ACCRETION BY THE SECONDARY IN η CARINAE DURING
THE SPECTROSCOPIC EVENT. I. FLOW PARAMETERS

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ABSTRACT

We examine the influence of the gravity of the companion (the secondary) to the massive primary star η Carinae on the winds blown by the primary and the secondary. The two winds collide with each other after passing through two respective shock waves, and escape the system while strongly emitting in the X-ray band. While during most of the 5.5 yr orbital period the companion’s gravity has a negligible effect on the winds, we find that near periastron the companion’s gravity may significantly influence the flow, and the companion might accrete from the primary’s wind under certain circumstances. Near periastron passage, the collision region of the two winds may collapse onto the secondary star, a process that could substantially reduce the X-ray luminosity. We suggest that such an accretion process produces the long, almost flat, X-ray minimum in η Car.

Subject headings: accretion, accretion disks — binaries: close — circumstellar matter — stars: individual (η Carinae) — stars: mass loss

1. INTRODUCTION

η Car is a very massive star that underwent a 20 yr long eruption, termed the Great Eruption, about 160 yr ago (Davidson & Humphreys 1997). A massive nebula, ~12 $M_\odot$ (Smith et al. 2003b), expelled during the Great Eruption, formed the Homunculus, an expanding bipolar nebula around η Car (Morse et al. 1998). A smaller bipolar nebula is believed to have been created during the Lesser Eruption of 1890 as well (Ishibashi et al. 2003). It is widely accepted now that η Car is a massive binary system with an orbital period of 5.54 yr (e.g., Damineli 1996; Damineli et al. 1997, 2000; Ishibashi et al. 1999; Corcoran et al. 2001a, 2001b, 2004a; Pittard & Corcoran 2002; Duncan & White 2003; Fernandez Lajus et al. 2003; Smith et al. 2004; Whitelock et al. 2004; Verner et al. 2005). The more massive companion of the η Car binary system will be referred to here as the primary, while the companion, probably an O-type star, will be referred to as the secondary.

The binary nature of η Car is inferred from the periodicity of the so-called spectroscopic event, the fading of high-excitation lines (e.g., Damineli et al. 2000). Since the spectroscopic event might in principle result from a mass-shell ejection by the primary star (Zanella et al. 1984; Davidson et al. 1999; Smith et al. 2003a; Martin & Koppelman 2004), it is also called a shell-ejection event. The main motivation to assume a shell ejection is in a single-star principle result from a mass-shell ejection by the primary star before the periodicity of the spectroscopic event was discovered (Damineli 1996). Based on theoretical calculations, Soker (2001b, 2003, 2004, 2005) argued that a single star cannot explain many of the properties of η Car; in particular, a single star cannot account for the bipolar shape of the Homunculus. Instead, these papers argued that the secondary accreted a large fraction of the mass that was expelled in the Great Eruption, forming an accretion disk and two jets, which shaped the wind of the primary into the bipolar Homunculus (Soker 2001b). Soker (2003, 2005) discusses the formation of an accretion disk in the present-day binary system, in agreement with some of the suggestions made by van Genderen et al. (1994, 1995, 1999).

Soker (2005) discussed three basic accretion phases, two of which occur in the present day. Both stars blow winds, with the primary’s wind having a much larger mass-loss rate and a lower velocity. The winds collide, and at one location the momentum fluxes of the two winds exactly balance each other, forming a stagnation point. The stagnation point is located close to, but not exactly on, the line joining the centers of the two stars, because of orbital motion. The shocked material of the primary star cools very fast (Pittard & Corcoran 2002; see eq. [1] in Soker 2003), namely, before the mass moves far from the stagnation point. The surrounding pressure then compresses the cooled post-shock gas to high densities. When the wind is blowing, radiation pressure on the escaping gas overcomes gravitational attraction. However, this might not be the case with the dense gas near the stagnation point. Gravitational force on the dense and slowly moving (much below escape velocity) gas there might become large enough to accrete part of the gas back onto one of the stars. Whether accretion occurs at all, and if it does, which of the two stars accretes most of the gas, depends on the accretion phase of η Car:

1. The Great Eruption.—The mass loss rate by the primary was very high, and the stagnation point was within the secondary’s Bondi-Hoyle accretion radius. Along the entire orbit the secondary steadily accreted mass with high specific angular momentum. An accretion disk was formed and two jets (or a collimated fast wind [CFW]) were launched (Soker 2001b). According to Soker (2001b), these jets shaped the two lobes that are now observed as the Homunculus.
2. Apastron passages.—Soker (2003) proposed that during present apastron passages the primary itself can accrete ∼5% of the mass lost over an entire orbit. This should not be confused with the accretion fraction of ∼50% by the secondary during the 20 yr span of the Great Eruption. The high specific angular momentum of the accreted gas implies the formation of an accretion disk around the primary. The primary star might blow a CFW.

3. Periastron passages.—Near periastron passages, which occur near the spectroscopic events, short accretion episodes might occur, possibly leading to pulsed ejection of two jets by the secondary (Soker 2005). This accretion process was proposed but not studied by Soker (2005). The secondary might also ionize a nonnegligible region in its surrounding neighborhood.

This paper presents further exploration of the nature of the wind interaction near periastron passage. In § 2, we compare the typical timescales and length scales of several processes during the periastron passage. The reader interested only in the main points of the proposed model can skip § 2 and go directly to § 3, where we discuss how these may allow an accretion event onto the secondary star for several weeks near periastron passage. Our main results are summarized and discussed in § 4.

2. RELEVANT LENGTH SCALES AND TIMESCALES AT PERIASTRON

2.1. Length Scales

We calculate the distance of the stagnation point of the secondary’s wind, taking into account the orbital motion; by orbital velocity, we mean the relative orbital velocity of the two stars. Following Usov (1992), we take the wind with the larger momentum flux, in our case the primary’s wind, to be plane parallel near the secondary star. Because of the orbital motion, the stagnation point of the secondary’s wind will be ahead of the line joining the centers of the two stars (Fig. 1). The radial (along the line joining the two stars) component of the relative velocity between the secondary star and the primary’s wind is $v_1 - v_r$, where $v_1$ is the radial component of the primary’s wind speed and $v_r$ is the radial component of the orbital velocity; $v_r$ is negative when the two stars approach each other. Because the stagnation point is very close to the secondary star, we assume that this is also the relative radial velocity of the stagnation point to the primary’s wind. The total relative speed between the stagnation point and the primary’s wind is

$$v_{\text{windl}} = \sqrt{v_r^2 + (v_1 - v_r)^2}^{1/2},$$  

where $v_r$ is the tangential component of the orbital velocity. We neglect any time delay between ejection of the wind by the primary and its collision with the secondary’s wind. This assumption is justified over most of the orbit, but not near periastron, where the orbital and the primary’s wind speed are almost equal. However, the uncertainties because of this assumption are less problematic than the uncertainties caused by our ignorance of the exact velocity profile of the primary’s wind, i.e., its acceleration over a distance of ∼3 AU, and a possible enhancement in the primary’s mass-loss rate near periastron. At the stagnation point the ram pressures of the two winds are equal,

$$\rho_1 v_1^2 = \rho_2 v_2^2,$$  

where $v_2$ is the secondary’s wind speed, assumed to be much larger than all other flow velocities in the problem.

The respective wind densities are

$$\rho_i = \frac{\dot{M}_i}{4\pi D_i^2 v_i}, \quad i = 1, 2,$$  

where $D_i$ is the distance from the respective star to the stagnation point (Fig. 1). Substituting the expressions for the winds’ densities in equation (2) gives

$$\frac{1}{D_1^2} \frac{v_1^2}{v_1^2} = \frac{\beta^2}{D_2^2},$$  

where

$$\beta \equiv \left(\frac{\dot{M}_2 v_2}{\dot{M}_1 v_1}\right)^{1/2}. $$  

As we see, $\beta$ depends on the orbital separation, because near periastron it is assumed that the primary’s wind does not yet reach its terminal speed.

Let $\phi$ be the angle measured from the secondary between the direction to the primary and that to the stagnation point (Fig. 1). The following trigonometric relation holds: $D_1^2 = D_2^2 + r^2 - 2rD_2\cos\phi$, where $r$ is the orbital separation, and $\cos\phi = (v_1 - v_2)/v_{\text{windl}}$. Substituting for $D_1$ in equation (4) gives an equation that can be solved for $D_2$:

$$D_2 = \beta r \left\{ \frac{v_1 - v_2}{v_{\text{windl}}} \right\} \frac{\beta^2}{v_2^2} + \frac{v_1^2 - \beta^2}{v_1^2} - \frac{v_1 - v_2}{v_{\text{windl}}} \beta \right\} \times \left(\frac{v_2^2}{v_1^2} - \beta^2\right)^{-1}. $$  

FIG. 1.—Top: Orbit and relevant velocities in rest frame of primary star of mass $M_1$. Drawn are the primary’s wind velocity $v_1$, the two stars’ relative orbital velocity $v_{\text{orb}}$, and $v_1$ and $v_2$, which are the radial and tangential components, respectively, of $v_{\text{orb}}$. Bottom: Geometrical definitions relevant to the flow near the stagnation point. The contact discontinuity is the surface where the two winds meet after they have passed the shock waves. The velocity directions in the bottom panel are as in the top panel (the secondary moves to the lower right).
Near periastron the gravitational influence of the secondary on the flow near the stagnation point should be considered. This effect on the undisturbed primary’s wind is characterized by the Bondi-Hoyle accretion radius,

\[ R_{\text{acc}} = \frac{2GM_2}{v_{\text{wind}}^2} = \frac{0.2}{30} \frac{M_2}{M_\odot} \left( \frac{v_{\text{wind}}}{500 \text{ km s}^{-1}} \right)^{-2} \text{AU}. \] (7)

The relative speed \( v_{\text{wind}} \) is scaled for periastron passage, where \( v_{\text{orb}} \approx 400 \text{ km s}^{-1} \) and \( v_1 \approx 300 \text{ km s}^{-1} \), since the primary’s wind has not yet reached its terminal speed of \( \sim 500 \text{ km s}^{-1} \). Martin et al. (2005), for example, take the primary’s wind speed to be \( v_1 \approx 500(r/3 \text{ AU}) \text{ km s}^{-1} \) for \( r < 3 \text{ AU} \), and \( v_1 \approx 500 \text{ km s}^{-1} \) for \( r \geq 3 \text{ AU} \). When approaching periastron, \( v_{\text{wind}} \) will be larger than its value when leaving periastron. This results in a much larger accretion radius when leaving periastron, as described in § 3.

As in the classical Bondi-Hoyle-Lyttleton accretion flow, the density and velocity of the inflowing gas increase as the gas approaches the gravitating point mass. This increases the ram pressure of the primary’s wind, and according to equations (2) and (3), the stagnation point distance from the secondary changes as \( D_2 \propto (\rho v^2)^{1/2} \). Using the density (from Danby & Camm 1957) and the velocity (from energy conservation) along the symmetry axis in the upflow direction (ahead of the stagnation point), one finds the stagnation distance when the secondary’s gravity is considered

\[ D_{\eta} = 2 \left( \frac{1 + R_{\text{acc}}}{D_2} \right)^{-1/4} \left[ 1 + \left( \frac{R_{\text{acc}}}{D_2} \right)^1 \right]^{-1/2} D_2. \] (8)

For the parameters used here, \( R_{\text{acc}}/D_2 \approx 0.5 \) near periastron, and hence \( D_{\eta} \approx 0.8D_2 \) near periastron. Some of the quantities derived here are plotted on Figure 2.

### 2.2. Timescales

As shown by several authors (e.g., Pittard & Corcoran 2002; Soker 2003), at all orbital phases the cooling time of the shocked primary’s wind is much shorter than the flow time of the gas out of the shocked region. The post-shocked primary’s wind in a large area near the stagnation point cools and is compressed by the ram pressure of the colliding winds. As shown by Soker (2005), cold and dense blobs of size \( r_b \approx 0.001D_2 \) will be accreted by the secondary. One of the uncertainties in this study involves the exact velocity profile and terminal velocities of the two winds. In particular, the radiation from one star can influence the velocity profile of the wind from the other star before the winds collide (Gayley et al. 1997). In any case, we estimate these uncertainties to be minor. The secondary star radius is about an order of magnitude smaller than the distance to the stagnation point. Hence, it is already at its terminal velocity, assuming that the radiation pressure of the primary does not slow down the secondary wind much. This is justified by (1) the calculations of Pittard & Corcoran (2002) of the X-ray emission, which show the secondary wind to be shocked at \( v_2 \approx 3000 \text{ km s}^{-1} \), with little variation in the shocked gas temperature over most of the orbit (Ishibashi et al. 1999), and (2) the location of the stagnation point far from the primary. For the primary wind we assume a simple velocity profile in § 2.3. Taking other reasonable wind properties for \( \eta \) Car from the literature will not much change the results of this paper.

In § 3 we propose that mass from the shocked primary wind is accreted for \( \sim 80 \) days near periastron passage. This phase begins with accretion of dense blobs. We estimate the blobs’ properties following Soker (2005). We consider a post-shock spherical blob of mass \( m_b \), density \( \rho_b \), and radius \( R_b \), located at a distance \( r_2 \) from the secondary of mass \( M_2 \). We make three assumptions:

1. Because of the short radiative cooling time mentioned above, the temperature of the blob is \( T_b \approx 10^4 \text{ K} \).
2. The blobs are in pressure equilibrium with the ram pressure of the secondary’s wind, so that at any distance \( r_2 \), \( \rho_b kT_b/\mu m_H = \rho_2 v_2^2 \), where \( \mu m_H \) is the mean mass per particle in the blob, and \( k \) is Boltzmann’s constant. Using the assumed values of \( M_2 = 10^{-3}M_\odot \text{ yr}^{-1} \) and \( v_2 = 3000 \text{ km s}^{-1} \) (see § 2.3) gives

\[ \rho_b = 5 \times 10^{-11} \left( \frac{r_2}{1 \text{ AU}} \right)^2 \text{ g cm}^{-3}. \] (9)

The mass in one blob is

\[ m_b = 7 \times 10^{-10} \frac{\rho_b}{10^{-10} \text{ g cm}^{-3}} \left( \frac{R_b}{0.01 \text{ AU}} \right)^3 M_\odot. \] (10)

3. The radiative pressure of the secondary plays a minor role. This is for two reasons. First, the ratio of the secondary radiation pressure to the ram pressure of the secondary wind is

\[ \xi = \frac{P_{\text{rad}}}{P_{\text{ram}}} = \frac{L_2/c}{M_2v_2} = 0.6 \left( \frac{L_2}{9 \times 10^5 L_\odot} \right), \] (11)

where the secondary luminosity is scaled according to Verner et al. (2005). In other words, the radiation pressure adds to the secondary ram pressure, but not more than 60%. The second reason is that near the stagnation region, where it is determined whether accretion takes place, the radiation pressure from the primary acts in the opposite direction to that of the secondary. As can be seen from Figure 2, the distance of the stagnation region from the primary (\( r - D_{\eta} \)) is \( \sim 3 \) times the distance to the secondary (\( D_{\eta} \)). However, the primary luminosity is estimated to be \( \sim 10 \) times that of the secondary. Since the radiation pressure drops as \( 1/(\text{distance})^2 \), we find that the radiation pressure of the primary, which pushes the gas toward the secondary, more than compensates for the secondary radiation pressure near the stagnation point. This holds as long as \( L_1 \geq 10L_2 \). Considering radiation pressure by both stars will actually increase the effects studied in the present paper.

The condition for accretion, therefore, is that the gravitational force

\[ f_g = \frac{G M_2 m_b}{r_2^2} \] (12)

on a blob must be larger than the force due to ram and radiation pressure of the secondary wind,

\[ f_{w2} = \rho_2 v_2^2 \pi R_b^2 (1 + \xi). \] (13)

Substituting the typical physical values used here and in equation (9) in the condition for accretion \( f_g > f_{w2} \) gives the constraint on the size of the accreted blob,

\[ R_b > 0.004(1 + \xi) \left( \frac{r_2}{1 \text{ AU}} \right)^2 \text{AU}. \] (14)
This shows that even very small blobs can be accreted. Substituting values in the last equation, we find $R_b = 0.007$ AU for $r_2 = D_{g2} = 1$ AU, when the orbital separation is $r \sim 4$ AU (accretion is not expected to start earlier than this time). According to equation (14), closer to the secondary even smaller blobs can be accreted. By equations (14), and (10) with (9), the minimum mass allowed for accreted blobs scales as $m_b \propto r_2^2$. In other words, as a blob moves toward the secondary, even if it breaks into smaller blobs, these may still be accreted if they are not too small. The blobs near the stagnation point can be somewhat smaller than the constraint given by equation (14) because the radiation pressure of the primary pushes toward the secondary.

The last point should be emphasized. If the flow structure were such that the primary wind streams undisturbed toward the secondary star, to the point where the secondary radiation pressure is larger than the primary radiation pressure $L_2/r_2^2 > L_1/(r - r_2)^2$,

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**Fig. 2.**—Several physical variables as a function of orbital phase (phase zero is at periastron). The left column covers the entire orbit, while the column on the right covers the time just prior to and after periastron. Top row: Orbital separation (in AU) and relative orbital speed of the two stars (in 10 km s$^{-1}$ on the left and 100 km s$^{-1}$ on the right). The angle $\theta$ is the relative direction of the two stars as measured from periastron (scale on the right in degrees). Second row: Cooling time of the shocked secondary’s wind near the stagnation point (top line; eq. [17]), the free-fall time from the stagnation point to the secondary (middle line; eq. [15]), and the flow time of the primary wind from the stagnation point (bottom line; eq. [16]). Third row: Distance of the stagnation point from the secondary (top line; eq. [8]), and the Bondi-Hoyle accretion radius of the secondary star (bottom line; eq. [7]). Bottom row: The ratio $R_{acc2}/D_{g2}$ (thin line) and $\tau_{f1}/\tau_{f2}$ (thick line).
then the radiation pressure might slow down or even expel the incoming primary wind, a process termed radiative braking (e.g., Gayley et al. 1997). However, the primary wind encounters a shock wave before radiative braking starts. Dense blobs are likely to form in this unstable interaction region. We conclude that the radiation pressure and ram pressure cannot prevent the accretion by the secondary of even small blobs.

As discussed in §3, if the stagnation region collapses, the maximum steady state accretion rate is \(M_{\text{acc}} \approx 10^{-6} \ M_\odot \ \text{yr}^{-1}\). When the stagnation region collapses before or near periastron, the accretion radius is smaller and the Bondi-Hoyle formula will give a lower accretion rate. However, if (as proposed in §3) the entire stagnation region collapses, we can still crudely take an accretion rate of \(\sim 10^{-6} \ M_\odot \ \text{yr}^{-1}\). As shown below, the accretion dynamical timescale is \(\sim 0.01 \ \text{yr}\). Using the typical blob mass from equation (10), we find that \(\sim 10\) blobs exist during the collapse of the stagnation point region. These blobs when accreted onto the secondary are assumed to vigorously disturb the acceleration region of the secondary wind, such that the secondary wind for all practical purposes ceases to exist. For the rest of the flat minimum phase, we assume that there is no wind collision and the accretion flow is of the Bondi-Hoyle-Lyttleton type. That the accreted primary wind almost completely shuts down the secondary wind is a strong assumption of the present paper, and must be checked in future calculations of the acceleration zones of accreting O stars, and observations of \(\eta\) Car during minimum.

The free-fall time from the stagnation point to the secondary is

\[
\tau_{\text{ff2}} = 1.05 \left( \frac{M_2}{30 \ M_\odot} \right)^{-1/2} \left( \frac{D_{\text{g2}}}{0.2 \ \text{AU}} \right)^{3/2} \text{day}. \tag{15}
\]

The outflow time from the stagnation point is somewhat longer than the flow of the undisturbed primary wind

\[
\tau_{\text{flow1}} \gtrsim \tau_{\text{ff1}} \equiv \frac{D_{\text{g2}}}{v_{\text{wind1}}} = 0.69 \frac{D_{\text{g2}}}{0.2 \ \text{AU}} \left( \frac{v_{\text{wind1}}}{500 \ \text{km s}^{-1}} \right)^{-1} \text{day}. \tag{16}
\]

The cooling time of the secondary’s wind is longer than the flow time \(\tau_{\text{flow2}}\); for \(v_2 = 3000 \ \text{km s}^{-1}\), \(\tau_{\text{flow2}} > \tau_{\text{ff2}} \equiv 0.1(D_{\text{g2}}/0.2 \ \text{AU})\). However, near periastron when the stagnation distance \(D_{\text{g2}}\) becomes small, the cooling of gas very close to the stagnation point becomes nonnegligible because very close to the stagnation point \(\tau_{\text{flow2}} \gg \tau_{\text{ff2}}\) (Usov 1991). The secondary’s wind is shocked to a temperature of \(\sim 10^8 \ \text{K}\), where the cooling function dependence on temperature is \(\Lambda = \Lambda_0 T^{0.4}\). The cooling time at constant pressure is \(\tau_{\text{cool2}} = (5/2)nkT(n_e n_p \Lambda)^{-1}\), where \(n, n_e, n_p\) are the total, electron, and proton densities, respectively. Substituting the numerical values gives for the cooling time of the shocked secondary's wind

\[
\tau_{\text{cool2}} = 1.1\left( \frac{M_2}{10^{-5} \ M_\odot \ \text{yr}^{-1}} \right)^{-1} \left( \frac{D_{\text{g2}}}{0.2 \ \text{AU}} \right)^2 \left( \frac{v_2}{3000 \ \text{km s}^{-1}} \right)^{-2.2} \text{day}. \tag{17}
\]

(Near periastron and for these parameters the post shock electron density is \(\sim 10^{10} \ \text{cm}^{-3}\) and the temperature is \(\sim 10^8 \ \text{K}\); the equalization time of ion and electron temperature for these values is \(\sim 0.1 \ \text{day}\) [Usov 1992], shorter than the cooling time.) The fraction of the shocked wind that cools to very low temperatures and is compressed is (eq. [7] of Usov 1991) \(\alpha \sim (\tau_{\text{ff2}}/\tau_{\text{cool2}})^2\), which is \(\sim 0.01\) near periastron passage. This implies that at periastron passage the shocked secondary’s wind within a distance of \(\sim 0.1D_{\text{g2}}\) from the stagnation point cools and is compressed. This will further reduce the support to the shocked primary’s wind against being accreted by the secondary.

### 2.3. Numerical Values

The parameters used in the present calculations are based on papers cited above (e.g., Pittard & Corcoran 2002; Martin et al. 2005). We take the primary’s wind to have a profile of \(v_1 = 500[1 - (0.4 \ \text{AU}/r_1)] \ \text{km s}^{-1}\), where \(r_1\) is the distance from the center of the primary; as this expression is a crude estimate of the acceleration zone of the primary’s wind, we can take \(r_1 = r\) at the stagnation point. At periastron \(r = 1.66 \ \text{AU}\), and \(v_1 = 380 \ \text{km s}^{-1}\), which is larger than the speed assumed by Martin et al. (2005). Based on the results of Pittard & Corcoran (2002) the secondary’s wind speed is taken to be \(v_2 = 3000 \ \text{km s}^{-1}\), and the mass-loss rates are assumed to be \(M_1 = 3 \times 10^{-4} \ M_\odot \ \text{yr}^{-1}\) and \(M_2 = 5 \times 10^{-5} \ M_\odot \ \text{yr}^{-1}\). The masses are \(M_1 = 120 \ M_\odot\) and \(M_2 = 30 \ M_\odot\) (Hillier et al. 2001), the eccentricity is \(e = 0.9\) (Smith et al. 2004), and the orbital period is 2024 days; hence, the semimajor axis is \(a = 16.64 \ \text{AU}\).

### 3. ACCRETION NEAR PERIASTRON PASSAGE

The relevant length scales and timescales derived in §2 are plotted as a function of the orbital phase, where phase zero is taken at periastron, in Figure 2. Note that phase zero corresponds to periastron passage. This is not to be confused with phase zero defined from observations of the intensities of different lines; in the latter cases, phase zero is assumed to be near periastron, but may not correspond precisely to periastron passage (see footnote 4 in Martin et al. 2005).

Two ratios determine the importance of the secondary’s gravity: the accretion radius to stagnation distance \(R_{\text{acc2}}/D_{\text{g2}}\), and the flow time of the shocked primary’s wind to the free-fall time from the stagnation point \(\tau_{\text{ff1}}/\tau_{\text{ff2}}\). Over most of the orbital motion these ratios are very small (Fig. 2, lower row), and gravity is negligible. However, very close to periastron these ratios become \(\sim 0.5\). For example, both these ratios are larger than \(\sim 0.25\) from 10 days before to 40 days after periastron passage. These large ratios suggest that accretion might take place near periastron passage. The asymmetry of these ratios around periastron fits well with the asymmetrical behavior of the event around periastron. For example, the X-ray emission after the flat minimum period does not return to its luminosity prior to the flat minimum period. This asymmetry results from the larger relative wind velocity \(v_{\text{wind1}}\) as the system approaches periastron and \(v_2 < 0\), than when the system leaves periastron and \(v_2 > 0\) (eq. [1]). This then influences the values of the stagnation distance (eq. [8]) and the accretion radius (eq. [7]), which determines the other properties, such as cooling time.

The exact flow structure requires 3D gas-dynamical simulations. However, we can suggest the following scenario already from the present results. Due to thermal instabilities (e.g., Stevens et al. 1992), dense large blobs are formed in the post-shock primary’s wind region near the stagnation point. These blobs are pulled to the secondary as periastron is approached. Very close, possibly \(\sim 10\) days prior, to periastron passage, the mass of the primary’s wind that is accreted is assumed to be large enough to shut down the secondary wind. The assumed shutdown must be nonlinear, because the mass accretion rate is smaller than the mass-loss rate of the secondary. As the secondary’s wind no longer reaches the previous stagnation region, the entire primary’s wind entering the Bondi-Hoyle accretion cylinder, i.e., having
an impact parameter smaller than the accretion radius, will be accreted by the secondary. In other words, the previously colliding winds region collapses onto the secondary. We emphasize that a key ingredient in the model is that blobs accreted near periastron passage shut down, or substantially weaken, the secondary wind, such that the accretion radius becomes larger than the stagnation distance. This allows more accretion that is assumed to shut down the secondary wind in several days. This process might occur even if during the time the first blobs are accreted the accretion radius is smaller than the stagnation distance.

In a steady state situation the mass accretion rate is \( M_{\text{acc}} = \pi R_{\text{acc}}^2 \rho \Omega v_{\text{wind}} \), which for the parameters used here reaches a maximum value of \( M_{\text{acc}} \approx 0.006 M_\odot \approx 10^{-6} M_\odot \) yr\(^{-1}\). This is not a high accretion rate, compared with the mass blown by the secondary. However, a short time after the region near the stagnation point collapses, this mass falls onto the secondary from one direction. As shown here and by Soker (2005), dense blobs can be accreted by the secondary. After the collapse of the stagnation point, the primary’s wind accelerates and its density increases as it approaches the secondary. It is assumed that the accreted cold gas prevents the secondary’s normal wind acceleration. As the stars separate, the mass accretion rate declines and the secondary wind reappears again, building the wind collision region, only when orbital separation increases again to several AU. The orbital separation is twice its periastron values (1.66 AU) after 20 days, and five times its periastron distance after 75 days. When the colliding winds region is built again, the X-rays reappear.

In a steady state accretion from a wind, the accreted mass possesses angular momentum. If the specific angular momentum is larger than that of a test particle performing Keplerian motion on the equator of the accreting star, an accretion disk might be formed. The condition for that is given by, e.g., Soker (2001a, their eq. [1]). Substituting the physical parameters used here, we find that in a steady-state accretion the accreted primary’s wind possesses specific angular momentum, which is too low by a factor of \( \approx 10 \) to form an accretion disk around the secondary.

One reasons for this is that in the Bondi-Hoyle type accretion flow, the accretion flow rearranged itself to accrete a small fraction (\( \approx 20\% \)) of the angular momentum entering the Bondi-Hoyle accretion cylinder. However, during the brief collapse of the material near the stagnation point this reduction does not happen. Furthermore, the collision region near the stagnation point is larger than the accretion radius, and the specific angular momentum is larger. Overall, we suggest that an accretion disk does form for a brief period, a few days, during the collapse stage of the colliding wind region (the collapse of the stagnation point). Such an accretion disk could blow a collimated fast wind.

4. DISCUSSION AND SUMMARY

In § 3, we have shown that near periastron the secondary’s gravity becomes a significant factor in determining the flow near the stagnation point. Along most of the orbit, the accretion radius \( R_{\text{acc}} \) (eq. [7]), which characterizes the influence of the secondary on the undisturbed primary wind, is much smaller than the distance of the stagnation point from the secondary \( D_{\varnothing} \) (eq. [8]). Along most of the orbit, the outflow time of the post-shock primary’s wind from the stagnation point vicinity \( (\tau_{\varnothing}, \text{eq. [16]}) \) is much shorter than the free-fall time of this gas to the secondary \( (\tau_{\varnothing}, \text{eq. [15]}) \). However, very close to periastron passage these two ratios, \( R_{\text{acc}}/D_{\varnothing} \) and \( \tau_{\varnothing}/\tau_{\varnothing} \) (Fig. 2, bottom row), increase to \( \approx 0.5 \). This shows that the secondary significantly alters the flow, such that it might accrete from the cool post-shock primary’s wind. We speculate that this could lead to the collapse of the wind collision region near the stagnation point. We further speculate that for a short time an accretion disk and a collimated fast wind might be formed. We note that the accretion phase is relatively short, and the steady state accretion rate relatively low \( (M_{\text{acc}} \sim 10^{-6} M_\odot \) yr\(^{-1}\); § 3). Therefore, the total accreted mass is negligible compared with the mass lost in the secondary wind.

In addition to the assumptions made in developing the model, there are some not well-determined binary parameters which introduce further uncertainties. We consider the results with regard to uncertainties in wind parameters quite robust. This is because the wind parameters are constrained by the X-ray properties (Pittard & Corcoran 2002). More than that, it is possible that the equatorial mass loss by the primary will increase near periastron passages (Pittard & Corcoran 2002; but see Soker 2005), thus enhancing accretion. The model is only slightly more sensitive to the eccentricity as long as \( e \approx 0.8 \). If \( e = 0.8 \) instead of \( e = 0.9 \), for example, we find that the maximum ratio of \( R_{\text{acc}}/D_{\varnothing} \) is 0.23 instead of 0.47 for \( e = 0.9 \). However, in the case \( e = 0.9 \), which is assumed in this paper, this ratio stays at a value of \( R_{\text{acc}}/D_{\varnothing} \approx 0.2 \) for \( \approx 65 \) days, about the length of the flat X-ray minimum. It is assumed here that for such a value, the mass accretion rate is high enough to shut down the secondary wind. In the case of \( e = 0.8 \), the inequality \( R_{\text{acc}}/D_{\varnothing} \approx 0.2 \) holds for \( \approx 40 \) days, but the inequality \( R_{\text{acc}}/D_{\varnothing} \approx 0.17 \) holds for \( \approx 65 \) days. We see that the conditions for \( \approx 65 \) days accretion are not much different in the two cases. This suggests that even for \( e = 0.8 \) accretion might occur, and our model can hold even for eccentricity as low as \( e = 0.8 \). Of course, this depends on the validity of our assumptions that (1) accretion starts for these values of \( R_{\text{acc}}/D_{\varnothing} \), and (2) the accretion can shut down the secondary wind.

The suggestion of the collapse of the interacting winds region is disrupted. Other researchers argue that the winds collision region (the stagnation point region) continues to exist during the X-ray minimum, e.g., Abraham et al. (2005), who based their claim on their suggestion that the 7 mm emission comes from this region.

In our proposed scenario, the collapse of the wind collision region to the secondary is behind the long and almost flat minimum in the X-ray emission, lasting \( \approx 60 \) days (0.03 of the cycle) (Ishibashi et al. 1999; Corcoran et al. 2001a, 2004a; Corcoran 2005). The X-ray emission results from the collision of the two winds from the two stars (Corcoran et al. 2001a; Pittard & Corcoran 2002), as in similar massive binary systems (Usos 1992), most prominently, the WR-O binary system WR 140 (Williams et al. 1990). In the interacting massive binary system WR 140, the X-ray minimum is not flat, and it can easily be explained by absorption of the X-ray emission by the dense wind of the WR star (Corcoran et al. 2004a; Pollock et al. 2005). In \( \eta \) Car, the flat minimum in X-ray emission seems to be intrinsic (Corcoran et al. 2000; Hamaguchi et al. 2005), that is, a reduction in emission measure of the X-ray emitting gas, and the minimum is not easy to explain only by absorbing material. The hard (\( \gtrsim 5 \) keV) X-ray emission drops by a factor of up to \( \approx 100 \) during the minimum, but there is still hard X-ray emission. In Akashi et al. (2006) this residual X-ray emission during minimum is explained in the framework of the accretion model.

After the collapse of the stagnation region, the primary’s wind collides with the secondary’s wind very close to the secondary, most likely within the acceleration zone of the secondary’s wind. In this type of wind interaction, the secondary’s luminosity slows down the primary’s wind before it encounters the shock wave (Gayley et al. 1997), substantially reducing the X-ray intensity and making it softer (Usos 1992). The slowing down
process by the secondary radiation pressure affects the location of the new stagnation point, but does not push it back to its original place. Hence, it cannot prevent the proposed collapse of the colliding wind region. In any case, this process is not relevant to \( \eta \) Car because the primary wind is too slow to contribute to the observed X-ray emission above 1 keV. This is true even if the gravitational acceleration by the secondary is taken into account (Akashi et al. 2006). The effects of the secondary’s gravity and radiation on this asymmetrical flow are very complicated, and are postponed for discussion to a future paper. In addition to the reduction in the mass of the colliding primary’s wind and its deceleration, a substantial fraction of the X-rays emitted by the shocked winds is absorbed by the dense accreted mass (Akashi et al. 2006).

We also speculate that the accretion process is connected to the peak luminosity of the He \( \text{ii} \lambda 4687 \) line. This peak occurs very shortly after the X-ray flux starts to decline (Steiner & Damineli 2004; Martin et al. 2005). In a future paper we will examine the possibilities that the He line is emitted by the cooling secondary’s wind as its cooling time \( \tau_{\text{cool}} \) (eq. [17]) becomes short (Fig. 2, row 2), or that the He line is emitted by a collimated outflow blown by the accreting secondary.

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