STRUCTURE AND ENERGY TRANSFER OF UNSTABLE HOT STAR WINDS

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Abstract. Due to the instability of the radiation line force, the winds of hot, luminous stars should show a pronounced time-dependence resulting from the nonlinear growth of initially small perturbations. Following the method of Owocki, Castor & Rybicki (1988), we describe the time-dependent wind structure obtained with an independently developed code. Under the central assumption of *isothermality*, our results are in very good agreement with the ones by Owocki et al. We find that the response of the wind to periodic base perturbations remains largely periodic, at least up to $r \approx 2 \ldots 3\ R_*$, with no clear evidence of stochastic behaviour.

In order to test the foregoing assumption of isothermality and to compute the X-ray emission from models of structured winds, we have also incorporated the *energy equation* into our simulations. We encountered the numerical problem that all radiative cooling zones collapse because of the oscillatory thermal instability (cf. Langer et al. 1981). We present a method to hinder this collapse by changing the cooling function at low temperatures. The resulting wind shows *resolved* cooling zones; but, for a supergiant wind relatively close to the star ($r \lesssim 10\ R_*$), the macroscopic wind structure is very similar to isothermal calculations. Most of the hot material is caused by shell-shell collisions.

Key words: Stars: early-type – Stars: wind – Hydrodynamics – Instabilities

1. Introduction

The radiation driven winds of hot, luminous stars are thought to be subject to a very strong instability, which was first described by Lucy & Solomon (1970). As is now widely believed, the nonlinear growth of this instability leads to strong shocks in the wind, which are the supposed origin of many observed time-variable aspects of these winds, namely: X-ray emission (Cassinelli & Swank 1983, Hillier et al. 1993, and references therein); blue edge variability, black absorption troughs, and discrete absorption components in UV spectral lines (Prinja & Howarth 1986, Kaper et al. 1992, Prinja et al. 1992; this volume); variability of optical lines (Fullerton 1991, and this volume), especially Hα (Ebbets 1982); and non-thermal infrared and radio emission (Abbott et al. 1984). Because of the complexity of the radiation line force (Owocki et al. 1988) and the *evolutionary* character of the hydrodynamic flow from small initial disturbances to a fully developed nonlinear wind structure, a numerical treatment is necessary in modelling the time-dependent behaviour.

In the following we describe some recent efforts (cf. Feldmeier 1993) in calculating this structure. Since both the character of the wind instability (Lucy & Solomon 1970, MacGregor et al. 1979, Carlberg 1980, Abbott 1980,

Astrophysics and Space Science **221**: 393–408, 1994.

Owocki & Rybicki 1984, 1985, 1986, Lucy 1984, Rybicki et al. 1990) and the method of computing the unstable radiation force (Owocki et al. 1988, Owocki 1991) are described in the literature, we will concentrate on the *results* of the calculations.

The major assumptions herein are that: i) the flow is one-dimensional spherically symmetric; ii) the diffuse part of the line acceleration is calculated by the "Smooth Source Function" method (Owocki 1991); iii) the wind is optically thin in the continuum, *i. e.*, Wolf-Rayet stars are not treated; iv) the creation of structure is triggered by photospheric sound waves propagating into the wind. In §3, we further assume that the flow remains *isothermal*, even in the presence of strong shocks. This is equivalent to assuming that the radiative cooling times are much shorter than the competing dynamical time scales. This is tested in §4 by taking the energy equation including radiative cooling into account.

2. Computational Method

We solve the *hydrodynamical* equations with a time-explicit Eulerian code that uses the van Leer (1977) or piecewise parabolic advection schemes (Colella & Woodward 1984). This method is by now rather standard and described extensively in the literature (*e. g.*, Reile & Gehren 1991, Stone & Norman 1992; and references therein). When *energy transfer* is included, we solve the *thermal* energy equation, which is non-conservative, and where the correct amount of thermal energy at shocks is created by artificial viscosity (Schulz 1964, Winkler & Norman 1986). The *radiative cooling* is accounted for by an $\alpha = -1/2$ power law fit (cf. Castor 1987) to the Raymond et al. (1976) cooling function of an optically thin gas.

The *direct* part (absorption) of the *radiation line force* is calculated following the method of Owocki et al. (1988). The *diffuse* part (reemission) is calculated by a *formal solution* with presupposed local source function. This source function – and *only* it – is taken from the Sobolev approximation. This is the kernel of the *Smooth Source Function* method of Owocki (1991). All the following calculations assume an *optically thin* source function for purely geometric dilution of the photospheric continuum radiation field. (Calculations with optically thick source function from local Sobolev theory are shown in Puls et al., this volume).

The stellar model parameters in the following calculations are for a typical O supergiant similar to ζ Pup: $M = 42\ M_\odot$, $R = 19\ R_\odot$, $T_{\rm eff} = 42,000$ K, $N_{\rm He}/N_{\rm H} = 0.16$ (Kudritzki et al. 1983). The wind parameters are taken to be $v_\infty = 2,000$ km/s and $\dot{M} = 3 \cdot 10^{-6}\ M_\odot/{\rm yr}$.

Typically, we use 3,000 logarithmically spaced grid points to cover 4 or 9 stellar radii.

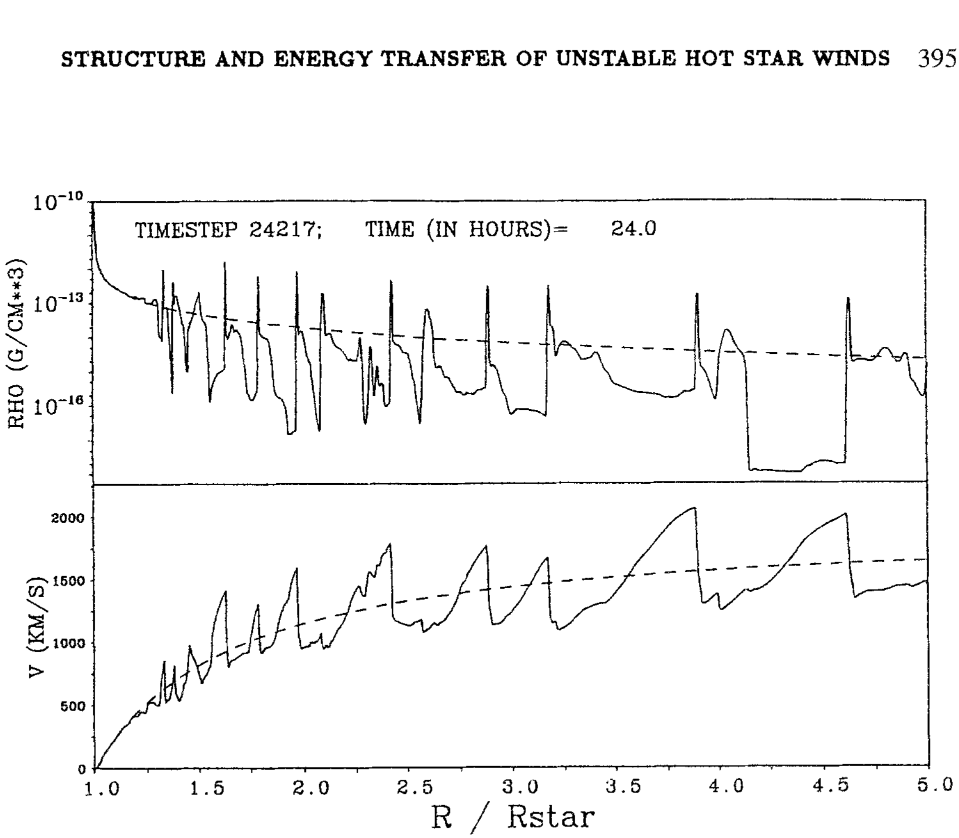

Fig. 1. Snapshot of the wind structure 24 hours after model start.

3. Isothermal Models

For the calculations in this section, the energy equation is trivially "solved" by setting $p = a^2\rho$, with p being the gas pressure and a the isothermal sound speed. The creation of structure is triggered by a sound wave of period 5,000 s (resulting in a wavelength of 2.8 times the photospheric barometric scale height H, where $H \approx 3.1 \cdot 10^{-3}\ R_*$) and amplitude $\delta\rho/\rho_0 = 1\%$ at the photosphere. The resulting flow, 24 hours of wind time after starting the model is shown in Fig. 1. Comparing this with the newest results by Owocki (see Fig. 1 in Puls et al. 1993, and this volume), we find remarkable agreement in spite of the different methods used and the strong instability. The wind consists of a pronounced shell structure, each shell being confined by a reverse shock (starward facing) and a forward shock (observer facing; but see below). A typical density contrast at the reverse shocks of $\rho_{\rm post}/\rho_{\rm pre} \approx 1,000$ corresponds to a Mach number of about 35, indicating *very* strong shocks (strong shocks being usually defined to have Ma $\gtrsim$ 5). The forward shocks are much weaker, and recent calculations indicate that they are not shocks at all, but *rarefaction waves*. Most of the wind's mass is concentrated in the

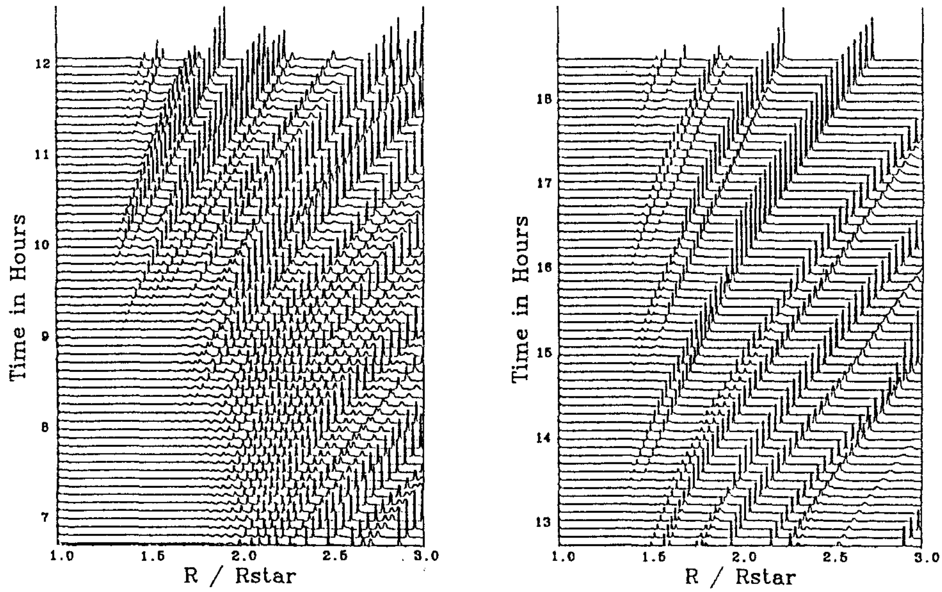

Fig. 2. Evolution of the mass loss rate of the time-dependent wind model over 12 hours.

very narrow (because of the high Mach numbers) shells. For a more extensive discussion of this structure we refer to Owocki et al. (1988) (cf. also this volume).

Fig. 2 shows the evolution of the mass loss rate in the unstable wind over a time interval of twelve hours, starting at $t \approx 7$ h. Up to about $t = 11 \ldots 12$ hours a lot of waves are seen. Their primary cause is not the photospheric sound wave but the (stationary) *initial condition*: the deviations of the latter from the actual stationary solution act as perturbations which are amplified by the instability of the radiation force. In the right panel, the wind has settled to a state of strong *periodic response* to the photospheric wave, with pronounced shells running out. Interestingly, the shells are created in pairs or triples, with a strong tendency for the members of these pairs and triples to collide and merge.

This multiple shell creation is connected with strong *overtone modes* present. The upper panel of Fig. 3 shows the time evolution of the normalized wind velocity over 60 cycles of the driving period at a fixed radius $r \approx 2\ R_*$. The lower panel shows the power spectrum of this signal. The photospheric wave corresponds to the peak of normalized height 1 at $2 \cdot 10^{-4}$ Hz, but the first 9 harmonic overtone modes of this fundamental frequency are

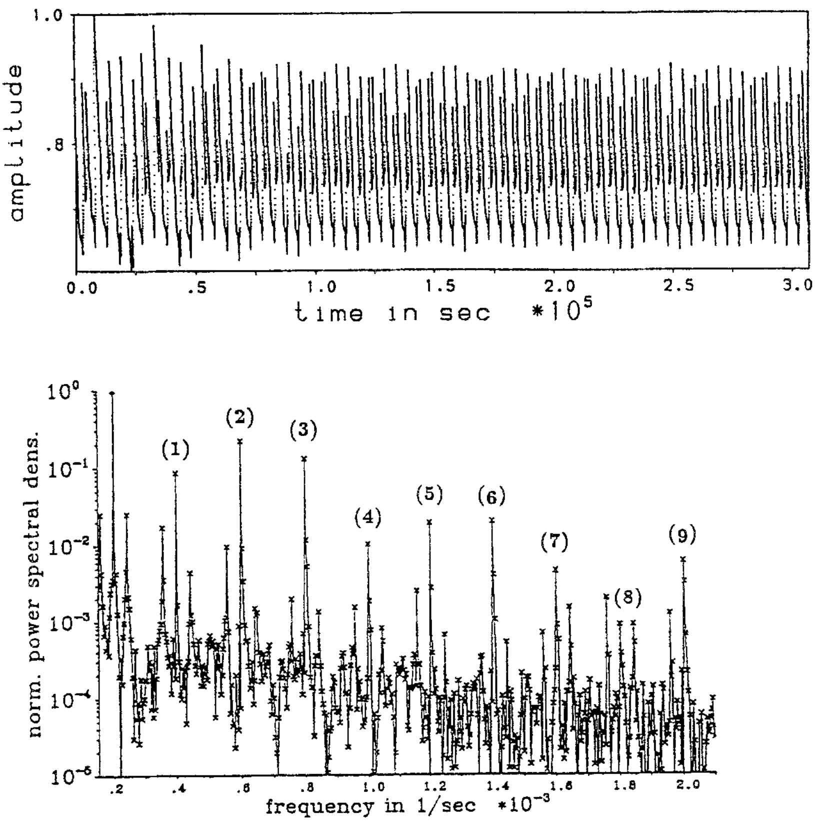

Fig. 3. *Upper panel*: Time evolution of normalized velocity at a fixed radius $r \approx 2\ R_*$. *Lower panel*: Power spectrum of this signal; numbers label the overtone modes.

also clearly seen. Such overtone modes are always created in nonlinear systems, and their cause here is the wave steepening due to the unstable growth. This wave steepening creates more and more high frequency components, their frequency being an integer multiple of the one of the fundamental wave (see Landau & Lifschitz 1991, p. 494).

Besides this photospheric sound wave and its overtone modes, no clear evidence of other frequencies is found in Fig. 3. In particular, there is no indication of a continuously filled domain of frequencies corresponding to stochastic behaviour (as is also rather obvious from the temporal signal itself). Therefore, the wind response to the periodic base perturbation remains (largely) periodic, at least up to $r \approx 2\ R_*$.

4. Inclusion of the Energy Equation

The motivation for including the energy balance in the calculation of the time-dependent wind structure is twofold. First of all, the assumption of dynamical isothermality has to be tested. Isothermality should not hold in regions of low density, that occur i) far away from the star or ii) in thin winds of, *e. g.*, OB Main Sequence stars like τ Sco, because the radiative cooling rates are small due to a lack of collisions. The second major motivation is that *X-ray emission* is one of the most direct indicators of variability in hot star winds, but it can only be accounted for in a proper manner by including the heating and cooling of the gas, *i. e.*, by resolving the radiative cooling zones.

Previous attempts to include the energy balance in the framework of time-dependent hydrodynamical simulations undertaken by Cooper & Owocki (1992) resulted in a serious shortcoming, in that radiative cooling zones could not be seen at all. A possible explanation of this behaviour was given by Owocki (priv. comm.), namely that the interaction of the diffusive errors of the advection scheme with the radiative cooling is an amplifying one: diffusive errors induce enhanced cooling, which in turn steepens flow structures leading to the build-up of new diffusive errors. We suggest another explanation, which also has the advantage of being able to explain the numerically observed collapse of *standing* radiative shocks with no advection.

The radiative cooling zone is subject to a thermal instability, that was discovered in numerical calculations of accretion onto white dwarfs by Langer et al. (1981). The linear stability analysis was done in a fundamental paper by Chevalier & Imamura (1982). The cause of this instability is the nonlinear dependence of the total length of the cooling zone on the immediate post shock temperature. The mechanism of the instability is explained, *e. g.*, by Chevalier & Imamura (1982), Gaetz et al. (1988), and Kinwah Wu et al. (1992). This instability is an *oscillatory* one (*i. e.*, it is an *over*stability, cf. Chandrasekhar 1961), which makes the position of the adiabatic shock at the very beginning of the cooling zone oscillate around its stationary rest position. Even if the numerical grid is fine enough to resolve the *stationary* cooling zone, the actual temporal cooling zone of minimum extent during this oscillation, *i. e.*, when the adiabatic shock is closest to the domain of cooled down material, may fall below the numerical grid resolution. As a result, the *reversible shock oscillation is turned into an irreversible collapse of the shock*. Once the radiative cooling zone is shorter than the grid length, the numerical scheme "forgets" about the existence of the cooling zone and the shock becomes a truly numerical isothermal one, which, of course, shows no cooling oscillation at all. (A more detailed description of this behaviour, as well as the corresponding numerical test calculations will be given in a forthcoming paper).

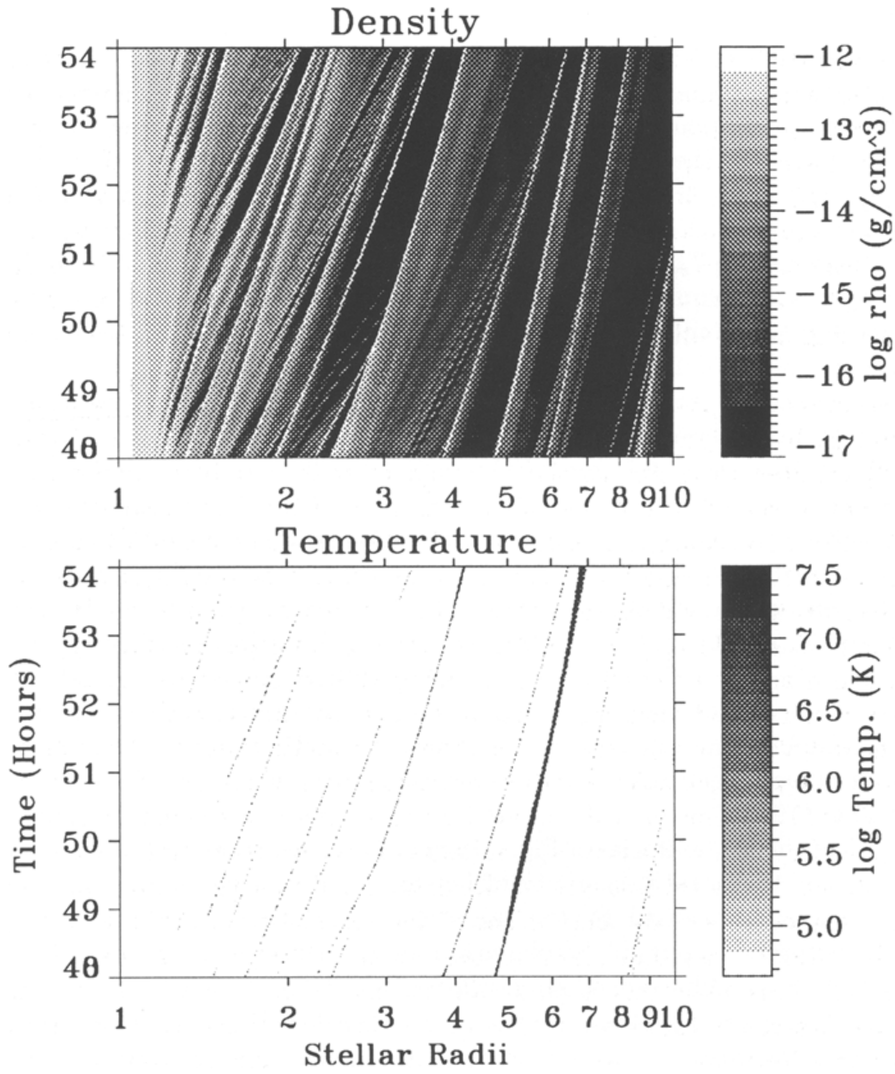

Fig. 4. Evolution of the wind structure over 6 hours, including the energy equation with radiative cooling. Cooling zones are not resolved.

Fig. 4 shows the evolution of an unstable wind (up to 10 R_* now) over 6 hours, from t = 48 to 54 hours after model start. The triggering sound wave is assumed to have T=10,000 s, with A=1 % at the photosphere. Although this calculation now includes the energy equation, the cooling zones are barely seen. The material cools as rapidly as it is heated, namely, on numerically unresolved length scales. The density plot of Fig. 4 again shows how frequent shell encounters are.

As a possible way to overcome this collapse, we suggest altering the cooling function at *low* temperatures so that there is no oscillatory instability in this region. If chosen properly, the stabilizing influence of the low temperature domain may then hinder the collapse. For a cooling function $\sim \rho^2 T^\alpha$, only exponents $\alpha \lesssim 1$ show the oscillatory instability (Chevalier & Imamura 1982, Imamura et al. 1984, Bertschinger 1986); thus, we assume that below some temperature $T_{\rm swi}$ the cooling function has an exponent $\alpha = 2$. At $T_{\rm swi}$ itself, the cooling function is assumed to be continuous.

Fig. 5 shows a calculation with $T_{\rm swi}/T_{\rm eff} = 10$, *i. e.*, $T_{\rm swi} = 420,000$ K. The tracks of the dense shells are very similar to the ones in Fig. 4, which indicates that the inclusion of radiative cooling has not changed the macroscopic dynamical behaviour of the wind. (Since nearly all cooling zones in Fig. 4 are collapsed, this model shows practically identical results to an *isothermal* one, and consequently we haven't shown the latter.) The cooling zones are narrow enough to justify the assumption of dynamical isothermality.

The temperatures behind strong shocks in Fig. 5 range from 10^6 K to some 10^7 K, giving a ratio $T_{\rm post}/T_{\rm eff} \approx 30\ldots1,000$ which is (much) larger than $T_{\rm swi}/T_{\rm eff} = 10$. Therefore, the influence of $T_{\rm swi}$ on the overall radiative cooling zone *should be small.* A calculation with twice the above value for $T_{\rm swi}$ results in almost the same wind structure. Further arguments in favour of the near independence of the results from the value of $T_{\rm swi}$ are given in Feldmeier (1993).

Most of the hot material in Fig. 5 is caused by *shell-shell* collisions. What is more, the temporal periodicity of hot and cold domains along the track of a shock is *not* caused by the thermal instability, but by these regular shell encounters. A possible explanation for the absence of the instability's oscillation is that the typical time (for a $T = 10,000$ s photospheric sound wave) between two shell encounters is a factor of ≈ 5 shorter than the period of this oscillation, so that the latter might not be seen because it is constantly disrupted by the collisions.

Besides the radiative cooling zones of shock heated material, there are a few broad regions of *thin* and *hot* material enclosed between two shells, namely in Fig. 5 starting at ($r = 3.2\ R_*$, $t = 51$ h) and at ($r = 4.5\ R_*$, $t = 50$ h). A possible explanation of this hot, thin gas is the following: In the course of two shells coming close to each other, the thin material between them may become *compressionally heated.* This enlarges its gas pressure, which possibly becomes large enough to drive the two shells apart again. The hot gas is then distributed between the shells where, because of its low density, it cannot cool efficiently, and therefore remains hot. This explanation is supported by the tracks of the shells in the density plot of Fig. 5, shortly

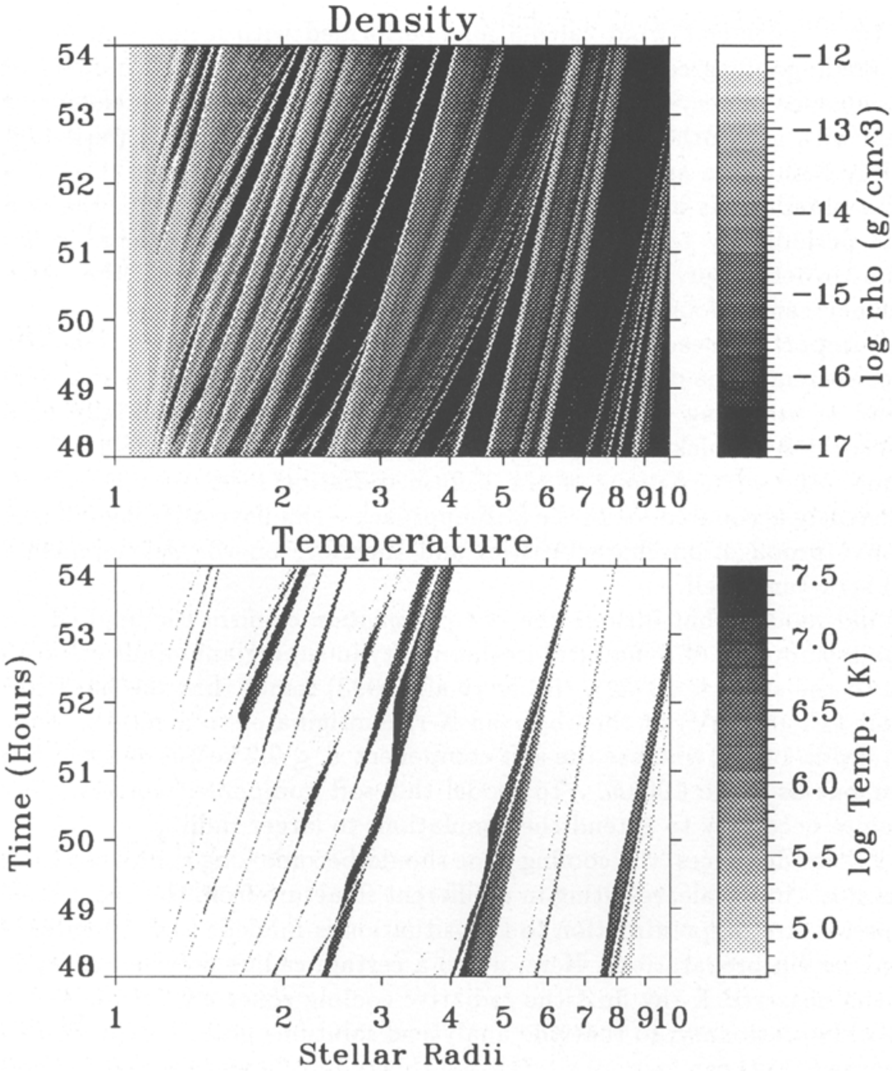

Fig. 5. Same as Fig. 4, but now for a cooling function with exponent $\alpha = 2$ instead of $\alpha = -1/2$ below $T_{swi}/T_{eff} = 10$.

before the hot, thin material shows up. These tracks show shell encounters where the inner shell is subsequently deflected again.

5. Conclusions and Future Work

Our time dependent wind calculations performed with a newly developed code are in good agreement with the results by S. Owocki. The wind consists of a sequence of very narrow, dense shells which show a strong tendency to collide with each other. (A speculation is, whether these encounters might possibly build up a *hierarchy*, cf. Hogan & Woods 1992.)

The simulations indicate that, at least up to $r \approx 2 \ldots 3\ R_*$, the wind reacts periodically to periodic base perturbations, with no clear signs of stochastic behaviour. As many as 50 harmonic overtone modes of the driving frequency can be found at these radii.

An important need at present is to develop a tractable way to calculate self-consistently the diffuse radiation field appropriate to a structured wind, instead of assuming a fixed, smooth source function. As originally found by Owocki & Rybicki (1985) and recently emphasized by J. Puls (cf. this volume), the *perturbations* of the diffuse radiation field – which are not consistently accounted for in the SSF approach – can have a strong influence on wave propagation characteristics and therefore on the development of wind structure itself.

Wind models that include the *energy equation* confirm the applicability of the assumption of dynamical isothermality in supergiant winds relatively close to the star ($\approx 10\ R_*$). Hillier et al. (1993) found that the hard component ($E \gtrsim 1$ keV) of the observed X-rays originates to a certain degree from inside 10 R_*, whereas the soft component $E \lesssim 0.2$ keV is emitted from as far out as 50 to 100 R_* . To model this soft component correctly, it is therefore necessary to extend the simulations to larger radii.

At these distances, the cooling time should become long compared to the dynamical time scale, resulting in a different structure from the one presented here. A first approximation to this situation is modelled by Kudritzki & Feldmeier (in preparation). Here, up to a certain radius determined by fitting the observed X-ray flux, the radiative cooling zones are assumed to be effectively stationary, so that the analytical solutions given by Chevalier & Imamura (1982) can be applied. Beyond this radius, however, radiative cooling is neglected, resulting in an adiabatically expanding and cooling shell, as was first analyzed by Simon & Axford (1966).

We intend to compute the X-ray emission from the time-dependent calculations in order to compare them with corresponding *ROSAT* observations. This work is just in preparation in our group, but examples of such comparison can be found in Cooper & Owocki (this volume).

Acknowledgements. It is a pleasure to thank J. Puls, S. P. Owocki, R. P. Kudritzki, A. W. A. Pauldrach, A. Fullerton, and U. Springmann for many discussions. Special thanks go to C. Reile who laid the foundation of the hydrodynamics code used here. This work was supported by the Max-

Planck-Institut für Astrophysik and the Deutsche Forschungsgemeinschaft under grant Pa 477/1-1. The calculations were performed on a Cray Y-MP at the Leibniz Rechenzentrum in Munich.

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DISCUSSION

Kaper: I would expect that the number of shell-shell collisions could depend on the period of your disturbances at the base of the wind. I mean, the shorter the distance between the shocks, the more collisions you would expect. So what happens to the X-ray production, then?

Feldmeier: That's right, yes, of course. This depends on the period of the photospheric sound waves, and it is the same second and third harmonic overtone mode which determines the frequency of the shell-shell collisions. So if you remember the calculations, we are seeing this pairwise and triple-wise creation of shells. The causes of this are the second and third overtone modes. So divide 5000 seconds by two or three, then you have the typical time scale between shell-shell collisions. Of course, this is at the moment a free parameter, yes.

Prinja: I have a related question to that, to some extent. You are getting a very impressive amount of structure there in a short time scale, and of course you're driving the wind terrifically in this 5000-second cycle. So, for example, if you wanted to get this picture that we've seen over the last few days of one very substantial structure per day–coherent structure per day–is it simply a case of just changing that one parameter to achieve that? Or would things become much more complicated? I mean, is it that straightforward? You just change the 5000-second cycle to achieve it?

Feldmeier: The structure should vary over, say, 12 hours. What's the question?

Owocki: Can you make one-day time scales by driving the base at one day?

Feldmeier: Yes.

Prinja: What would happen if you did that? Would you see....

Feldmeier: I haven't done the calculations, but I think nearly the same structure should be developed, yes.

Prinja: The same amount?

Owocki: Can I answer the question? The problem with sound waves is that they don't propagate at a frequency below the acoustic cut-off frequency. So the problem is you can't drive this with a very low-frequency wave and get anything to go up in the wind. That is, beyond the acoustic cut-off period, the wavelength of the sound waves becomes greater than the scale height of the atmosphere. Now in 1-D simulations you have no alternative to get waves up from the photosphere into the wind except through sound waves. It's the only wave you have, and you can't drive those more slowly. So, you have to do something, basically, more *ad hoc*. What I do is to put a perturbation already in the wind at 100 km/s. Those can make much larger scale structures that would be [at a longer time scale]. I would identify that with some sort of larger structure, like a magnetic field or something disturbing the wind.

Feldmeier: But this argument depends on the production of sound waves, on the frequency of creation of sound waves. But there are other mechanisms besides sound waves....

Owocki: Sure. If you had 2-D or 3-D then you could make gravity waves propagate up, or NRP–nonradial pulsations–or something could propagate and seed them, but [in 1-D all you have are sound waves].

Feldmeier: I should do the calculations, yes, but I haven't I done them up to now.

Heap: In your pictures showing shell-shell collisions you can see various shells merging, and sometimes it looks as if shells crossed. Might these apparent crossing be related to Lex Kaper's finding of two maximum DAC velocities in the wind of ζ Oph ?

Owocki: Feldmeier's plots show density versus space; Lex's observations show time versus frequency or velocity. The two aren't the same. You can have things crossed just because there are two structures in the wind that happen to go through the same velocity, but they're separated by a large distance. I would interpret what's happening basically as the higher overtones are creating all this higher structure, and there's structure coming off the shells as well. I wouldn't interpret this as crossing, but you can look at it how you want. The point I'm saying is that you see these crossing in synthetic spectra; I showed an example of that earlier and I'll show it again.

Cassinelli: Early on you showed a picture that looked very much like the OCR velocity structure. That was for the isothermal case?

Feldmeier: Right, yes.

Cassinelli: So after you included this energy equation, it looks quite different.

Feldmeier: I can show the plots again. That's the isothermal calculation out to five stellar radii, and here is the corresponding calculation with the switching temperature. Plotted are the density and velocity structures, and one sees again all the shell collisions. The structure is...

Cassinelli: It's not much different.

Owocki: It's qualitatively very similar.

Feldmeier: It's the same, but now the cooling zones are resolved.

Cassinelli: Now, if you did include conduction, do you think it would change?

Feldmeier: Maybe a little. As I know from the MacFarlane and Cassinelli (1989) paper, conduction should be important at temperatures higher than approximately 2×10^7 K. I think this limit is hardly reached here. But I will include it in the future.

Brown: On that point, this is a very extreme situation. It's not directly comparable, but in the solar atmosphere, the fact that you see these X-ray loops, and you don't see small ones, is not only because you don't have a good enough telescope. There are scaling laws, and if you make the loop too small, the conduction (because it goes like one over l squared where l is the loop length) will not allow high temperature loops to exist. It kills off any tendency to go to higher temperatures by a radiative instability. So I suspect that if you put conduction in, probably your smaller, very hot regions might tend to disappear and be smoothed out. I suspect that would be the effect. It would be interesting, though, to take a look at it.

Lamers: I understand that if you want to calculate this you need to start with some instability at the photosphere and you have to put in some period. So, isn't it possible in reality that you don't really need to start at the photosphere, but you can have just some very small random motions somewhere in the wind and they amplify, and then you can have all the time scales?

Feldmeier: That should be so, yes. I haven't done the calculations but I'm sure that if you don't start too far out in the wind, where you're close enough to the photosphere so that the instability is still quite large, you will see the resemblance. I think it doesn't depend on the photospheric sound waves–that it's just a simple way to control this developing structure–but rather it depends on the amplitude and the period.

Lamers:I have another question, if I may. I see a difference between your calculations and those of Stan. Your shocks develop at a slightly larger radius and a slightly higher velocity. Is the density gradient or the density amplitude that you start with the main determining factor?

Feldmeier: Right, yes. First of all, there's a regime of exponential growth of the instability, and the location in the wind where a given density contrast occurs is dependent on this assumed base perturbation. So when you take 25 percent base perturbation, you will reach this limiting value of the instability, say, quite quickly, and with a small amplitude it will take a longer time then, yes.

Owocki: Just a comment. The fact that that simulation was just done was in order to make lots of structures so that we could do the line synthesis. But, you know, it was an arbitrary number. We just put in a big perturbation to make lots of structure.

Lamers: Then it's important for how you use Hα observations to derive mass loss rates.

Owocki: Yes!

Lamers: That's important for [interpreting observations].

Owocki: Another crucial point that should be emphasized–that can't be emphasized enough–is that the details of what you assume for the smooth source function in that region also determine how much structure forms there. It's supposed to go to one-half [I_{star}] and thus stabilize [the base flow], but it doesn't quite go to a half. So you get instability right down to the photosphere and you won't have to drive the sound waves - in my simulations anyway. We have to reconcile this difference. In my simulations, I don't have to put in sound waves to get structure starting at about 1.6 stellar radii. It happens all by itself and never goes away. In his, he has to put it in. I don't know what the difference is.

Henriksen: Yes, all right. I like this work. I just want to say that I think this is establishing whether it goes to a chaotic situation. That's perhaps very important to the discussions. And it seems to me that you simply haven't run your calculations far enough yet. Your dynamics will still be affected by your driving frequency. And there are various ways to estimate this. There's this thing called velocity – which is essentially the frequency divided by something like the flow time. I wonder if you look at that it might tell you where you stand.

Feldmeier: So, what's your opinion? Do these calculations have to be done further in time?

Henriksen: Yes.

Feldmeier: I think nothing will happen, because as well, this Fourier time signal was, I think, almost for 83 complete cycles.

Henriksen: You're seeing frequency doubling, am I correct?

Feldmeier: I haven't shown the frequency spectrum of the subharmonics but there is clearly no sign of period doubling. There is a little peak there for frequency. I don't know what it is. But it's certainly not at half the frequency.

Henriksen: There are two things to look at. One is this Rossby number, and secondly the range of scale.

Feldmeier: But the critical number for these calculations must depend in some sense on the radiation field. This is not included in the Rossby number.

Henriksen: I wonder if you have done this analysis on the adiabatic calculation.

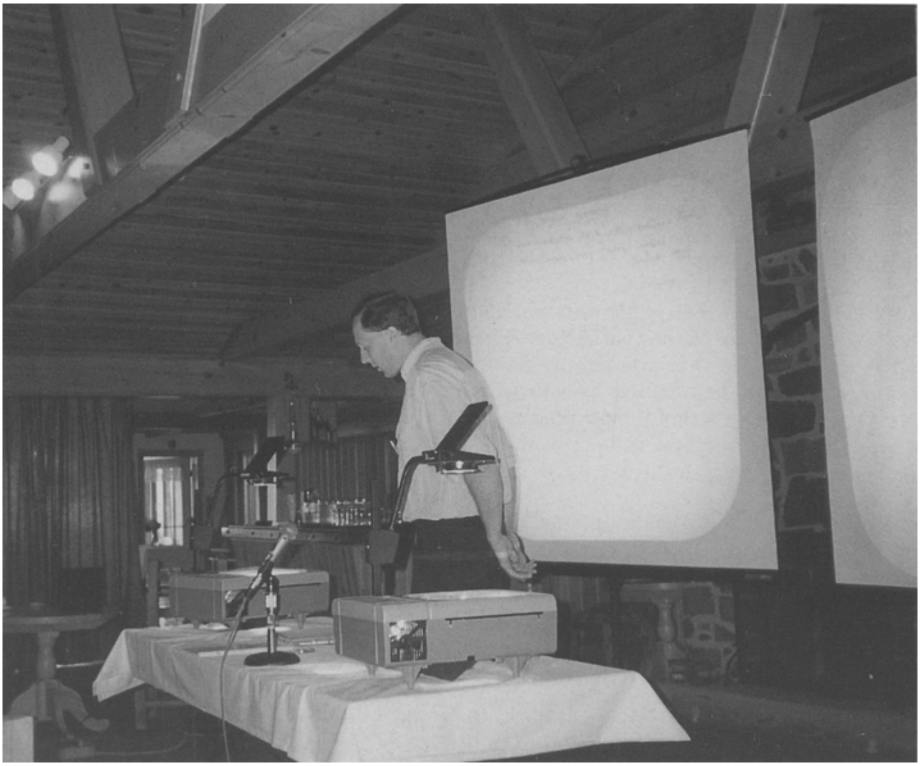
Feldmeier: Not yet. But I think it will not show a very different power spectrum. But I have to do it.

Cherepashchuk: Do you include as a cooling factor Compton scattering of optical photons by hot electrons? (The optical radiation field is very strong in the star and this is a strong cooling factor.)

Feldmeier: No.

Cherepashchuk: It's very important. [In a] Wolf Rayet plus O binary system, [it] keeps the field of radiation an order of magnitude [cooler?] by the flow of optical photons through the shockwave.

Feldmeier: There is only a parameterized cooling function at the moment, and, true, one should include a physical cooling function.



Achim Feldmeier