Can the line-driven instability form BAL QSO clouds?

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Abstract. We discuss whether a radiation-hydrodynamic instability of line-driven winds can lead to the formation of BAL QSO clouds. We find that, (1) under conditions as in BAL regions, the instability-generated clouds have filling factors of $\geq 10^{-4}$. This should be too large to keep the gas at the observed low ionisation stages. (2) A large fraction of the gas is compressed into the clouds only at some height in the wind, via collisions with tiny clouds. Shock confinement is therefore efficient in this model.

1. Introduction

Weymann et al. (1985) deduced volume filling factors for the gas which forms absorption troughs in BAL QSOs as small as 10^{-6} or 10^{-7} (namely, $\epsilon \approx 10^{-7}$ $N_{20}\, L_{46}^{-1}\, R_{\rm pc}\, T_{4.5}$). Such tiny clouds must be externally confined to hinder them from quickly expanding due to internal pressure.

Some of the recent ideas proposed to overcome the problems connected with any such *massive* confining medium include: (1) to reduce the number of ionizing photons by X-ray absorption in dense gas located between the central source and the BAL region (Murray et al. 1995; cf. also Shlosman et al. 1985); (2) to identify the BAL absorbers with trails of dust intermixed with gas, where the dust stems from evolved mass-loss stars in the nuclear star cluster (Scoville & Norman 1995); (3) to postulate a massless confining medium like magnetic fields (Rees 1987; de Kool & Begelman 1995) or cosmic rays (Begelman et al. 1991).

As a further possible alternative, we discuss here a hydrodynamic instability (Lucy & Solomon 1970; Owocki & Rybicki 1984) which may generate and confine clouds in line-driven flows. This instability and the flow structures it causes are the subject of many recent calculational and observational efforts concerning O star winds.

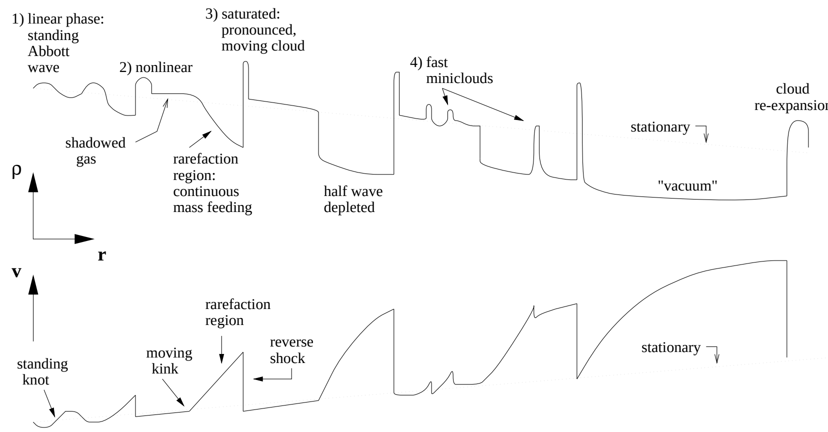

Figure 1. The radial wind structure which develops from the instability.

2. The line-driven instability

If a fluid parcel in a line-driven stellar wind experiences a positive, radial velocity perturbation, the parcel is Doppler-shifted out of the absorption shadow of gas lying closer to the star. As this raises the line force, the parcel is accelerated further, i.e., the original perturbation is *amplified* (Lucy & Solomon 1970).

Since velocity perturbations perpendicular to the incident photon flux are damped away (Rybicki et al. 1990), the instability-generated structure should be mainly radial. Fig. 1 shows a sketch of the wind structure as found in numerical O star wind simulations (Owocki et al. 1988; Feldmeier et al. 1997). Here, a sequence of dense and narrow clouds develops in the flow (or actually shells in present 1-D simulations). The clouds are confined by a strong reverse shock on their inward side, which decelerates and compresses streams of fast, inner gas. At some height in the wind, this continuous mass feeding ceases; however, somewhat later small cloudlets start to form in the remaining intercloud gas, which are accelerated by the line force through the void space ahead of them until they collide with the next outer pronounced cloud. This gives rise to a second stage of gas compression and confinement.

In the following we assume a flow geometry of purely radial wind rays (cf. Arav & Li 1994; Arav et al. 1994), and frequently use the velocity law of Castor et al. (1975; CAK in the following) for a smooth wind from a point source of radiation, $v_0(r) = v_\infty \sqrt{1 - r_s/r}$, with r_s the subsonic start-off radius of the ray, and v_∞ the terminal speed. However, the estimate for the overdensity should not be too different if instead, e.g., the disk wind model of Murray et al. (1995) is used. For these authors notice that the streamlines rapidly approach radial lines, since in most wind regions $v_r \gg v_\phi \gg -v_\theta$. Here r_s is the start-off radius from the disk.

3. Estimate of maximum overdensity

The gas compression takes place in strong reverse shocks terminating rarefaction regions of highly accelerated, thin gas. To find an upper estimate for the overdensity behind the shock, o, we assume radiative cooling to be efficient, so that the shocks can be viewed as being effectively isothermal (cf. below). Shock properties like velocity jumps, etc., can then be found from examining the rarefaction regions. The latter occur due to the instability steepening the thermal band until the band becomes optically thin ("de-shadowing", cf. above). The large line force acting on this optically thin gas can be determined as follows. Let $g_{t,\kappa} \equiv \kappa F/c$ be the force-per-mass due to an optically thin line of strength κ, and $N(\kappa)$ the line distribution function. Then,

$$C \equiv \frac{g_{\rm thin}}{g_{\rm line}} = \frac{\int d\kappa\, N(\kappa)\, g_{t,\kappa}}{\int d\kappa\, N(\kappa)\, g_{t,\kappa}\,(1-e^{-\tau})/\tau}. \tag{1}$$

Since the contribution from optically thick lines to the total force is only slightly larger than that from optically thin lines (CAK), we perform the integral in the denominator only for $\tau \leq 1$, hence $(1-e^{-\tau})/\tau \approx 1$. Using the CAK distribution function, $N(\kappa) = \kappa_0^{1-\alpha}/\kappa^{2-\alpha}$, gives

$$C \lessapprox \left(\frac{\kappa_m}{\kappa_1}\right)^{\alpha} = \tau_m^{\alpha}, \tag{2}$$

where $\tau(\kappa_1) = 1$, and κ_m is for the strongest line present. For a standard value of $\alpha \approx 2/3$, this implies $C < 10$ for BAL QSO flow conditions. Using finally the CAK velocity law in $g_{\rm line} \equiv v_0 v_0'$,

$$g_{\rm thin}(r) = \frac{C\, r_s\, v_\infty^2}{2r^2}. \tag{3}$$

We assume now that the rarefaction region is narrow, $h \ll r$, and so is nearly planar, and that it is roughly stationary in an *inertial* frame moving with at the smooth wind speed $v_0(r)$. The rarefaction starts off from (nearly) smooth wind, v_0 and ρ_0 (see the discussion of "Abbott knots" below). From momentum conservation, we take the cloud speed to be equal to $v_0(r)$. From (3), the jump velocity at the reverse shock is then,

$$v_j(r) = \frac{\sqrt{C\, r_s\, h}}{r}\, v_\infty. \tag{4}$$

Taking a fixed mass flux, the stretching of the accelerating fluid parcels gives the rarefaction ratio of preshock density to smooth wind density,

$$\frac{\rho_{\rm pre}}{\rho_0} = \frac{1}{\sqrt{C}}. \tag{5}$$

Applying the jump conditions at a strong isothermal shock gives then for the cloud overdensity (with a the isothermal sound speed),

$$o \equiv \frac{\rho_{\rm cloud}}{\rho_0} = \sqrt{C}\, \frac{r_s\, h}{r^2} \left(\frac{v_\infty}{a}\right)^2. \tag{6}$$

To achieve large overdensities, one should therefore seek to maximize h. However, the unstable growth-rate drops with the perturbation wavelength, λ (where $h \approx \lambda$), and one has to ascertain that the growth time of such a long-scale perturbation is still short against the flow time. From Eq. (1) of Owocki (1994), the growth-rate for $\lambda \gg v_{\rm th}/dv/dr$ is,

$$\Omega = \pi^2 v_{\rm th} \, \frac{v}{dv/dr} \, \frac{1}{\lambda^2}. \tag{7}$$

We assume that perturbations of lengthscale λ_s and amplitude $O(10^{-1})$ relative to the mean flow (e.g., sound waves or turbulence) are introduced at the subsonic wind base in order to have enough time to grow. In general, this should also make sense in the framework of the model of Murray et al. (1995). There one encounters first the instability of the vertical wind driven by the disk's UV radiation field, and starting off from subsonic velocities; later the streamlines bend over into radial directions due to illumination by the central source, and the perturbations grow further due to the line-driven instability caused by this latter radiation field.

The wavelength of a perturbation stretches according to $\lambda(r) = \lambda_s \, v(r)/a$ in the accelerating wind, and its amplitude grows $\propto \exp(\int \Omega \, dt)$. The maximum λ_s that is still able to grow from the base amplitude at r_s to saturation at some radius r is then given by

$$\int_{r_s}^{r} \frac{\Omega}{v} \, dr \gtrsim 10. \tag{8}$$

If we assume that pronounced BAL clouds are formed by the time half the terminal velocity is reached, then $r \approx 1.5 \, r_s$. Using BAL QSO parameters $a = 20$ km/s, $v_\infty = 30,000$ km/s, this implies $\lambda_s \lesssim 2 \times 10^{-5} \, r_s$, which, when inserted into (6), and assuming $C = 10$, gives

$$o = 10^5 \text{ to } 10^4 \quad \text{for} \quad r = 1.5 \text{ to } 5 \, r_s, \tag{9}$$

where the latter radii span the dynamically most interesting region. The corresponding cloud filling factors, $\sim 1/o$, are therefore one or two orders of magnitude *larger* than those derived by Weymann et al. (1985).

Cloud re-expansion. Even then, (9) is an *over*estimate since, after initial shock compression, the clouds will re-expand due to internal pressure. This can be estimated as follows. Since the internal sound crossing time is much smaller than the flow time, the clouds are always in hydrostatic equilibrium. According to the equivalence principle, the line acceleration of the clouds can be replaced by gravity pointing in the opposite direction, acting on a static 'atmosphere' then. The scale height of the latter is a measure for the cloud size,

$$H = \frac{a^2}{g_{\rm line}} = \frac{2r^2}{r_s} \left(\frac{a}{v_\infty} \right)^2, \tag{10}$$

where again the CAK velocity law has been used. By assuming that all the gas from within a rarefaction region of extent h is compressed into the cloud, it follows that

$$o \approx \frac{h}{H} = \frac{r_s \, h}{2r^2} \left(\frac{v_\infty}{a} \right)^2. \tag{11}$$

Comparing this with (6), we see that cloud re-expansion reduces the density achieved by shock compression by a factor of $2\sqrt{C} \approx 5$, i.e., the filling factors are now *two or three* orders of magnitude larger than those derived by Weymann et al. (1985). Actually, they are now even sufficiently large that radiative cooling could be efficient enough to allow for such compressions.

Yet, the above argument depends crucially on the assumption that the acceleration of clouds is effective over the whole dynamical region. One may wonder whether instead *almost all the wind mass* is compressed into the clouds already at low radii, e.g., up to $r = 1.5\,r_s$. Besides a loss of cloud acceleration and confinement, this would also affect the overdensity directly, since outer rarefaction regions would no longer start off from smooth wind conditions but instead from highly depleted gas, implying low densities at the reverse shocks.

We argue in the next section why this should *not* happen, and that instead the above estimate is essentially correct.

4. A closer look at cloud formation

The occurence of steep rarefaction regions that terminate in strong reverse shocks can be viewed as unstable growth of long-scale, radiative-acoustic waves. Abbott (1980) showed that such waves propagate inwards at a velocity of $\approx -v_0$ (neglecting terms of order the sound speed), i.e., they stand with respect to the central source.

Consider then those *knots* of an inward propagating wave (i.e., density and velocity are in anti-phase) at which the velocity gradient is *steeper* than in the smooth flow. From Fig. 1 one sees that these knots are the footpoints of rarefaction regions. Unstable growth steepens then the velocity gradient to larger radii of the knot; whereas the gradient to smaller radii remains at roughly its stationary value, since shadowing effects hinder unstable growth there. This opposite behavior transforms the *knot* into a *kink*, i.e., short-scale Fourier modes develop there.

The signal speed of inward propagating short-scale perturbations is however the *sound* speed again (Owocki & Rybicki 1986), as opposed to being the *wind* speed for long-scale radiative-acoustic waves. The nonlinear, comoving kink, as opposed to the standing knot, is therefore "shut" for further mass and momentum transfer into the rarefaction region and the cloud, and, in the limit $a \to 0$, a contact discontinuity may develop at the kink. Notice that, for λ_s from the foregoing section, the nonlinear kink occurs below or around $1.5\,r_s$.

As indicated before, numerical simulations show how this gas reservoir below the kink is used up at larger heights in the wind in a second compressional phase. Tiny cloudlets which remain at roughly stationary wind densities form in this gas, and penetrate through the kink. The cloudlets are optically thin (however, approaching $\tau_m = 1$), wherefore they are strongly accelerated through the empty rarefaction region by the thin line force (3), until they finally collide with the next outer, *overdense* cloud. (The X-ray flash from these collisions was suggested as origin for the observed X-ray emission from O stars; Feldmeier et al. 1997). In simulations of O star winds, the cloudlets reach out as far as $\gtrsim 5\,r_s$, by which radii the flow has essentially reached v_∞.

Summary. We conclude that the line-driven instability in a radial BAL flow may compress the wind gas into clouds of filling factors 10^{-3} or 10^{-4}. This falls short by two or three orders of magnitude to explain the filling factors of Weymann et al. (1985). A promising direction for future work seems therefore to consider if a *combination* of the present approach with alternative proposals to overcome the confinement problem (Murray et al. 1995; Arav et al. 1994) may allow to relax some of the strong assumptions of the latter models.

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Discussion

Arav: Are the overdense clouds driven by the line force?

Feldmeier: Not directly, since the clouds are very optically thick, and the incident flux is blocked at their surface. However, the clouds gain momentum from the radiatively driven gas which is compressed into them.

Voit: These instabilities arise from line-of-sight shielding. Should we expect them to be coherent perpendicular to the line of sight?

Feldmeier: Rybicki et al. (1990) showed that the flow is *not* subject to the line-driven instability in directions perpendicular to the incident flux. However, the Kelvin-Helmholtz instability, etc., should generate 3-D structure.