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The condensation of dust around η Carinae

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Summary. A study is made of the present condensation process of solid material around η Car. We show that this high luminosity ($5 \times 10^6 L_\odot$) star with a mass of $160 M_\odot$ is losing mass since 1840 at the rate $\dot{M} = 7.5 \times 10^{-2} M_\odot$/yr. To drive this mass flow, a power is required of $8 \times 10^6 L_\odot$, which may account for the intrinsic fading since 1840 by 1 mag. Since 1856 dust has been condensing from the circumstellar gas at distances in excess of $R_i = 3.7 \times 10^{14}$ m from the star, initially at the rate $\dot{M}_\mathrm{d} = 1 \times 10^{-4} M_\odot$/yr and presently at a somewhat higher rate. This follows from an interpretation of the light curve with a model of the time evolution of the dust envelope, which accounts for the wavelength-dependent extent of the infrared source. The average density at R_i of about 1×10^{12} m^{-3} falls short by orders of magnitude to explain the condensation of solid material. The possibility is discussed of strong inhomogeneities in the density, which are stabilized against turbulent perturbations by a stellar magnetic field. Such a field would have a strength that is normal for Ap and Am stars. We furthermore give chemical arguments on the nature of the condensate and conclude that it consists in part of disordered silicate clusters. Using the pertinent optical properties and a model of the temperature stratification of the dust envelope, we fit the observed infrared spectrum for clusters with a size of about 1 μm. Such large grains give an almost grey circumstellar extinction of 3–4 mag in the visual and ultraviolet, which fits well in an independent interpretation of the extinction of η Car.

1 Introduction

Only few observations pertain to the nucleation and growth of solid particles from cosmic gas clouds. Such processes have been important in the early history of the solar system, where they form 'one of the most obscure chapters' (Reeves 1972). Cloud parameters like abundance, temperature, pressure and the state of ionization, play a role in determining the chemical nature of the condensate. It is useful, therefore, to scrutinize the few documented astronomical events in which solid particles have been formed.

The present paper discusses the case of η Car. After its eruption of 1840 this star is losing gas, from which by now a thick envelope of dust has been condensed. This view is based on the existence of an expanding envelope of gas with a very complex structure called *homunculus* (Gaviola 1950). Irregular features move at high speed from the centre of this nebula, whereas no direct light from the star can be seen. After 1856 the star became fainter as a result of the formation of solid material in the *homunculus*, which is heated by the star and makes the object one of the brighest infrared sources in the sky (Neugebauer & Westphal 1968). Obviously, the light curve contains important clues to the condensation process.

In Section 3, this curve will be interpreted, using spectral data on the gas and infrared data on the dust, in terms of a condensation rate. The information on gas densities in the envelope is fitted into the picture of a continuous mass flow (Section 4). Then, in Section 5, temperatures and pressures in the gas cloud are related to the actual condensation. Chemical arguments will be discussed (Section 6) predicting the nature of the condensed particles and the main arguments and conclusions are resumed in the final section. First, however, we have to discuss some basic data.

2 Basic data

The star η Car is located near the centre of the OB-association Trumpler 16 and therefore it probably belongs to it. Photometric studies of the association have yielded distances of 2.5 kpc (Thé & Vleeming 1971) and 3.4 kpc (Feinstein, Marraco & Muzzio 1973). Observed motions in the gas cloud surrounding the star suggest a nova-distance of about 2.3 kpc (Gehrz *et al.* 1973), so let us adopt the smaller distance of 2.5 kpc.

The spectrally integrated energy flux is about $2.7 \times 10^{-8}\,\mathrm{W\,m^{-2}}$, the main part being emitted in the infrared, namely $2.6 \times 10^{-8}\,\mathrm{W\,m^{-2}}$ (Sutton, Becklin & Neugebauer 1974). At the above distance this amounts to a luminosity $L_* = 1.9 \times 10^{33}$ W ($5 \times 10^6\, L_\odot$).

The ultraviolet flux seems to peak as if the star has a temperature of 1.5×10^4 K, but this value cannot describe the overall spectrum well (Pottasch, Wesselius & van Duinen 1976). A higher temperature of about 3×10^4 K has been derived by Davidson (1971) from the Hβ luminosity. If the reasonable assumption is made that the gas cloud is ionized by stellar far-ultraviolet photons, the presence of emission lines of Fe III, N III and Ne II requires a temperature of 3×10^4 K or higher indeed. In a star with an extended atmosphere as is the case of η Car, the brightness temperature at the ionization edges of these elements differs from the effective stellar temperature (Castor 1974). Therefore a direct estimate of T_* from the line spectrum is difficult. Let us nevertheless adopt $T_* = 3 \times 10^4$ K.

The mass of the star can be estimated from the binding condition that $M_* > \mu L_*/(4\pi c G)$. Here μ is the mass absorption coefficient in the stellar atmosphere, which for the above T_* is essentially given by electron scattering: $\mu = 4 \times 10^{-2}\,\mathrm{m^2\,kg^{-1}}$; G is the gravitational constant. This gives $M_* > 2.9 \times 10^{32}$ kg ($145\, M_\odot$). In the case of a considerable mass loss (see Section 4), the ratio of the radiative to the gravitational acceleration is probably close to unity (*cf* the extended atmosphere models discussed by Cassinelli 1971 and Castor 1974). We thus estimate M_* to be some 10 per cent heavier than the above limit or $M_* = 3.2 \times 10^{32}$ kg ($160\, M_\odot$). The same value is obtained by extrapolating the mass–luminosity relation for heavy stars (Stothers 1972) to the absolute bolometric magnitude -12.0 of η Car. If the star has lost 6 per cent of its mass by now (Section 4), this relation suggests an average intrinsic fading by 7×10^{-4} mag/yr only.

Table 1. Basic data on the star and its present extinction.

		η Car
distance		2.5 kpc
luminosity		5×10^6 $L_\odot$
effective temperature		3×10^4 K
(radiative) mass		160 $M_\odot$
(radiative) radius		80 $R_\odot$
surface gravity		6.8 m s^{-2}
escape velocity		3×10^5 m s^{-1}
terminal velocity		9×10^5 m s^{-1}
E_{B-V}	interstellar	0.4–0.5 mag
	circumstellar	0.7 mag
A_V	interstellar	1.2–1.5 mag
	circumstellar	3.8 ± 1 mag
mass envelope		10 $M_\odot$

The above luminosity and temperature define a radius $R_* = 5.6 \times 10^{10}$ m ($80\,R_\odot$) which, together with the mass estimate, leads to a surface gravity $GM_*/R^2 = 6.8\,\mathrm{m\,s^{-2}}$ and a gravitational escape velocity $v = (2GM_*/R_*)^{1/2} = 8.8 \times 10^5$ m/s. The true escape velocity v^* is reduced by radiation pressure and becomes $v^* = v\,[1 - \mu L_*/(4\pi c G M_*)]^{1/2} \simeq 3 \times 10^5$ m/s. The observed gas velocities are larger (see, further, Section 4). Table 1 lists the above data and the main results of the following discussion of the extinction.

The colour excess E_{B-V} of η Car can be obtained from the intensity of emission lines, not affected by self absorption, in different spectral regions. A value of $E_{B-V} = 1.2$ mag was derived by Pagel (1969) and Viotti (1969) from [Fe II] lines, whereas Lambert (1969) suggested an even higher value. This conclusion has been called a mistake by Pottasch *et al.* (1976), who discussed the absorption in the 2200 Å band towards η Car. They derive that $E_{B-V} = 0.5$ mag, which refers to truly interstellar extinction (*cf.* Feinstein *et al.* 1973), and suggest a difference between the interstellar dust and the circumstellar material. In discussing the controversy, Leibowitz (1977) has confirmed the conclusion of Pagel and Viotti and has referred to the case of θ^2 Ori, another star for which the 2200 Å band is 'too weak'. There is no controversy if we attribute $E_{B-V} = 0.7$ mag to circumstellar dust and show that this material fails to absorb selectively at 2200 Å.

In order to do so we summarize an argument involving the grain size, which has been given in full detail by Andriesse (1977). To produce an absorption band, a solid particle should not be larger than $1/\kappa$, where κ is the absorption coefficient. If it is much larger, the particle is opaque, so that the absorption cross-section is equal to the geometrical cross-section and, hence, independent of the wavelength. If it is much smaller, the particle is transparent, so that the absorption cross-section is proportional to κ and, hence, dependent on the wavelength in general. It can be shown that in the maximum of an absorption band $\kappa = \Omega^2/(c\gamma)$, *irrespective* of the chemical nature of the particles. Here Ω is a typical plasma resonance frequency of valence electrons ($\simeq 10^{16}\,\mathrm{s^{-1}}$) and γ the damping in this electron system ($\simeq 0.1\,\Omega$); c is the speed of light. One finds $\kappa = 3 \times 10^8\,\mathrm{m^{-1}}$.

Thus, if we require complete transparency, only particles smaller than $1/\kappa \simeq 3 \times 10^{-9}$ m are able to produce the 2200 Å band. If we allow for some absorption in the particles, we can relax the limit to about 1×10^{-8} m. The case of such particles, much smaller than the common grains, has found support in recent investigations of the *interstellar* medium (Witt 1977; Andriesse 1978). However, it is unlikely that they exist in an environment which obviously favours grain growth. As such we also characterize *circumstellar* clouds with their relatively high gas density. We thus expect that the dust grains around η Car are much larger than 10^{-8} m, which makes them opaque to the ultraviolet around 2200 Å. In Section 6 their size will be determined to be ~ 1 μm.

The total circumstellar extinction A_V^{cs} can now be estimated from the circumstellar E_{B-V}. We adopt the ratio (5.4 ± 1.5) for the total to selective extinction, which is suggested by the slope of the recent trajectory of η Car with its nebula in the ($B-V$, V) diagram (Thé, private communication) assuming unchanged stellar colours. Thus $A_V^{\mathrm{cs}} = (3.8 \pm 1)$ mag. In view of the argument below, we believe that the true value is 3.0 rather than 3.8 mag. Aitken *et al.* (1977) recently obtained $A_B = 5.5$ mag from an analysis of the Brackett α line. Since A_B^{is} is about $(3+1)\ 0.45 = 1.8$ mag, A_B^{cs} would be 3.7 mag, which is circumstellar E_{B-V} more than the above A_V^{cs}. Pecker (1972) has suggested that one can estimate the circumstellar extinction in the ultraviolet from the ratio of infrared to ultraviolet radiation. His procedure gives about 4. The data thus yield the picture of an extinction which increases towards shorter wavelengths. Grains of 1 μm will not absorb better when the wavelength goes from V to B, so that the increasing extinction has to be ascribed to a scattering effect. The case must be intermediate between pure Rayleigh scattering, which heavily reddens, and pure Tyndall scattering, which does not redden. In fact, for an average grain size of twice the yellow wavelength, there is still a tendency for the scattering cross-section to increase from V to B (e.g. van de Hulst 1957) and this tendency is amplified by the multiplicity of the scattering in the optically-thick *homunculus*.

3 Light curve

The visual light curve of η Car is shown in Fig. 1, which is based on the observations by Innes (1903), O'Connell (1956) and Feinstein & Marraco (1974). The dot represents the present bolometric magnitude based on infrared data (Sutton *et al.* 1974). Remarkable are the light maxima during the bright phase around 1843 and the maximum around 1888. The spectral evolution around 1888, rediscussed by Witney (1952), is nova-like. Even more

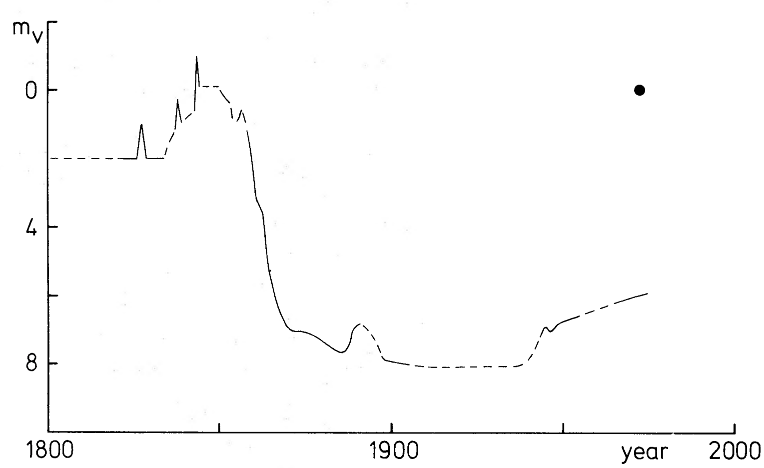

Figure 1. Visual magnitudes of η Car in the past and the present bolometric magnitude based on infrared data (filled circle). The full parts of this light curve are well documented, the dashed parts are uncertain.

remarkable is the steady fading of the star from 0.5 mag in 1856 to 6.9 mag in 1868. The average rate is 0.57 mag/yr. Similar, though faster, fadings have been observed in other peculiar stars, like RCrB. The decrease of brightness has to be ascribed to the condensation of dust, which started to obscure the star from 1856 on.

Convincing evidence for this is the strong circumstellar infrared emission discovered by Neugebauer & Westphal (1968). The spectrum of this emission is continuous, with the suggestion of a silicate feature (Robinson, Hyland & Thomas 1973), which proves that it is caused by solid particles. Combining the two ideas of a nova with mass ejection and dust formation in the ejecta, Davidson (1971) proposed the following scenario. In 1840 'some instability caused the sudden ejection of material'; when this material 'had expanded to the point where most of its radiation was at visual wavelengths' η Car was brightest; 'later, it seemed to fade as it became yet cooler and as dust grains began to form' (1856). By absorbing almost all stellar photons the dust is heated, so that the bolometric magnitude of 0.0, given by the infrared radiation, measures the present photon flux from the star.

It is instructive to compare this 0 mag with the past magnitudes of the star. We expect a total visual extinction due to interstellar dust of 1.2 to 1.5 mag (*cf.* Section 2). At its bright phase around 1843 the star thus had a bolometric magnitude of -1.0 to -1.5, as in Davidson's scenario the bolometric correction is zero here. Before 1840 we need to apply a bolometric correction of about -2 mag, if the stellar temperature at that time was comparable to its present value. This again suggests a bolometric magnitude of about -1. Thus the star is fainter now by about 1 mag. We shall relate this factor to the mechanical power required to eject gas at the rate discussed in the next section. If the star is fainter than before 1840 by about 1 mag, the present magnitude of about 6 suggests a *circumstellar* extinction of $A_V^{cs} \simeq (6-2-1) = 3$ mag. This is close to the value (3.8 ± 1) mag, estimated from the present colour excess.

The consistency of the above budget calculations encourages us to interpret the difference between the varying visual magnitude and the visual magnitude before 1840, decreased by 1 mag, as the time-dependent optical depth $\tau(t)$ of the circumstellar dust (except for the factor 1.086). We shall now relate $\tau(t)$ to the condensation of the dust using a model of the time-dependent geometry of the dust envelope.

The ingredients for this model can be obtained from the present distribution of the dust, suggested by infrared photometry, and from the gas dynamics. Andriesse (1976) found that, for distances R in excess of R_i, the mass density of the dust is

$$m(R)/m(R_i) = (R/R_i)^{\alpha} \quad (1)$$

with $R_i \simeq 3.7 \times 10^{14}$ m and $\alpha = 1.5$. Presently, $m(R_i) \sim 10^{-20}$ kg m^{-3}. (The model referred to has been worked out for a distance to the star of 3.4 kpc. The ultraviolet and infrared energy fluxes scale in the same way when another distance is adopted. Therefore the calculated temperature profile in the envelope remains applicable.) For distances R smaller than R_i, the dust would be too hot to exist and the gas too hot to allow for the condensation of dust.

The total mass of dust is

$$M_{\mathrm{d}} = 4\pi \int_{R_i}^{R_0} dR\, R^2\, m(R), \quad (2)$$

where R_0 is the outer radius of the envelope. On the other hand we have for the visual optical depth

$$\tau = \pi a^2 Q \int_{R_i}^{R_0} dR \frac{m(R)}{4/3\ \pi a^3 \rho}, \quad (3)$$

where Q is the visual extinction efficiency ($\simeq a\omega/c$ with ω the visual frequency of 3.4 : $10^{15}\ s^{-1}$ if the grains are small compared to c/ω, and $\simeq 2$ if they are much larger tha c/ω), a the average radius of a dust particle and ρ its average mass density of 2×10^3 k m^{-3}.

Using (1) we find the following relation between the dust mass and the visual optica depth

$$\frac{M_d}{\tau} = \frac{16\pi}{3} \frac{\alpha+1}{\alpha+3} \frac{\rho c}{\omega} \frac{R_0^{\alpha+3} - R_i^{\alpha+3}}{R_0^{\alpha+1} - R_i^{\alpha+1}}, \qquad (4$$

provided $a < c/\omega$. In the other case the factor $\rho c/\omega$, which is the inverse opacity, has to b replaced by $a\rho/2$. If the grains are 1 μm in size, this factor is about 5 times larger. Thu equation (4) gives a lower limit to M_d/τ. If a power law like (1) has been relevant for a times since 1856, we can determine M_d as a function of time from $\tau(t)$, provided the tim evolution of R_0 is known. The simplest assumption on R_0 is

$$R_0 = R_i + vt \qquad (5$$

with t the time elapsed since 1856 and v the velocity of expansion of the envelope of abou 9×10^5 m/s (see next section).

For small times equations (4) and (5) can be expanded and give

$$\dot{M}_d = \dot{\tau}\, \frac{16\pi}{3} \frac{\rho c}{\omega} R_i^2. \qquad (6$$

Inserting for $\dot{\tau} = 1.75 \times 10^{-8}\ s^{-1}$, which refers to the initial fading of 0.57 mag/yr, we fin $\dot{M}_d = 7 \times 10^{18}$ kg/s. Initially, the grain size should have satisfied the condition $a < c/\omega$, s this is a true rate rather than a lower limit.

For long times we can evaluate (4) with (5) and find the increasing curve shown in Fi 2. This curve depends on α, which here is taken to be 1.5 in accordance with the infrare profile of the envelope. We now derive $\tau(t)$ from the light curve as indicated above an multiply this quantity with the curve for M_d/τ. The result for M_d as a function of time shown in Fig. 2 also. Ignoring the dashed (uncertain) parts of the curve for M_d, we find

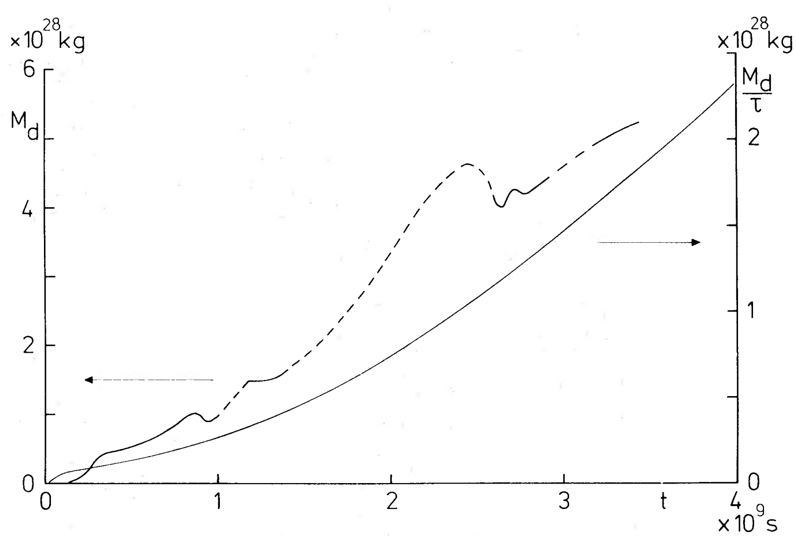

Figure 2. Time evolution of the ratio of total dust mass to visual extinction M_d/τ given by equation (: and the derived increase of the total dust mass M_d (see text).

practically *linear* relation with the time. The average increase of the dust mass $\dot{M}_d$, which is the condensation rate, is 1.8×10^{19} kg/s. One should take into account that this is a lower limit, because at present the grain size may exceed c/ω. Being of the same order of magnitude as the initial rate, it suggests that the dust has been condensing continuously since 1856 at an almost constant, or perhaps slightly increasing, rate of the order of

$$\dot{M}_d \simeq 1 \times 10^{19} \text{ kg/s } (1.5 \times 10^{-4} M_\odot/\text{yr}). \quad (7)$$

4 Mass flow

In this section we discuss indirect arguments for a continuous mass flow from the star since 1840.

The first argument is the improbability of the opposite, namely that the mass is lost in one short period only. Presently, the total dust mass given by (7) is 4×10^{28} kg. The ejected amount of gas has to be then at least 4×10^{30} kg, in view of the cosmic abundances (see also below). To lift off this mass from η Car, one needs $> 2 \times 10^{42}$ J of mechanical energy. This energy cannot possibly have been produced as such without the simultaneous production of about as much radiative energy. Such an energy pertains to a supernova. But the light curve of η Car is completely different from that of a supernova. Furthermore, the gas should form a shell around the star, which looks very different from the actual nebula with its bright core.

Dismissing the loss in one short event as very improbable, we can for another reason exclude that the mass is lost in a small number of short periods, associated with the maxima in the light curve. The energy of these less energetic events is still $\sim 10^{42}$ J, which would point to a quick recurrence of supernova-like instabilities. These are hard to conceive in the context of explosive nucleosynthesis. It is difficult, also, to explain how the outflow can be suddenly stopped after these instabilities. The Kelvin–Helmholtz relaxation time of the star is about 1300 yr. This leads us to believe that the expulsion of gas is in a smooth flow rather than in one or a few discrete lumps. Further indirect arguments are the similarity with the gas dynamics of P Cyg and the gas density in the immediate neighbourhood of the star. They will be discussed below in the context of an assumed constant outflow.

We turn back to the condensation rate. The dust is a condensate of the fraction of atoms in the circumstellar gas cloud that are able to form chemical bonds. Probably these atoms are mainly Mg, Si and Fe on the one hand, forming bonds with C, N and O on the other hand. The cosmic mass fraction of these atoms sums up to about 1 per cent of H. (The abundances in the gas envelope of η Car may deviate significantly from the cosmic values, as a large mass loss brings the products of the CNO cycle to the surface of the star (de Loore, De Grève & Vanbeveren 1977), which should increase the relative abundance of the metals). A constant mass flow from the star should thus be at least 100 times larger than the condensation rate. This gives $\dot{M} > 1 \times 10^{21}$ kg/s. Actually it may be larger by an order of magnitude, to allow for inefficiencies in the condensation process.

Let us now adopt that, on the average, the nuclear energy production in the star is unchanged during the whole period considered. The fact that the star is now intrinsically 1 mag fainter than before 1840 implies then, that only 40 per cent of this energy production is radiated and that 60 per cent is non-radiatively employed in the expulsion of gas. The mechanical power would therefore be larger than the present radiative power, $P = (0.6)/(0.4)L_* = 2.9 \times 10^{33}$ W. The mas flow driven by this power is $\dot{M} = PR_*/(GM_*) \simeq 7 \times 10^{21}$ kg/s. We take this as an indication of the real order of magnitude.

In an interesting paper, Hoyle, Solomon & Woolf (1973) have discussed the physics of such a continuous mass loss. The driving force would be the excess nuclear energy produced

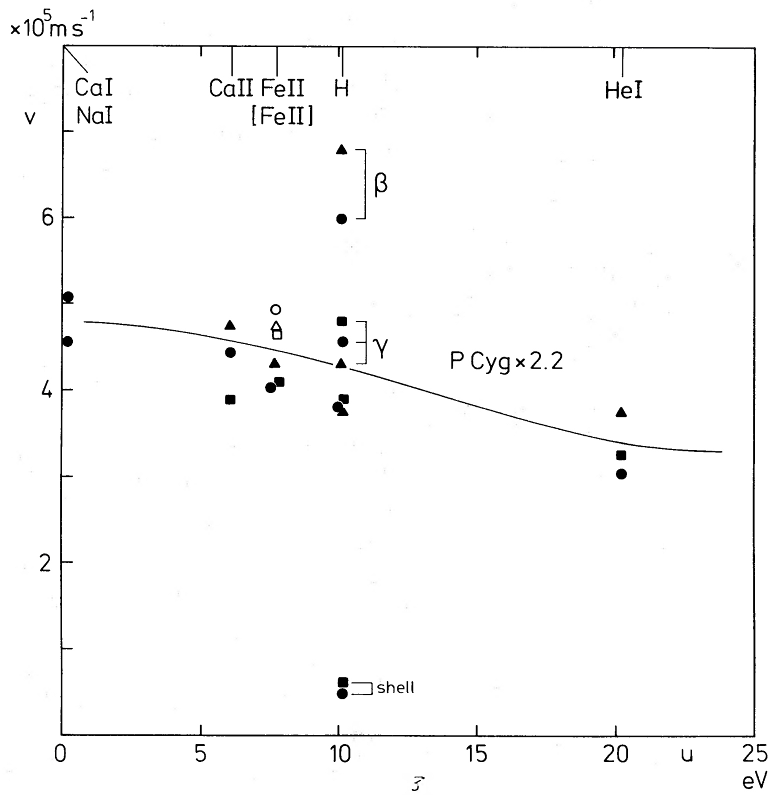

Figure 3. Radial velocity of absorption lines in η Car plotted versus the total excitation energy, in 194[illegible] (circles, from Viotti 1968), in 1953 (squares, from Thackeray 1953) and in 1961 (triangles, from Alle[illegible] & Dunham 1966). The open symbols refer to forbidden, the closed to permitted transitions. Radia[illegible] velocities of the central reversal (shell) of Hα and Hβ are also indicated. The curve gives the trend found i[illegible] P Cyg (de Groot 1969) at 2.2 times smaller velocities.

in the CNO-cycle as a result of temperature variations in stellar pulsations; the latter are onl[illegible] little damped in a weakly-bound massive star. This excess can be taken up 'by increasin[illegible] the dynamical energy of oscillation of the star, i.e. by increasing the amplitude; or, b[illegible] propagating a shock wave to the surface and lifting off the photospheric layers, thereb[illegible] producing mass loss'. The second mechanism takes over from the first when the amplitude become large. For the mechanism of the outflow of gas, which will not be discussed furthe[illegible] here, references can be made to P Cyg. In Fig. 3 we plot the velocities v of absorptio[illegible] lines of η Car as a function of the total excitation energy u of the pertinent atoms or ions The lines at the higher energies, formed close to the star, have lower velocities than the line at the lower energies, formed far from the star. This suggests an outward acceleratio[illegible] of the gas in the neighbourhood of the star, similar to the situation of P Cyg found by d[illegible] Groot (1969) and recently discussed by Barlow & Cohen (1977). According to Hoyl[illegible] *et al.*, the excess nuclear energy is capable of producing a continuous mass flow $\dot{M} \sim 6$ [illegible] $10^{-9} (L/L_{\odot}) M_{\odot}$/yr. Inserting the complete $L = 1/(0.4) L_*$ of η Car, we obtain $\dot{M} \sim 5$ [illegible] 10^{21} kg/s.

Judging the strength of the arguments in the above three independent approaches t[illegible] $\dot{M}$, we think that the best estimate for the mass flow is

$$\dot{M} \simeq 5 \times 10^{21} \text{ kg/s } (7.5 \times 10^{-2} M_{\odot}/\text{yr}). \quad (8$$

For our problem it is important to note that, thus, the efficiency of the condensatio[illegible] process is relatively *high*.

Let us now consider the speed of the outflow. The radial velocities shown in Fig. 3 refer to the inner 2 arcsec of the *homunculus.* Higher velocities are measured farther out, e.g. Thackeray (1961) describes a feature at 8.5×10^5 m/s. Lamers, van den Heuvel & Petterson (1976) suggested that there might be a constant ratio between the terminal velocity of gas from a star with mass loss and the escape velocity. If their ratio is applied to η Car, we get $2.8\,v^* \simeq 8.4 \times 10^5$ m/s. Let us therefore take 9×10^5 m/s as the final expression velocity of the gas.

If the outflow is constant, the mass density of the gas will obey

$$m_g(R) = m_g(R_0)(R_0/R)^2. \qquad (9)$$

The total mass M_g, lost until now, is given by equation (8) and the time elapsed since 1840; it is 2×10^{31} kg ($10\,M_\odot$). On the other hand M_g follows from the integral over $4\pi R^2 m_g(R)$ until the outer edge of the envelope $R_0 \simeq 3 \times 10^{15}$ m; this yields $M_g = 4\pi R_0^3\, m_g(R_0)$. The resulting mass density at R_0 is $m_g(R_0) \simeq 5.9 \times 10^{-17}$ kg m^{-3}, At $R_i \simeq 3.75 \times 10^{14}$ m, where the envelope of dust begins, we get 3.8×10^{-15} kg m^{-3}, which corresponds with an atomic density of about 2×10^{12} m^{-3}. These values refer to *average* densities at a certain distance.

For distances smaller than R_i the gas will be ionized by the ultraviolet flux from η Car. For distances larger than about $2R_i$ the gas is neutral, because the dust (with a total ultraviolet optical depth of the order of 4) will have effectively absorbed the ionizing photons at this distance. We thus expect that the gas densities, derived from lines of circumstellar ions, are of the order of, or higher than, the atomic density at $2R_i$ of 5×10^{11} m^{-3}. In fact, Rodgers & Searle (1967) have obtained a density of 3×10^{12} m^{-3} for what they call the cooler ionized region (the electron temperature is here 2×10^4 K); for the hotter (electron temperature 5×10^4 K) ionized region they require densities in excess of 10^{15} m^{-3}. In a study of scattering effects in the [N II] lines, Craine (1974) has arrived at a density of 3.8×10^{11} m^{-3} for what he calls the shell region (electron temperature 1×10^4 K). These data depend in a sensitive way on the value of E_{B-V}. Adopting $E_{B-V} = 1.2$ mag, Leibowitz (1977) recently showed that one requires the higher density of $\sim 10^{14}$ m^{-3} in combination with an electron temperature of $\sim 10^4$ K to fit both the forbidden and the permitted lines. Aitken *et al.* (1977) have concluded from the Brackett α line that 'η Car is surrounded by a small radiation bounded H II region of diameter less than 1.5 arcsec (5×10^{14} m) and electron density $\gtrsim 10^{12}$ m^{-3}'. Although partly in conflict with each other, these data are not in conflict with our calculation and thus indirectly confirm the hypothesis of a constant mass flow and the high value of the mass loss.

5 The condensation

The outflow of gas (equation 8) and the condensation of dust (equation 7) match each other in the sense that the condensible atoms soon appear in the condensate. Apparently the conditions for condensation are well fulfilled. This cannot be easily explained, however, in terms of the above picture of the gas envelope.

To condense solids, one roughly needs gas temperatures around 1700 K and gas pressures above 10^{-2} N m^{-2} (10^{-7} atm). We can substantiate this as follows. Equilibrium properties of a high-temperature cosmic condensate may be inferred from vaporization data of a sample of lunar rock (De Maria *et al.* 1971). They apply to various metals and oxides, of which the vapour pressure p, in N m^{-2}, is approximately described by

$$-\log p = 1 + T_c/T. \qquad (10)$$

The value of T_c depends on the species found in the vapour of the sample and is, in K, 1925 (Al), 1850 (Ca), 1800 (Cr), 1725 (Mn), 1675 (SiO), 1625 (Mg), 1525 (O_2), 1475

(Fe) and 1350 (Na). Further references to thermodynamic data of high-temperature condensates are given by Salpeter (1977). It follows from equation (10) that, in order to condense a species at half the temperature T_c, one can allow for an order of magnitude less pressure. Actually the vapour pressure is lower than the pressure needed to deposit atoms on a solid (*cf.* Yamamoto & Hasegawa 1977), but let us for the moment ignore deviations from the equilibrium. The pressures given by equation (10) are used merely as a convenient representation of the relative fluxes of particles on an atomic scale; particle densities $N = p/(kT)$, where k is Boltzmann's constant, could have been used as well. We thus require that, typically,

either

$$\left.\begin{array}{l} T = 1700\ \mathrm{K\ and}\ N = 5 \times 10^{17}\ \mathrm{m}^{-3}, \\ \mathrm{or} \\ T = 850\ \mathrm{K\ and}\ N = 5 \times 10^{16}\ \mathrm{m}^{-3}. \end{array}\right\} \tag{11}$$

The determination of gas temperatures, as a function of the distance to the star, is a problem in itself, as the above data by Rodgers & Searle (1967), Craine (1974) and Leibowitz (1977) are partly contradictory. In Fig. 4 we sketch a model, which is based on the following considerations.

Track A: from the photosphere to a distance of about 5 R_*, there is a slight increase due to dissipation in the supersonic turbulent flow. (The Reynolds number $v/\langle v^2\rangle^{1/2}$. L/l will exceed 10^8; the first ratio of the speed of outflow and the velocity dispersion will exceed unity, and the second ratio of the size of the envelope and the mean free path,

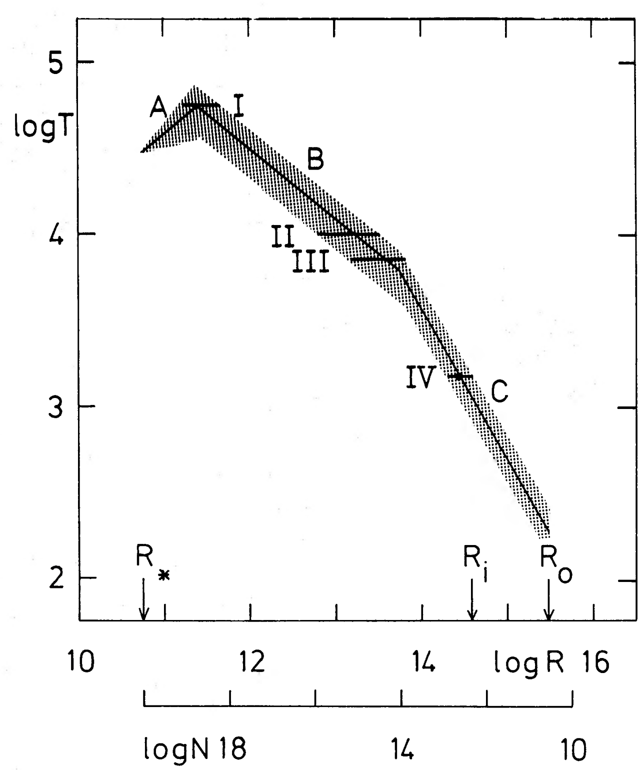

Figure 4. Model of the gas temperatures (see text). The units are K for the temperature T, m for the distance R and m^{-3} for the particle density N. N has been obtained from equation (11). The hatched region suggests uncertainties.

which around R_i is about 10^7 m, will be of the order of 10^8. Ter Haar (1972) argues that the critical Reynolds number for 3-dimensional flow is 10^7. As the above number is higher, we expect the flow of gas to be turbulent. One easily verifies that it is supersonic as well (*cf.* the temperatures in Fig. 4).) Such a temperature inversion seems to be common to OB stars and it is required to explain the temperature level I of the 'hotter ionized region' described by Rodgers & Searle. The density of this region around $5\,R_*$ would be of the order of $10^{18}\,m^{-3}$ (equation 9), which more than satisfies the condition that it should exceed $10^{15}\,m^{-3}$.

Track B: from $5\,R_*$ to the distance where the density becomes less than $10^{14}\,m^{-3}$, there is a temperature decrease according to $R^{-2/5}$. This dependence follows from the Whang–Chang theory of coronal expansion (*cf.* Hundhausen 1972), in which high densities constrain electrons and ions to have equal temperatures. The temperature level II refers to the 'shell region' of about 10^4 K, described by Craine to have a density of 3.8×10^{11} m^{-3}, but which, according to Leibowitz, has a density of $\sim 10^{14}\,m^{-3}$. Our model suggests a density of $10^{15}\,m^{-3}$, with an uncertainty of an order of magnitude. We have marked a level III of 7000 K, below which the opacity of the gas drops sharply from its high value $\mu = 4 \times 10^{-2}\,m^2\,kg^{-1}$ due to free-electron scattering. At the pertinent distance $R_{\rm III}$ of 3×10^{13} m, the optical depth in the gas

$$\tau_g = \int_{R_*}^{R_{\rm III}} dR\; \mu m_g(R) \tag{12}$$

becomes 0.02, as can be verified by inserting equation (9). (Note that, if Craine's 'shell region' of 10^4 K would extend to 10^{15} m (5 arcsec), as he holds, it has to be optically thick. This would have prevented the observation of the 'hotter ionized region' of Rodgers & Searle. It would also contradict the findings of Aitken *et al.* (1977) about the size of the H II region ($\lesssim 5 \times 10^{14}\,m \simeq 1.5\,R_i$). Therefore the 'shell region' cannot extend much farther than 10^{14} m.)

Track C: where the density comes below $10^{14}\,m^{-3}$, the electrons and ions are no longer constrained to have equal temperatures. These relatively low densities also apply to the solar corona and the plasma farther out from the Sun, where temperature differences have been established (Hundhausen 1972). The equipartition time, calculated from electrostatic forces alone (e.g. Spitzer 1968), which at the density $10^{14}\,m^{-3}$ is still much shorter than 1 s, may not apply to the case of turbulent flow with strong magnetic effects (see below). From this density on outwards we follow the temperature dependence $R^{-6/7}$, which is predicted by the two-fluid model of the solar wind (Sturrock & Hartle 1966) for the ions; for the electrons one has a slower decrease with $R^{-2/7}$.

In this model, the ion temperature reaches 1700 K (level IV) around 2.5×10^{14} m, slightly before R_i. The model further predicts a temperature of about 200 K towards the outer edge R_0, which allows for the condensation of volatile solids. Thus one can understand reasonably well why the gas temperatures allow for the condensation between R_i and R_0. Can the grains be substantially cooler than 1700 K around R_i? New observations by Koornneef (1976) give a brightness temperature of about 250 K at 10 μm in a circular diaphragm of 3.2 arcsec, which corresponds to a distance of $1.5\,R_i$; if the optical depth at 10 μm is 0.3 (Gehrz *et al.* 1973), the dust temperature would be about 1000 K near the inner edge. The radiative equilibrium temperature of about 2000 K at R_i, as calculated by Andriesse (1976), depends linearly on the energy of the average lattice resonance, which is uncertain by a factor of 2. The temperature model presented in Fig. 4 has a similar uncertainty and it actually yields about 1000 K rather than 1700 K at R_i. Thus there is some ground for cooler grains than 1700 K at the inner edge so that the second condition in equation (11) may apply. The real problem is now to understand how the density at R_i ca be as high as 5×10^{16} m^{-3}.

The average gas density N, fixed by equation (9), has been indicated in Fig. 4: at R_i it i 2×10^{12} m^{-3}. It should be kept in mind that it refers to a *total* particle density, wherea equation (11) refers to a *partial* particle density, namely that of the species involved i the condensation. The latter is at least 3 orders of magnitude smaller than the former. Thu there is a discrepancy of at least 7 orders of magnitude between the average gas densit and the density required for condensation.

To explain the factual condensation we infer that the gas cloud has a sponge structure with strong inhomogeneities of the density. Although speculative, we would like to discus this hypothesis in some detail.

Let us first refer to the observed structure of the *homunculus*. Gaviola's (1950) 1-s exposur of its core shows irregularities in the form of prominences, whereas permanent dense region ('condensations') have been traced for decades (Ringuelet 1958). It would be useful t study the spatial fine structure of the envelope with a large aperture and millisecond photog raphy. Such a study might reveal filaments and thus the presence of magnetic fields in th plasma. Turbulence may lead to dense regions, but only as transients. The stability of th 'condensations' cannot be understood, therefore, without a stabilizing magnetic field The origin of this field cannot be in the gas. If there is a seed field, turbulence will amplif it until equipartition is reached between kinetic and magnetic energy (J. Kuijpers, privat communication; Nagarajan 1971). But such a field, nourished by the turbulence, canno supersede it by freezing transient density inhomogeneities. The hypothetical field, fo which in order to withstand turbulent perturbations

$$B^2/(2\mu_0) > m_g \langle v^2 \rangle /2, \qquad (13$$

therefore has to be of *stellar* origin. It might pinch the circumstellar plasma around R_i an farther out, so that a density contrast is built up of 10^7 and more. This conjecture bring us close to the views by Alfvén & Arrhenius (1975) on a circumstellar plasma.

Condition (13) is satisfied if, at R_i, $B > 1.5 \times 10^{-5}$ Wb m^{-2} (0.15 gauss); let us adop $B(R_i) = 5 \times 10^{-5}$ WB m^{-2}. Mestel (1968) has discussed the magnetic field around a rotatin magnetic star with mass loss. He derives that, far from the star, $B \propto m_g R$ which, by virtu of equation (9), leads to $B \propto R^{-1}$. Note that for this dependence equation (13) is satisfie for any R. Using this as a scaling law to estimate the field at the stellar surface, we obtai $B(R_*) \simeq 0.3$ Wb m^{-2}. This may not be unreasonable, as it is common to Ap and Am star which, admittedly, are relatively small and do not involve as much magnetic energy a η Car. The energy of the hypothetical field is

$$U_m = \int_{R_*}^{R_0} dR \, 4\pi R^2 \frac{B^2(R)}{2\mu_0} = \frac{2\pi}{\mu_0} R_*^2 R_0 B^2(R_*), \qquad (14$$

and this yields 4×10^{42} J, being half the energy associated with the mass flow (Section 4 Thus, if a field exists that can stabilize the required density inhomogeneities, its strengt would imply strong interference with the flow dynamics, e.g. by transferring energy t the gas in accelerating it.

We conclude this speculative subject by a comment on the observability of the magneti field. In the case of η Car, the lines we see all originate above the photosphere. Let us co sider the lines closest to the star, where the Zeeman effect would be largest. Lines fro the 'hotter ionized zone' (level I in Fig. 4) might be formed at $B(R_*)/5 \simeq 6 \times 10^{-2}$ Wb m^{-2} for which the relative displacement would be of the order of 10^{-5}. But this effect is swampe in the kinetic line broadening of about 10^{-3}.

To sum up. We can understand without grave difficulties why the gas is cooled to condensation temperature at the onset of the dust envelope. We cannot understand the pertinent gas density, unless we invoke regions of a very much higher than average density. Such regions require a rather strong magnetic field as a stabilizer which, however, escapes direct observation and which is proffered here only as a possibility.

6 The condensate

So far the 8–13 μm spectrophotometric observations by Robinson *et al.* (1973) give the only hint to the nature of the condensed particles. We reproduce the observed hump in the spectrum in Fig. 5 and quote: 'The shape of this feature is similar to the excess emission found in M stars, which has been identified with the emissivity curves of certain silicates..... The exact nature of the emitting particles is however indeterminate, since the spectral energy distribution is influenced by, among other things, particle size effects, optical depth and temperature stratification as well as grain composition'. In this section we will try to disentangle these effects, so that a conclusion can be reached about the nature of the condensate. Some time ago, Donn & Sears (1963) predicted that the condensate should initially be a fluffy, disordered material. The general arguments for such a prediction can be found in the following comment on cosmic condensation processes.

As a first approximation, the nucleation in a gas can be based on the equilibrium theory of phase transitions (Hirth & Pound 1963). It is an irreversible process, however, and its kinetics should differ from transition rates of atoms between the gas and the already existing

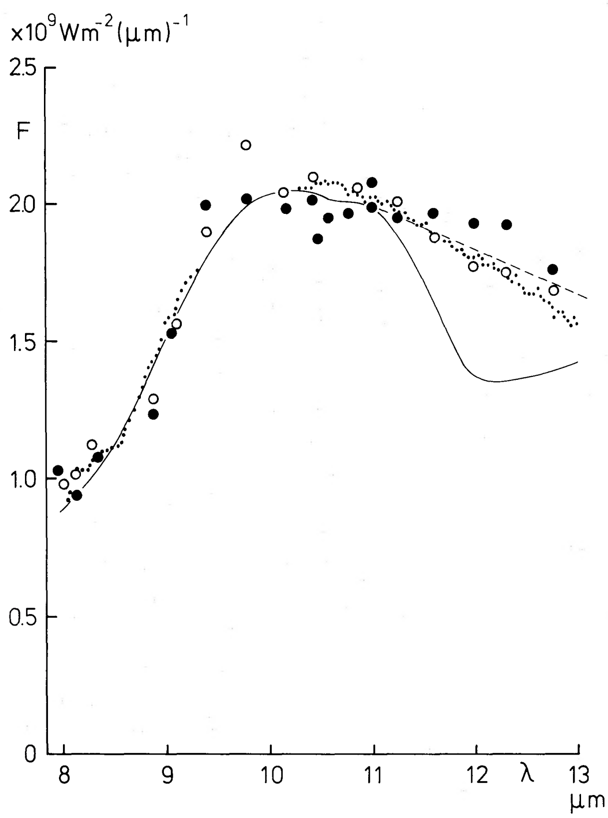

Figure 5. Infrared spectrum of η Car. The filled and open circles refer to two different series of observations by Robinson *et al.* (1973). The dots refer to new observations by Aitken & Jones (1975). The full and dashed curves are calculated from the full and dashed curves in Fig. 6 (see text at the end of Section 6).

condensate at equilibrium. Details of the physico-chemical mechanisms, by which the gas–solid transition is effectuated, will be important (Salpeter 1974; Donn 1976). One wil meet metastable states between the gas phase and the solid phase which, under circum stellar conditions, can have long relaxation times. To give an example, we examine the hot gas cloud of Al and O and its condensate.

Grossman & Larimer (1974) propose that this condensate consists of corundum, Al_2O_3 It is assumed that a gas phase reaction occurs

$$Al_2O_3 \rightleftarrows 2\,Al + 3O \qquad (15)$$

with a reaction constant K; the latter is taken from equilibrium theory and thus is typically temperature-dependent according to $\log K \propto -1/T$. If the partial pressures p_{Al} and p_O fulfil

$$p_{Al}^2\, p_O^3 > K, \qquad (16)$$

condensation of corundum is expected to follow. This would occur around $T = 1680$ K at a total gas pressure of $10\,N\,m^{-2}$. However, on vaporizing corundum, Drowart *et al.* (1960) had detected Al, O, AlO, Al_2O, O_2 and some Al_2O_2 in the gas phase, but not Al_2O_3. The polyatomic molecules found have a small heat of formation, in contrast to the very stable Al_2O_3.

In any binary gas mixture, like the one of Al and O, the polyatomic molecules that form easiest will be the most abundant. From Tsuji's (1973) extensive calculations of chemica equilibria at low pressures (for atmospheres of cool stars), we learn that both AlO and Al_2O form, the latter being the more abundant below 2000 K; the author did not include Al_2O_2. In passing, we note that at the low pressures considered by Tsuji, the vibrationa transition probabilities sometimes exceed the collisional excitation probabilities; thus also in this case the use of equilibrium theory is not always warranted. It should be con cluded that around 1700 K the relevant gas phase reaction is

$$Al_2O \rightleftarrows 2\,Al + O \qquad (17)$$

instead of (15). The condensate will therefore differ from corundum; the latter minera can form from the condensate only after compression and some heating.

In a gas cloud of cosmic composition, the nucleation comprises more than two atomic species. Donn (1976) has discussed the hierarchy of the many-atom processes, leading to a compound condensate in the form of a *cluster*. The cluster is understood to be a generally unstable group of various atoms. A representative cluster could be $Fe_7Mg_6(Ca, Al, Na)_1$ $(SiO)_5$, which will not be stable above about 1400 K. Between 1400 and 1700 K the cluste would be partially transformed (decayed) into the stable Al_2O_3 and $Ca_2Al_2SiO_7$ and into a number of metastable silicates. In such a cluster the free energy is not at a minimum and the atomic order, imposed by the chemical bonds, extends over a short range only Under circumstellar conditions the decay will be slow, because the required energy exchange with the gas is proportional to the density, so that the clusters will be different from any of the usual terrestrial minerals. Also, the temperatures at which the different atoms become enclosed in the clusters will differ from those at which the associated terrestrial mineral form.

Zaikowski, Knacke & Porco (1976) have investigated type I carbonaceous chondrites which might contain some initial high-temperature condensate. These meteorites consis primarily of a fine grained (~ 0.1 μm) *hydrous silicate* assemblage with the silicate occurring as serpentine, chlorites or montmorillonites. Determining the 8–13 μm transmission o these materials, they found in all the cases a striking resemblance with the broad absorption feature centered on 9.7 μm as observed towards the Becklin–Neugebauer object. Earlier

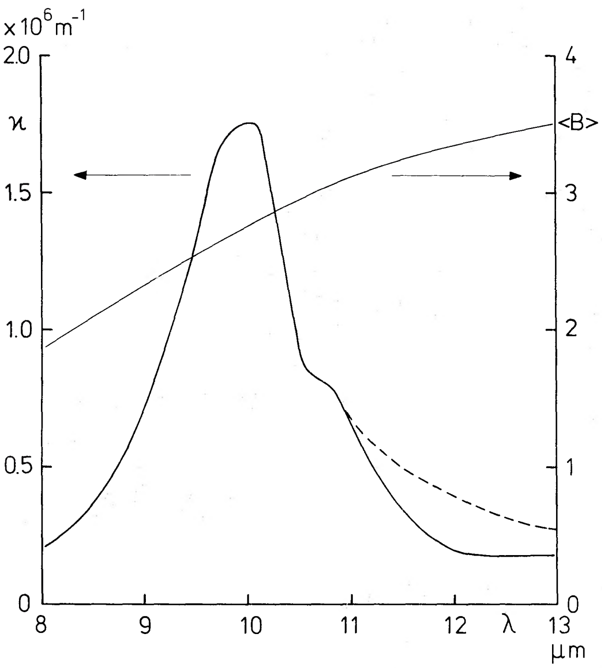

Figure 6. Absorption coefficient κ of a silicate gel around 10 μm, derived from data by Day (1974, 1975) and the geometrical average over the Planck function of the dust envelope $\langle B \rangle$ given by equation (19).

Day (1974) showed that a synthetic *silicate gel*, basically Mg_2SiO_4 without long-range order, gives a very similar absorption feature. Day & Donn (1978) recently reported about the formation of such solids in the laboratory. Physically, the broad absorption effect in both the hydrous silicate and the silicate gel should be due to lattice disorder, by which the vibrational Si–O resonances are spread. This is just what is expected to occur in thermodynamically unstable clusters.

Turning now to the condensate around η Car, we anticipate it to consist of disordered silicate clusters, at least partially. Fig. 6 shows 8–13 μm absorption coefficient $\kappa(\lambda)$ for these clusters, based on Day's (1974) transmission data of the gel. (We have assumed that the transmission is proportional to $\exp(-\kappa)$ and that the absolute value of κ is of the order of olivine (Day 1975); the estimated uncertainty in the absolute value is –50 and +100 per cent.) If a is the average radius, these clusters have an emission efficiency of

$$Q(\lambda) = 1 - \exp[-4/3\,\kappa(\lambda) a]. \tag{18}$$

According to Gehrz *et al.* (1973) the 10-μm optical depth in the envelope of η Car is 0.3, so that we can consider it to be virtually transparent. To calculate the spectral energy distribution, the only problem left is thus the temperature stratification in the envelope.

The dust temperature T around η Car has been calculated by Apruzese (1975) and Andriesse (1976) as a function of the distance R. We shall adopt the analytical approximation of the latter, as it comes close to the numerical results of the former for $R \gtrsim 2R_i$; furthermore, it should be more realistic near the inner edge R_i. (The constant $(R_0/R_*)^2\pi/18(\hbar\omega_r/kT_*)(\omega_r/\Omega)^4$, which appears in this approximation, is taken equal to 40.) The geometrical average over the Planck function of the dust envelope is

$$\langle B(\lambda) \rangle = \frac{1}{M_d}\int_{R_i}^{R_0} dR\, 4\pi R^2 m(R) B[\lambda, T(R)] \tag{19}$$

where $m(R)$ is given by equation (1). The result is shown in Fig. 6 also.

Using this geometrical average, we get the infrared spectrum $Q(\lambda)\langle B(\lambda)\rangle$. Keeping it in arbitrary units we have fitted this spectrum to the 8–11 μm data by Robinson *et al.* for a size of $a = 3\,\mu$m (see Fig. 5). This fit is good. However, the 11–13 μm data cannot be fitted with the κ of Fig. 6. One needs a slight enhancement of the absorption in this wavelength region (dashed curves in Figs 5 and 6), which is physically possible for slightly different disordered silicates. Uncertainties in the κ-scale of Fig. 6 suggest that the range of sizes, compatible with the curve in Fig. 5, is 1–5 μm. Having neglected optical depth effects in the dust envelope itself, we should interpret this as a border range. Clearly only particles with a size of the order of 1 μm can give sufficient internal extinction to explain the 'suppressed hump' in the spectrum.

We conclude that the condensate around η Car contains disordered silicate clusters with an average radius of the order of 1 μm.

7 Conclusions

In this paper we have brought together arguments from various disciplines, which may help to solve the general problem of condensation of solid matter around a star. We have taken η Car as an example, in view of the rather stringent constraints that can be derived for the circumstellar gas density and the rate of dust formation. Although in many respects this star is exceptional; the conclusions of Sections 5 and 6, that the condensation requires strong inhomogeneities in the gas envelope and that the initial condensate contains disordered silicates, can be of a more general significance. In particular, they stress the possible importance of magnetic compression for the condensation in the solar nebula.

Let us resume the arguments. We have first established the present amount of dust in the immediate neighbourhood of the star by discussing the colour excess E_{B-V}. This allowed us to make a distinction between the interstellar extinction showing the 2200 Å band and the circumstellar extinction without such a band. It was inferred that the circumstellar particles are much larger than 10^{-8} m. Then the light curve of the star has been interpreted in terms of a time-dependent optical depth of the dust envelope. Using a model of the geometrical evolution of this envelope, based on infrared photometry, we have derived that the dust has been condensing at the almost constant rate of 10^{19} kg/s.

As a second step we have discussed the physical conditions of the gas cloud from which the dust originates. This includes theoretical arguments about the outfloe of gas, estimated to be about 5×10^{21} kg/s since 1840, which implied that the condensation process is quite efficient. Furthermore, a model is discussed for the gas temperature. The conclusion was reached that those temperatures can explain the condensation, but that the densities are far too low.

The third, speculative, step has been to postulate dense regions, stabilized by a stellar magnetic field against turbulent perturbations. The required field strength would be that of 'normal' magnetic stars, but the total magnetic energy would be hardly less than the kinetic energy of the outflow. A possibly filamented structure of the gas cloud, which might indirectly confirm this hypothesis, would be difficult to observe, whereas the Zeeman effect is unobservable.

As a fourth and final step we have reviewed chemical arguments on the nature of the condensed particles. This has been included to show that it makes no sense to adopt optical characteristics of terrestrial or lunar minerals: the latter have been processed so that thermodynamically unstable configurations are removed. Adopting a plausible model of the temperature stratification of the dust envelope, we have shown that disordered silicate clusters can explain the infrared spectrum. The relatively large size of the clusters fits well in the high efficiency of the condensation.

This argument as a whole seems to us internally consistent and at the same time consistent with all observations available. It is hoped that it is helpful for the interpretation of dust formation processes around other stars.

Acknowledgments

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