant is roughly a factor of 10 slower than that of DN decays. This can be seen directly in the famous plot of four BH SXT outburst light curves which has appeared in so many review articles (cf. Tanaka & Shibazaki 1996, Fig. 3). In general, all the time scales associated with outbursts in the SXTs are much longer than those associated with outbursts in the DNe. This seems to be true for both the NS SXTs and the BH SXTs, therefore arguing for an irradiation effect and not a scaling with central object mass to account for the differences — assuming the basic limit cycle model to be at work in both systems. Another possibility, which I will argue for below, is that irradiation accounts for the difference in going from DNe to NS SXTs, but differences in central object mass $M_1$ account for differences in the light curves in going from DNe to BH SXTs.

In his talk, Andrew King argued that the irradiation effect is strong for the BH SXTs in outburst (cf. King & Ritter 1997=KR). The irradiation flux scales as $r^{-2}$, whereas the viscous flux scales as $r^{-3}$ for a steady state disk. So, all other factors being equal, one would indeed expect irradiation to become important at large radii where the bulk of the optical flux is produced. KR argue that irradiation prevents the cooling front from forming and thereby keeps the entire disk in the hot, ionized state. Cannizzo (1994) pointed out that, if the limit cycle model is operating and the disk is in outburst, the profile of aspect ratio $h/r$ versus $r$ will acquire a convex shape. This would shield the outer disk from direct irradiation. It is not clear, however, how reliable the theoretical $h/r$ values are for actual disks. Also, indirect irradiation from, for example, a scattering corona above the disk, could alter this simplistic picture of the shielding.

The general scenario of the viscous decay advocated by KR has been studied in detail by many workers (e.g. Bath, Clarke, & Mantle 1986, Lyubarskii & Shakura 1987, Cannizzo, Lee, & Goodman 1990=CLG, Mineshige, Yamasaki, & Ishizaka 1993). In view of the complex calculations required to model the limit cycle instability in accretion disks, it is interesting to note that the viscous decay process — in which the disk mass can only decrease by virtue of mass loss onto the central object — is remarkably simple: the evolution can be expressed analytically. CLG present a convenient power law form for the decay of the disk mass $dM_{\text{disk}}/dt \sim t^{-q}$, where $q = (38 + 18a - 4b)/(32 + 17a - 2b)$ and the Rosseland mean opacity $\kappa = \kappa_0 \rho \alpha T^b$. This scaling assumes $\alpha$ to be constant. The parameter $q$ is relatively insensitive to the opacity: for electron scattering ($a = 0, b = 0$) we get $q = 19/16$, whereas for Kramer’s opacity ($a = 1, b = -3.5$) we get $q = 5/4$. If one plots this functional form for the decay on the same scale as the 1975 outburst of A0620-00 as seen in X-rays, one sees that it does
not at all resemble the observed decay. The reason for this failure can be seen trivially in the structure of the diffusion equation for surface density, $\partial \Sigma / \partial t \propto \text{fcn}(\nu \Sigma)$, where $\nu \propto \alpha_{\text{midplane}}$. For $\alpha$ constant, one typically has $T_{\text{midplane}} \propto \Sigma^{2/3}$ (this is exact if the opacity is electron scattering), so that $\partial \Sigma / \partial t \propto \Sigma^{5/3}$. For exponential decays one needs $\partial \Sigma / \partial t \propto \Sigma$ so that $\Sigma(t) = \Sigma_0 e^{-t/t_0}$. One can ensure exponential decays by having $\nu$ be constant, therefore requiring either an isothermal disk if $\alpha$ is constant, or $\alpha \propto 1/T_{\text{midplane}}$. This scaling seems unphysical and would have rather dire consequences for the limit cycle picture.

What are the equivalent constraints imposed on $\alpha$ by starting with the assumption that the exponential decays are due to a cooling wave traversing the disk and shutting off the flow onto the central object? Vishniac & Wheeler (1996=VW, see also Vishniac 1997) present an analytical theory for the functional form of the decay of the light curve in an outburst caused by the limit cycle instability. Their model relies on deviations from steady state flow conditions within the inner, hot portion of the disk during the decay from maximum light. VW’s theory makes some detailed predictions about the flow patterns within the hot portion of the disk which are substantially borne out in detailed testing (Cannizzo 1998b). One may use their analytical formalism to derive an expression for the rate of decay of the [hot] disk mass for $\alpha$ constant. If one plots this functional form for the decay on the same scale as the 1975 outburst of A0620-00 as seen in X-rays, one sees that it does not at all resemble the observed decay. The reason for this failure can be understood by considering the scaling of the cooling front velocity with time. For $\alpha$ constant, the cooling front has a faster-than-exponential speed (when expressed as a function of radial position) leading to a faster-than-exponential decay of the light curve. This is in contrast with the purely viscous decay model, in which $\alpha = \text{constant}$ leads to a power law (slow-than-exponential) decay. Therefore, to get the cooling wave model to work, we have just the opposite problem as with the viscous decay model: now we need to have $\alpha$ decrease as the outburst proceeds. Cannizzo et al. (1995=CCL) tested functional forms for $\alpha$ and concluded that the exponential decays seen in the BH SXTs could be produced with a cooling wave provided $\alpha = \alpha_0 (h/r)^n$, where $\alpha_0 \simeq 50$ and $n \simeq 1.5$. This form for $\alpha$ was first used by Meyer & Meyer-Hofmeister (1983, 1984) in their modeling of DN outbursts, although they did not justify it in detail observationally. In CCL’s model, the $e^{-\text{folding decay time}} |d \ln M_{\text{disk}} / dt|^{-1}$ is given by $t_e \simeq 0.4 G M_1 \alpha_0^{-1} c_s^{-3}$, where $c_s$ is a constant $\sim 16 \text{ km s}^{-1}$. (As a minor aside, the value $\alpha_0 \simeq 50$ presupposes $M_1 \sim 10 M_\odot$ in the BH SXTs. A recent compilation of inferred SXT BH masses by Bailyn et al. (1996) shows a strong peak in the distribution at $\sim 7 M_\odot$. If true, this would lower $\alpha_0$ to $\sim 35$.) The scaling $t_e \propto M_1$ would account naturally for the factor of $\sim 10$ difference between DNe and BH SXTs. It would not explain NS SXTs, where irradiation appears to be an important factor in determining the outburst characteristics (van Paradijs 1996).

It is interesting to note that neither the viscous model nor the cooling wave model for the decays of outbursts naturally produces an exponential decay if $\alpha$ is constant. One must assume that one of the models is correct, and then infer the constraint on $\alpha$ required to get the model to work. That fact that $\alpha = \text{constant}$ fails fits in with the theme of the talk by John Hawley. He stressed that we only use “$\alpha$” to enable standardized comparison of results between
different groups of workers, and that the mere usage of “α” as a concept can engender a false sense of understanding or foster predisposed notions about the physics of viscous dissipation and angular momentum transport and their (untenable) connection with turbulence. The physical dependence of α on local or global physical conditions within the disk may require many years of MHD modeling to ascertain. Such efforts to date have only gone as far as determining \( \alpha \simeq 0.01 - 0.1 \), generally consistent with values inferred from studies of DN and SXT outbursts. As a further note, we stress that the aforementioned constraints derived from the exponential decays only apply to \( \alpha_{\text{hot}} \) which is pertinent to the ionized gas on the upper stable branch of the S-curve, for it is the hot part of the accretion disks that give rise to the optical and soft X-ray fluxes we observe. The viscosity parameter \( \alpha_{\text{cold}} \) along the lower branch of the S-curve of relevance to the neutral gas is set primarily by the recurrence time scales for outbursts (Cannizzo, Shafter, & Wheeler 1988). It may prove difficult to constrain a functional form for \( \alpha_{\text{cold}} \).

There is one additional consideration involving the application of the law \( \alpha \approx 50(h/r)^{1.5} \) to DNe. CCL showed, by varying the outer radius of the disk in a series of models, that the e−folding time for outbursts is relatively insensitive to \( r_{\text{outer}} \) and therefore \( P_{\text{orbital}} \). This result stands in contrast to previous disk instability computations which showed general consistency between theory and observation as regards the Bailey relation (S84, Cannizzo 1994). A resolution of this contradiction may lie within recent work by Gu & Vishniac (1998), who redo the vertical structure computations (which give us the S-curve) taking \( \alpha = 50(h/r)^{1.5} \). Rather than prescribe \( \alpha \) to be constant as in previous works, Gu & Vishniac iterate to get \( \alpha \) for each value of \( \dot{M} \) in a series of models at fixed \( r \) and \( M_1 \). They find that, due to an inherent inverse scaling of \( h/r \) with \( M_1 \), the \( \alpha \) values along the upper branch of the S-curve tend to be larger for \( M_1 = 1M_\odot \) than for \( M_1 = 10M_\odot \). In fact, \( \alpha_{\text{hot}} \) exceeds the typical values which S84 found are required to get the Bailey relation. Therefore \( 50(h/r)^{1.5} \) cannot be the entire scaling law: whatever physical mechanism produces this scaling (if it is in fact correct) must saturate at some \( (h/r)_{\text{crit}} \) so that \( \alpha_{\text{crit}} \sim 0.1 - 0.2 \). The existence of a saturation value to \( \alpha \) would restore the Bailey relation for DN parameters.

While we have many convincing examples of quantifiably exponential decays in the DNe and SXTs, do we have any good examples of power law decay? As mentioned earlier, the flat topped outbursts in SS Cyg are supposedly viscous rather than thermal, but unfortunately the dynamic range in flux level of a flat topped outburst is not large enough to reveal what the functional form of the decay is. Only the WZ Sge stars (including WZ Sge itself) have enough dynamic range during their superoutbursts to be able to see the slower-than-exponential form characteristic of viscous decay, a feature also seen in models for superoutbursts (Osaki 1996). The only other example I have been able to find is V1057 Cyg (Cannizzo 1996a), an FU Ori star.

7. Long Term Light Curves of the SXTs

The \( \alpha \) value inferred by CCL was based solely on the decay properties of the BH SXTs. The obvious next step is to run models covering complete cycles of outburst and quiescence for parameters appropriate to A0620-00 to see if the
outbursts bear any resemblance to those observed. Two outbursts have been seen in A0620-00 (Nova Mon 1917 and 1975). Cannizzo (1998b) performs long runs using the time dependent model described in CCL and finds that the outburst properties are quite different from those observed. Although the decays of the outbursts have the \( \sim 30 \) day e–folding time scale, the outbursts occur every \( \sim 5 \) yrs and have rise times much slower than those observed. Most of the outbursts are also low amplitude, \( \dot{M}_{\text{peak}} \sim 10^{15} - 10^{17} \text{ g s}^{-1} \), in contrast to the value \( \sim 10^{18} \text{ g s}^{-1} \) inferred for A0620-00. The 1975 outburst of A0620-00 made it a sustained \( \sim 50 \) Crab source for \( \sim 2 \) wks in 3–6 keV X-rays as observed by Ariel 5. Much fainter outbursts, had they existed, would have been observed with a host of X-ray satellites over the past \( \sim 20 \) yrs. Therefore one must conclude that the standard model does not work very well.

It has been known for some time that the quiescent observations of DNe and SXTs are not consistent with theory. One sees substantial EUV and soft X-ray fluxes in quiescence, and yet in the standard limit cycle disk instability model one expects negligibly small rates of accretion onto the compact object. Some mechanism extrinsic to the standard model must evaporate or remove gas from the inner disk and deposit it onto the compact object in such a fashion that it does not alter the quiescent \( \Sigma(r) \) profiles to the extent that an outburst is triggered. Meyer & Meyer-Hofmeister (1994) proposed that a hot, “coronal siphon flow” exists in the inner disk. Hot electrons in a corona conduct heat down to the disk photospheric layers and drive a flow onto the central object. Liu, Meyer, & Meyer-Hofmeister (1995) show general consistency between theory and observation for certain DNe. While the exact physics discussed by Liu et al. may or may not be relevant, it is clear that, in order to have a complete theory for the light curves of DNe and BH SXTs covering not only outbursts but also quiescence, some evaporative process must be at work. Therefore for full generality we should include the effects of evaporation or mass removal from the inner disk in our long term computations.

C97b finds that when evaporation is included, the frequent, small amplitude outbursts which are triggering at small radii in the standard model are eliminated. One is only left with major outbursts \( \dot{M}_{\text{peak}} \sim 10^{18} \text{ g s}^{-1} \) occurring every \( \sim 50 - 100 \) yrs, as observed. By eliminating the inner disk, the rise times of the outbursts are also shortened somewhat because the triggering must occur further out in the disk. (We move from having type B to type A outbursts.) The rise times are still slower than observed, however. Recent time dependent disk instability computations for the galactic microquasar GRS 1915+105 (Hameury et al. 1997) find that, for parameters relevant to this long orbital period/ high mass transfer rate system, the rise times are significantly affected by the altered \( \Sigma_{\text{quiescence}}(r) \) profile produced by introducing evaporation. It is not clear why our modeling of A0620-00 differs in this respect.

The observed 1975 outburst of A0620-00 had a peak visual magnitude corrected for extinction \( M_V(\text{peak}) = +0.7 \) (vPMc), a total energy \( \delta E = 3 \times 10^{44} \text{ ergs} \) (Chen, Shrader, & Livio 1997) which gives an accreted mass \( \delta M \sim 5 \times 10^{24} \) g (assuming a rest mass to energy conversion efficiency \( \epsilon \sim 0.06 \)), and a peak rate of accretion within the disk \( \sim 10^{18} \text{ g s}^{-1} \). These values are all close to those found by C97b using the standard model plus evaporation from the inner disk at a rate \( \dot{M}_{\text{evap}} \sim 10^{14} \text{ g s}^{-1} \). We may estimate how strong irradiation
is at the peak of the outburst by directly comparing observation with theory. This should provide a more specific estimate for A0620-00 than was given in the previous section. Our computed $M_V \simeq +1$ for a face-on disk. Taking a 60° inclination thought to be appropriate for A0620-00 reduces the optical flux by $\sim \cos 60^\circ = 0.5$ which adds $\sim 1$ mag to $M_V$ (peak). This difference is similar to that found by previous workers who modeled $B$ and $V$ for the 1975 outburst of A0620-00 using nonirradiated time dependent models for the accretion disk limit cycle, and compared their flux levels directly with A0620-00 (Huang & Wheeler 1989, Mineshige & Wheeler 1989). Therefore $\Delta M_V$ between observation and theory (without irradiation included) is $\sim 1$ mag, in line with our previous estimate obtained by comparing the Warner and vPMc relations. As a final check, C97b also runs models with irradiation included (as described in Mineshige, Tuchman, & Wheeler 1990) and finds that, within the inner hot portion of the disk existing at the time corresponding to maximum optical light in the outburst, the ratio of the local irradiation temperature to the local (viscous) effective temperature is $\sim 0.3 - 0.4$. Therefore, were we to recompute the visual magnitude with irradiation included, this would increase the effective temperature in the outer portion of the hot, inner disk by $\sim 1.3 - 1.4$. Since the visual flux scales as $\sim T_{\text{eff}}^2$ for the temperature region of interest (vPMc), this would increase the optical flux by $\sim 2$ and account for the $\Delta M_V \sim 1$ discrepancy between the observed value of $M_V$ (peak) and the value taken from the nonirradiated models. If most of the area of the entire disk (going out to $\sim 0.7$ of the Roche lobe of the primary) were strongly irradiated as in KR so that the cooling front was prevented from forming and propagating, then the irradiation-induced optical flux would far exceed that observed.

The model of C97b which includes the standard model plus evaporation for A0620-00 has $\dot{M}_{\text{evap}} \sim 10^{14}$ g s$^{-1}$, and yet the value inferred by McClintock, Horne, & Remillard (1995) for the quiescent value of $\dot{M}$ in A0620-00 inferred from ROSAT observations is $\sim 10^{10} - 10^{11}$ g s$^{-1}$. This assumes an accretion efficiency $\epsilon \sim 0.1$. The discrepancy might be accounted for in the models of Narayan, McClintock, & Yi (1996) which assume a much smaller efficiency $\epsilon \sim 10^{-4}$ for the quiescent state. On the other hand, if the model of Narayan et al. were correct, one would expect a large difference $\sim 10^3$ between the quiescent flux level of soft X-rays in going from the BH SXTs to the NS SXTs because, in the NS SXTs the infalling material strikes the neutron star surface and thus $\epsilon \sim 0.1$, whereas for the BH SXTs most of the thermal energy is supposedly advected through the event horizon. The observed quiescent fluxes in the NS and BH systems are actually quite similar, however (Chen, Shriver, & Livio 1997), despite claims to the contrary (Narayan, Garcia, & McClintock 1997). Although evaporation is attractive in some aspects, it appears to be problematic in other ways. C97b discusses other possible resolutions to the failure of the standard model (i.e., without resorting to evaporation) in the BH SXTs.

8. Conclusion

We present an overview of accretion disks in transient systems, with concentration on the limit cycle instability model which has been developed extensively to account for the outbursts seen in DNe and SXTs. In much of our discussion we
use the two brightest members of these classes in our time dependent modeling – SS Cyg for DNe and A0620-00 for SXTs. Our conclusions are drawn primarily from comparing theory with observation for these two systems. Since A0620-00 is a black hole candidate, our conclusions regarding the BH SXTs have been developed in more detail than for the NS SXTs.

Our conclusions are as follows:

(i) If the instability triggers near the inner disk edge it produces a slow rise outburst; if it triggers some distance away from the inner edge it produces a fast rise outburst. The differences are due to the rate of build-up in surface density near the compact object, and the rate of enlargement of a hot area in the disk which produces the optical flux.

(ii) Until “correct” flux distributions are used for the accretion disk, one cannot determine whether or not there exists a problem with the model as regards the delay of the FUV (i.e., flux from $\sim 1000 - 2000 \, \text{Å}$) with respect to the $V$ during a fast rise outburst.

(iii) For both DNe and SXTs there exists a relation between $M_V(\text{peak})$ and systemic parameters – $P_{\text{orbital}}$ for the DNe and $(L_X/L_{\text{Edd}})^{1/2} P_{\text{orbital}}^{2/3}$ for the SXTs. The absolute level of $M_V(\text{peak})$ between the two classes indicates brighter disks in the SXTs which is due in part to irradiation and in part to larger disks. The underpinnings of the Warner relation derive from the concept of the “maximum mass” which can be stored in the quiescent state before the next outburst is triggered (Cannizzo, Shafter, & Wheeler 1988).

(iv) The exponential decay in soft X-ray flux of outbursts seen in SXTs places constraints on the viscosity parameter $\alpha$. For constant $\alpha$, both the viscous and cooling wave models fail: the former produces a slower-than-exponential decline in $dM_{\text{disk}}/dt$, whereas the latter produces a faster-than-exponential decay. To avoid these unpleasant attributes and obtain agreement with observations, one would require $\alpha$ to increase with time as the outburst decays for the viscous decay model, and to decrease with time as the outburst decays in the cooling wave model. A calculation of $\alpha$ from first principles would help decide which option is more physically plausible.

(v) In the cooling wave model for the decay, one requires $\alpha = \alpha_0 (h/r)^n$, where $\alpha_0 \sim 35 - 50$ and $n \simeq 1.5$. For $n = 1.5$ the $e$–folding decay time for the mass contained within the inner, hot portion of the disk is $t_e \simeq 0.4 G M_1 \alpha_0^{-1} c_s^{-3} \propto M_1$, where $c_s \simeq 16 \, \text{km s}^{-1}$. The fact that the $e$–folding decay time of a system like A0620-00 is $\sim 30$ days whereas that for a system like SS Cyg is $\sim 3$ days would be a consequence of the ratio $M_1(\text{BH})/M_1(\text{WD})$. The outbursts in the NS SXTs (where $M_1 \sim 1 M_\odot$ as for DNe) are complicated by the effects of irradiation and would not presumably follow this law.

(vi) The “standard model” (i.e., $\alpha = 50 (h/r)^{1.5}$) of SXTs produces major outbursts with recurrence times and $\dot{M}_{\text{peak}}$ values similar to those observed, but one also sees frequent, low amplitude outbursts which were not seen in A0620-00. The introduction of evaporation of matter from the inner disk gets rid of the small outbursts and leaves only the major ones, but does not reduce the rise times of the major outbursts to a value as small as that seen in the 1975 outburst of A0620-00. This is the only shortcoming of our A0620-00 model and may prove to be minor. Alternatively, another factor such as a lower $\alpha_0$ value in quiescence may suppress type B outbursts. Hameury et al. (1997) find better
agreement between theory and observation as regards the rise time characteristics at different wavelengths using evaporation in disk instability computations for GRS 1915+105.

(vii) By comparing theory and observation for the value of $M_V$ at maximum light in A0620-00, we find a shortfall of $\sim 1$ mag in the theory which seems plausibly to be accounted for by introducing self-irradiation of the disk into the model, in agreement with previous workers.

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