ACCRETION ONTO THE COMPANION OF η CARINAE DURING THE SPECTROSCOPIC EVENT. II.
X-RAY EMISSION CYCLE

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ABSTRACT

We calculate the X-ray luminosity and light curve for the stellar binary system η Car for the entire orbital period of 5.54 yr. By using a new approach we find, as suggested in previous works, that the collision of the winds blown by the two stars can explain the X-ray emission and temporal behavior. Most X-ray emission in the 2–10 keV band results from the shocked secondary stellar wind. The observed rise in X-ray luminosity just before minimum is due to the increase in density and subsequent decrease in radiative cooling time of the shocked fast secondary wind. Absorption, particularly of the soft X-rays from the primary wind, increases as the system approaches periastron and the shocks are produced deep inside the primary wind. However, absorption cannot account for the drastic X-ray minimum. The 70 day minimum is assumed to result from the collapse of the collision region of the two winds onto the secondary star. This process is assumed to shut down the secondary wind, and hence the main X-ray source. We show that this assumption provides a phenomenological description of the X-ray behavior around the minimum.

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

1. INTRODUCTION

The X-ray light curve of η Carinae has a period of 5.54 yr with a deep, nonzero minimum of ~70 days (Corcoran 2005) every period. Some variability is observed during the minimum (Hamaguchi et al. 2005). The minimum is more or less coincidental with the fading of many visible emission lines and a decline in the IR. It is widely accepted that η Car is a binary system (e.g., Damineli 1996; Damineli et al. 1997, 2000; Ishibashi et al. 1999; Corcoran et al. 2001a, 2001b, 2005; Pittard & Corcoran 2002; Duncan & White 2003; Fernandez Lajus et al. 2003; Smith et al. 2004; Whitelock et al. 2004; Steiner & Damineli 2004; Verner et al. 2005; Iping et al. 2005), and we refer to it as such. The more massive companion of the η Car binary system (the one that caused the Great Eruption of ~1840; see the description of the eruption in Davidson & Humphreys 1997) is referred to here as the primary, while the companion, probably an O-type star (Verner et al. 2005), is referred to as the secondary. The aforementioned periodic brightness minima are believed to occur near periastron passage and are generally termed the spectroscopic event (e.g., Damineli et al. 2000).

Corcoran et al. (2001a) and Pittard & Corcoran (2002) demonstrated that the collision of the winds blown by the two stars can account for the X-ray light curve. However, they had problems in accounting for the X-ray minimum. The X-ray properties of η Car are reviewed in § 2.1. In examining the properties of the colliding winds of the two stars, Soker (2005a, 2005b, hereafter Paper I) has suggested that for several weeks near periastron passages the secondary accretes mass from the primary. This suggestion of accretion is motivated by observations, which show accretion from winds in many other types of binary systems, e.g., symbiotic systems (e.g., Skopal 2005), and from numerical simulations showing the likelihood of accretion from a wind (e.g., Jahanara et al. 2005), in which in some cases an accretion disk is formed (Mitsumoto et al. 2005). This accretion, Soker suggested, shuts down the secondary’s wind, and hence the X-ray emission, leading to the X-ray minimum. This scenario is described in § 2.2. However, Soker did not explore the X-ray light curve around the minimum. This is the subject of this paper. In § 3 the X-ray emission of the colliding winds is examined. The X-ray emission during the accretion phase is examined in § 4. Our discussion, in which the model for the X-ray minimum is summarized according to the results of §§ 3 and 4, is presented in § 5.

2. X-RAY MINIMUM

2.1. Observed X-Ray Minimum

The main relevant properties of the X-ray emission near the minimum are as follows:

1. Minimum.—The X-ray emission of η Car between 2 and 10 keV is normally at a level of $4 \times 10^{34}$ ergs s$^{-1}$. This emission is believed to result from the collision of the two winds from the two stars (Corcoran et al. 2001a; Pittard & Corcoran 2002), as in similar massive binary systems (Usov 1992), most notably the WR-O binary system WR 140 (Williams et al. 1990). There is an almost flat minimum in the X-ray emission of η Car, lasting ~70 days (0.035 of the cycle; Ishibashi et al. 1999; Corcoran et al. 2001a, 2005; Corcoran 2005). The interacting massive binary system WR 140 shows a similar X-ray light curve to that of η Car, but the minimum is very short and not as flat. Hence, it can be explained by absorption of the X-rays by the dense wind of the WR star (Corcoran et al. 2005; Pollock et al. 2005). Conversely, in η Car the X-ray minimum is almost flat and appears to be a reduction in observed X-ray emission measure, which is difficult to explain by absorption alone (Ishibashi et al. 2003; Hamaguchi et al. 2005).

2. Behavior prior to minimum.—The X-ray intensity increases prior to the start of the X-ray minimum and then drops sharply to minimum (Corcoran 2005). In this paper phases are measured from periastron passage, which may or may not be the same as

1 Digital data were taken from http://lheawww.gsfc.nasa.gov/users/corcoran/eta_car/etacar_rxte_lightcurve.
the start of the X-ray minimum (Martin et al. 2006). The decline starts before the maximum of the He ii $\lambda$4686 line (Martin et al. 2006) and the maximum of the IR (Whitelock et al. 2004). The X-ray decline starts just after phase 0.98 (or $-0.02$); in the 2003 minimum there was an X-ray flare at phase 0.99 (Corcoran 2005). The minimum starts just before phase zero and ends at phase $\sim0.03$ (or $\sim1.03$). The sharp decline in the IR and He ii $\lambda$4686 line starts at phase $\sim0$.

3. No hardening during the minimum.—For $\sim40$ days into the minimum, the RXTE-observed $2-10$ keV X-ray luminosity is $\sim3 \times 10^{33}$ ergs s$^{-1}$ (for a distance of 2.3 kpc to $\eta$ Car), with large relative variations, e.g., spikes (Corcoran 2005). However, the RXTE field of view includes both the central source and the surrounding nebula. The X-ray emission by the extended nebula is much softer than emission from the stellar source (Hamaguchi et al. 2005), and hence RXTE finds a softer spectrum during the minimum (Corcoran 2005), which is not intrinsic to the central source. Hamaguchi et al. (2005) observed $\eta$ Car 24 days into the minimum (on 2003 July 22) with XMM-Newton.

Despite the huge decline, the X-ray spectrum does not become harder. This shows that absorption cannot be the main reason for the X-ray decline. After another several days the X-ray luminosity starts to increase for a while (Corcoran 2005) and becomes harder (Corcoran 2005; Hamaguchi et al. 2005). This again contradicts an absorption effect in which a rise in flux (less absorption) would be accompanied by softening. We note that the softening in the RXTE plot of Cororan (2005) starts a little after the decline starts. If the hardness variation was only due to the central source diminishing, while the surroundings remain constant, then the total radiation observed by RXTE should have become softer at the same time the decline started. This shows that although absorption is not the main factor during the minimum, it may still play a role during the decline to and the rise out of the minimum.

About 30 days into the X-ray minimum, RXTE detected an increase in X-ray flux to $\sim6 \times 10^{33}$ ergs s$^{-1}$. The spectrum becomes even a little harder (Corcoran 2005). This is followed by a slow decline back to $\sim4 \times 10^{33}$ ergs s$^{-1}$ and by softening of the emission before the steep rise out of the minimum.

4. Flares.—Flares occur during most of the orbit. Their effect in the RXTE 2–10 keV band is a $\sim10\%$ increase in luminosity. This may indicate stochastic variations in one or two of the wind parameters (velocity, density, inhomogeneity), and/or instabilities in the wind interaction region. The flare amplitudes increase somewhat with the X-ray average luminosity after phase $\sim0.6$ (Corcoran 2005). However, during the maximum period before decline, there are very strong variations, with “flares” that increase the intensity by up to $\sim50\%$ (Corcoran 2005). Such strong flares most likely indicate much stronger instabilities in the winds interaction process. Such instabilities may lead to the formation of dense blobs in the postshock primary wind region, which eventually are favorable for accretion by the secondary (Paper I).

5. Cycle-to-cycle variation.—Two minima were observed in X-rays by RXTE. The X-ray luminosity at the maximum emission just prior to the X-ray minimum in 2003.5 is larger than that in 1997.9 (Corcoran 2005). The qualitative behavior is similar in the two minima, however, with similar timescales.

2.2. Proposed Model

In order for absorption to account for the deep X-ray minimum, Corcoran et al. (2001a) were required to assume that the mass-loss rate from the primary in $\eta$ Car increases by a factor of $\sim20$ for 80 days following periastron. This model also requires the secondary to be behind the primary during periastron (Corcoran et al. 2005). Ishibashi (2001) first showed that a semimajor axis orientation perpendicular to our line of sight provided the best fit to the asymmetry in the X-ray light curve before and after the event, if the secondary is behind the primary after the event. We note that in a recent paper Smith et al. (2004) argue that the semimajor axis of the $\eta$ Car binary system is indeed more or less perpendicular to the line of sight, but the companion is not behind the primary near or after periastron; rather the secondary is behind the primary before periastron passage. Falceta-Goncalves et al. (2005) and Abraham et al. (2005) build a different model, in which the secondary is in front of the primary during periastron; after the event, if the secondary is behind the primary after the event, the model requires that we observe the $\eta$ Car system in the orbital plane. However, from the structure of the Homunculus we believe the orbital plane is tilted by $\sim48^\circ$ from an edge-on view (an inclination angle of $i = 42^\circ$; Smith 2002). It appears that the model of Falceta-Goncalves et al. (2005) cannot work for an inclined system, although they do not give enough details for a conclusive assessment.

In the model proposed here the collapse of the wind collision region onto the secondary is the reason for the sharp decline in the X-ray intensity. Figure 1 shows the proposed orbit in the top panel and the stagnation point region in the bottom panel. Most of the $2-10$ keV X-ray emission outside the X-ray minimum comes from the shocked secondary wind (e.g., Pittard & Corcoran 2002). When the two stars approach periastron, the ratio of the accretion radius of the secondary $R_{\text{sec}}$ to the distance of the secondary from the stagnation point $D_{22}$ increases substantially, reaching a value of $\sim0.1$ at 40 days before periastron passage and $\sim0.4$ at periastron. At the same time the free-fall time of cold postshock primary wind mass elements becomes shorter, and it is no longer much larger than the flow time out of the stagnation.
region. As pointed out by Soker (2005b), the evolution of these two ratios strongly suggests that the stagnation point region collapses onto the secondary ~0–40 days before periastron passage, and the secondary starts to accrete the primary wind. This process is assumed to prevent the secondary from accelerating its wind, and hence shutting down the main source of X-ray emission, causing the minimum. The gradual collapse of the winds is associated with a rise in density and column density, which is responsible for the temporary hardening of the X-ray spectrum before minimum (Corcoran 2005).

As distance increases after periastron passage, the accretion rate decreases, allowing the acceleration zone of the secondary wind to build up again ~60–70 days after periastron passage. As a result the wind reappears and so does the X-ray emission. Right after the minimum the binary separation is still small and X-rays are emitted from deep in the primary wind, again through a relatively high column density, which again results in a (temporarily) hard spectrum. The progressive hardening and softening of the X-ray spectrum just outside the deep minimum is consistent with an absorption effect. Later on, as the secondary emerges from periastron, the radiation becomes softer, returning to its normal orbital spectrum.

At ~30 days into the minimum the X-ray radiation becomes harder for ~25 days. This is accompanied by an increase in X-ray luminosity; this is observed in both the 1997.9 and the 2003.5 minima (Corcoran 2005). The small rise in X-ray luminosity ~30 days into the minimum might be explained with our model in the following, somewhat speculative way. For a short time during the accretion phase the specific angular momentum of the accreted mass rises, such that the accretion onto the secondary is concentrated in the equatorial plane and part of the accreted mass is blown along the polar directions. This moderate-velocity polar outflow runs into the ambient primary wind and is subsequently shocked, which results in weak extra X-ray emission during the minimum.

Based on the results of Soker (2005b) and observations, we consider the X-ray minimum to start ~0–30 days (phase ~0.015–0) before periastron passage. In this paper we take the accretion to start 20 days before periastron (orbital phase 0.01) and to end 60 days after periastron (orbital phase 0.03). The X-ray minimum itself starts a little after accretion starts and lasts several days less than the assumed 80 day accretion period. Any starting phase of the X-ray minimum in the range of approximately ~0.04–0 is acceptable in our model.

3. COLLIDING WINDS

3.1. Structure of the Flow

3.1.1. Binary Orbit

The binary system parameters are as in Paper I. The stellar masses are $M_1 = 120$ and $M_2 = 30 M_\odot$, the eccentricity is $e = 0.9$, and orbital period 2024 days; hence, the semimajor axis is $a = 16.64$ AU and periastron occurs at $r = 1.66$ AU (for the definition of these and other terms, see Table 1). The mass-loss
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 $r$ (in AU) and the relative orbital speed of the two stars $v_{\text{orb}}$ (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: Radiative cooling time of the shocked secondary’s wind near the stagnation point $\tau_{\text{cool}}$ (top line), and the typical time $\tau_{f2}$ for the shocked secondary wind to flow out of the winds interaction zone (bottom line). Third row: Distance of the stagnation point from the secondary when the gravity of the secondary is included $D_{g2}$ (top line; see Fig. 1) and the Bondi-Hoyle accretion radius of the secondary star $R_{\text{acc2}}$ (bottom line). Bottom row: Velocity of the primary wind relative to the stagnation point $v_{\text{wind1}}$ (thick line) and the ratio $v_{\text{wind1}}/\tau_{\text{cool2}}$ (thin line). For more detail, see Paper I.

In the second row the thick line represents the typical time $\tau_{f2}$ for the shocked secondary wind to flow out of the shocked region (winds interaction zone), while the thin line depicts the radiative cooling time of the shocked secondary wind $\tau_{\text{cool2}}$. In the third row of Figure 2 the distance of the stagnation point from the secondary, $D_{g2}$, and the Bondi-Hoyle accretion radius of the primary wind by the secondary star $R_{\text{acc2}}$ are drawn. In the fourth row the velocity of the primary wind relative to the stagnation point $v_{\text{wind1}}$ is depicted by the thick line, while the thin line represents the ratio of $\tau_{f2}/\tau_{\text{cool2}}$. For more detail on these quantities see Paper I.

### 3.1.2. Colliding Wind Geometry

The colliding wind region is schematically drawn in Figure 3. The winds from the two stars produce two respective shock waves. The shocked gas flows away from the stagnation point along the contact discontinuity—the surface where the two momentum fluxes exactly balance each other and which separates the two postshock flows. The gas is heated by the shock waves to temperatures of $\sim 10^7-10^8$ K and generates X-ray emission. The distances $D_1$ and $D_2$ from the binary components to the stagnation point are calculated from the equation

$$\rho_1 v_{\text{wind1}}^2 = \rho_2 v_{\text{wind2}}^2.$$  \hfill (2)

Here $\rho_1$ and $\rho_2$ are the preshock densities of the two winds, $v_{\text{wind1}}$ is the preshock speed of the primary wind relative to the stagnation point (assumed to move with the secondary), and $v_{\text{wind2}}$ is the preshock speed of the secondary wind. All quantities are calculated at the stagnation point. Using equation (2) and the relation $D_1 + D_2 = r$, where $r$ is the distance between the stars, the distances $D_1$ and $D_2$ can be found. When the effect of the gravity of the companion is included, the newly calculated distance of the stagnation point to the secondary, $D_{g2}$, decreases slightly. At periastron $D_{g2} \approx 0.8D_2$. 

rates are $\dot{M}_1 = 3 \times 10^{-4} M_\odot$ yr$^{-1}$ and $\dot{M}_2 = 10^{-5}$ $M_\odot$ yr$^{-1}$. Ishibashi (2001) proposed that $M_2 \approx 40 M_\odot$. A more massive secondary implies even more gravity by the secondary, favoring accretion near periastron passage even further. According to Smith et al. (2003) the total primary mass-loss rate is higher than $3 \times 10^{-4} M_\odot$ yr$^{-1}$ but with a higher mass-loss rate in the polar directions and lower mass-loss rate toward the equatorial plane. Since most of the X-ray emission comes from regions near the equatorial plane, we take the mass-loss rate as quoted above. The primary wind velocity profile is

$$v_1 = 500[1 - (0.4 \text{ AU}/r_1)] \text{ km s}^{-1},$$  \hfill (1)

where $r_1$ is the distance from the center of the primary.

The secondary wind speed is taken to be $v_{\text{wind2}} = 3000$ km s$^{-1}$. The orbital separation $r$, the relative orbital velocity of the two stars $v_{\text{orb}}$, and the angle $\theta$ of the position of the secondary relative to the semimajor axis during periastron (see Fig. 1) are plotted in the top row of Figure 2.
The asymptotic half opening angle \( \phi_a \) of the contact discontinuity, which is defined in Figure 3, is (e.g., Eichler & Usov 1993):

\[
\phi_a \sim 2.1 \left( 1 - \frac{\beta^{4/5}}{4} \right) \beta^{2/3} \simeq 1 \simeq 60^\circ,
\]

where

\[
\beta \equiv \left( \frac{M_2 v_2}{M_1 v_1} \right)^{1/2}.
\]

For the parameters used here \( \beta \simeq 0.4 \).

### 3.2. X-Ray Emission

Usov (1992) gives an expression for the X-ray luminosity \( L_X \) later used by Ishibashi et al. (1999). However, this expression cannot be applicable here, because if we substitute the velocities and mass-loss rates typical of \( \eta \) Car at an orbital separation of \( r < 4 \) AU, the total X-ray emission is more than the total kinetic power of the secondary wind. This is indeed a problem in the model of Falceta-Goncalves et al. (2005). We therefore need to derive a different expression for the X-ray luminosity. The expected increase in the intrinsic luminosity with orbital separation is \( 1/r \) (Usov 1992), much more than that observed. The more moderate increase of \( L_X \) with decreasing orbital separation was attributed by Ishibashi et al. (1999) and Corcoran et al. (2001a) to an increase in absorption accompanying the increase in X-ray intrinsic luminosity with decreasing radius. We assume the same.

Usov (1992) uses a cooling function \( \Lambda \propto T^{1/2} \) that is appropriate for hot gas at \( T \simeq 2 \times 10^7 \) K. Since the shocked primary wind of \( \eta \) Car is much cooler, we need to use a more appropriate cooling function. We define \( F_{\text{AB}}(T_s) \) to be the fraction of the X-ray flux emitted by the gas shocked to the temperature \( T_s \) in the range \( A–B \) (in keV), out of the total flux emitted between 0.01–100 keV. The X-ray luminosity between 2–10 keV depends strongly on \( T_s \). Using the APEX plasma database (Smith et al. 2001), we have calculated \( F_{\text{AB}}(T_s) \) for a range of temperatures corresponding to preshock wind velocities (relative to the shock) of 500 km s\(^{-1} \) \( \leq v \leq 2000 \) km s\(^{-1} \). We then fitted \( F_{\text{AB}}(T_s) \) with a parametric form of the bremsstrahlung emissivity function to obtain

\[
F_{210} = 0.65 \exp \left( -2.325 \times 10^7 / T_s \right)
= 0.65 \exp \left[ -6.6 \left( \frac{v}{500 \text{ km s}^{-1}} \right)^2 \right].
\]

#### 3.2.1. Primary Wind

Close to the X-ray minimum, which occurs near periastron passage, the shocked primary wind is very dense and its cooling time is much shorter than the flow time (e.g., Pittard & Corcoran 2002; Soker 2003). Therefore, we assume that any postshock primary wind material cools instantaneously by emitting all of its thermal energy in the form of radiation. The emitted spectrum is taken to be that of gas at the postshock temperature at the stagnation point. Since the shock front is oblique, the shock velocity and the temperature away from the stagnation point, in reality, are lower. Moreover, as the shocked primary wind loses energy, its temperature decreases. Consequently, on average, the spectrum is typical of temperatures lower than the intermediate postshock temperature. For temperatures typical of the shocked primary wind, a small decrease in the emitting gas temperature substantially reduces its contribution in the band above 2 keV. Therefore, it should be clear that our treatment overestimates the contribution of the primary wind to the X-ray emission in the 2–10 keV band.

From the shape of the shock wave for our relevant parameters (§3.1) we estimate that the primary’s wind segments with \( \phi \simeq 40^\circ \) pass through a strong enough shock and heat to a high enough temperature to contribute to the X-ray emission in the 2–10 keV band (see Fig. 3). This implies that the rate of mass entering the strong shock region is \( k_1 M_1 = (1 - \cos 40^\circ) M_1 / 2 = 0.1 M_1 \). The preshock velocity of this mass is about equal to the relative velocity of the primary wind to the stagnation point, \( v_{\text{wind1}} \), and it is plotted in the fourth row of Figure 2. The calculation of \( v_{\text{wind1}} \) includes the wind velocity relative to the primary and the relative orbital velocity between the two stars. The acceleration of the primary wind is included by taking a lower value of \( v_1 \) close to the primary star, which increases with distance from the primary up to a terminal velocity at large distances of \( v_1 = 500 \) km s\(^{-1} \) (see Paper I for more details).

The relevant velocity for the primary wind at the stagnation point is \( v_{\text{wind1}} \simeq 500 \) km s\(^{-1} \) (see the fourth row of Fig. 2). For a preshock velocity of \( v_{\text{wind1}} = 400, 500, \) and 600 km s\(^{-1} \), corresponding to postshock temperatures of \( T_s = 2.2 \times 10^6, 3.5 \times 10^6, \) and \( 5 \times 10^6 \) K, we find \( F_{210} = 2.1 \times 10^{-4}, 1.1 \times 10^{-3}, \) and \( 5.5 \times 10^{-3} \), respectively. In other words, only a small fraction of the energy is emitted in the X-ray band. The contribution of the primary wind is therefore:

\[
L_{X1} \simeq 2.4 \times 10^{33} \left( \frac{k_1}{0.1} \right) \left( \frac{F_{210}}{0.001} \right) \left( \frac{M_1}{3 \times 10^{-4} M_\odot \text{ yr}^{-1}} \right) \times \left( \frac{v_{\text{wind1}}}{500 \text{ km s}^{-1}} \right)^{-2} \text{ ergs s}^{-1}.
\]

The X-ray luminosity \( L_{X1} \) as a function of orbital phase is plotted as a thin solid line in the top row of Figure 4. The dotted line in the same figure shows the X-ray flux in the entire 0.01–10 keV range, namely, taking \( F_{210} = 1 \) in equation (6) (note the different scaling there). Considering photoelectric absorption (mostly at 2–3 keV), the observed luminosity reduces to \( L_{X1} \simeq 10^{33} \) ergs s\(^{-1} \) during most of the orbit, with a possible intrinsic contribution of up to \( L_{X1} \simeq 10^{34} \) ergs s\(^{-1} \) as the system approaches periastron, where \( v_{\text{wind1}} \) is high (Fig. 2). On the other hand, close to periastron, X-ray absorption is at its peak. The most important point, however, is that the primary wind is very slow, and therefore the postshock temperature is low and most of the X-ray emission is soft with a negligible contribution above \( \sim 3 \) keV (Pittard & Corcoran 2002). Furthermore, the X-ray emission below 3 keV is more strongly absorbed, and one must conclude that during most of the orbit the major fraction of the observed 2–10 keV X-ray emission comes from the secondary wind. (We emphasize again: the primary wind does emit in the 2–3 keV band, as seen in the top row of Fig. 4; however, absorption in this band is high and most of this flux never reaches the observer. In addition, for reasons stated at the beginning of this subsection, the treatment here overestimates the contribution of the shocked primary wind to the 2–10 keV band.)

#### 3.2.2. Secondary Wind

The secondary wind is much faster (\( \sim 3000 \) km s\(^{-1} \)) than the primary wind, and the relevant temperature range is therefore \( \sim 0.5–1.3 \times 10^6 \) K, for which we find an average value of \( F_{210} = 0.46 \). At the high temperatures and low densities typical of the secondary wind, and in contrast with the primary wind, the radiative cooling time of the postshock secondary wind material...
The solid line shows the intrinsic luminosity of the shocked primary wind only in the 2–10 keV band $L_{X1}$ (eq. [6]), and the thick line shows the intrinsic luminosity of the shocked secondary wind only in the 2–10 keV band $L_{X2}$ (eq. [8]) with $k_3 = 2$, both in units of $10^{36}$ ergs s$^{-1}$. Most of the contribution to $L_{X1}$ is in the 2–3 keV band, explaining its large attenuation. Middle row: X-ray transmission factor through the primary wind for the 2–10 keV band, calculated with XSPEC. The solid lines show the transmission using the column density calculated by eq. (9) (or eq. [10]), while the dashed lines show the transmission for an inclination $i = 42^\circ$ and periastron orientation $\omega = 180^\circ$ (eq. [11]); $k_y = 0.5$ in all cases. The two top lines show the transmission of the X-ray emitted by the shocked secondary wind, while the two bottom lines show the transmission of the X-ray emitted by the shocked primary wind. These lines emphasize the small difference between the two absorbing gas geometries considered here and the high absorption of the X-ray emitted by the shocked primary wind. Bottom row: Calculated X-ray luminosities of the shocked two winds in the 2–10 keV band when attenuation by eq. (9) is included; the thin and thick lines, respectively, represent the attenuated X-ray luminosities of the shocked primary and secondary winds. These quantities should be compared with observations outside the X-ray minimum. The horizontal line in each of the two panels in the bottom row marks the time during which, according to our model, the secondary wind does not exist or is highly suppressed. It spans the time starting a little before the minimum and ending a little after the minimum, phase $-0.01$ to $0.035$.

$\tau_{\text{cool2}}$, is much longer than the flow time out of the wind-collision region. This implies that the shocked secondary wind region is large, because the shock front is at a large distance from the contact discontinuity, e.g., as Pittard et al. (2002) simulate for the massive binary system WR 147. This is schematically drawn in Figure 3.

For the purpose of studying the X-ray minimum, it is adequate to take the contribution of the secondary wind to the X-ray luminosity as follows. We assume that about half of the mass blown by the secondary star is shocked in a shock front perpendicular to the wind velocity. We then assume that the X-ray luminosity is determined by how much of the thermal energy of the shocked gas is radiatively emitted as X-rays before the gas cools adiabatically. The radiative emission lasts for timescales of the order of the radiative cooling time $\tau_{\text{cool2}}$, while adiabatic cooling takes place on the flow timescale $\tau_{f2}$. Therefore, the total X-ray energy emitted is a fraction $k_2 \tau_{f2}/\tau_{\text{cool2}}$ of the thermal energy of the postshock gas. Here $\tau_{f2} \equiv D_{\Delta}/v_2$ is the characteristic flow time of the shocked wind out of the interaction region, where $D_{\Delta}$ is the distance of the stagnation point from the secondary. The value of $\tau_{f2}/\tau_{\text{cool2}}$ is plotted in the bottom row of Figure 2. The contribution of the shocked secondary wind material to the X-ray luminosity in the energy range $A–B$ can therefore be written as:

$$L_{X2} = \frac{1}{4} M_2 v_2^2 F_{AB}(T_2) \frac{\tau_{f2}}{\tau_{\text{cool2}}} k_2. \quad (7)$$

Substituting typical values in equation (7) gives for the 2–10 keV range:

$$L_{X2} = 6.5 \times 10^{34} \left( \frac{F_{210}}{0.46} \right) \left( \frac{M_2}{10^{-5} M_\odot} \right) \left( \frac{v_2}{3000 \text{ km s}^{-1}} \right)^2 \times \left( \frac{\tau_{f2}/\tau_{\text{cool2}}}{0.01} \right) k_2 \text{ ergs s}^{-1}. \quad (8)$$

As shown by Pittard & Corcoran (2002), Kelvin-Helmholtz instability modes develop on the contact discontinuity. This instability indirectly enhances the X-ray emission by two effects: (1) the large “tongues” slow down the velocity of the shocked secondary wind, increasing the effective outflow time of shocked secondary wind near the contact discontinuity; and (2) the Kelvin-Helmholtz instability mixes hot, shocked secondary-wind gas with cool, dense shocked primary-wind gas. This mixing causes a small fraction of the shocked secondary wind to cool to low temperatures, releasing most of its energy (and not just a fraction of $\tau_{f2}/\tau_{\text{cool2}}$). These effects, as well as our ignorance of the exact values of the secondary wind speed and mass-loss rate, are incorporated in the parameter $k_3 > 1$, which we fit to match the observed X-ray luminosity. The luminosity $L_{X2}$ according to equation (8) with $k_3 = 2$ is plotted as the thick line in the top row of Figure 4, where it can be seen to dominate $L_{X1}$. The same result holds in Figure 5.

3.3. X-Ray Absorption

If the X-rays from $\eta$ Car are observed through the primary wind, absorption is inevitable, but its magnitude depends on the wind parameters, flow parameters, and system orientation. First,
the soft (2–5 keV) and hard (7–10 keV) bands are generally emitted from different regions. The hard X-rays come mainly from regions closer to the stagnation point, where the shock wave is strong and postshock temperature is $\sim 10^8$ K. This region is marked by upper-case bold-face X’s in Figure 6. Away from the stagnation point, the gas cools adiabatically and the X-rays become softer. Thus, we expect the column density toward the hard X-ray regions to be somewhat higher than that toward the soft X-ray region. Although the soft X-ray regions are less obscured, absorption of soft X-rays is high enough to explain the X-ray hardening effect just prior and just after the X-ray minimum. As our main goal is to explain the X-ray light curve, including the minimum, in this work we do not model the 2–5 and 7–10 keV bands separately, but treat the entire 2–10 keV band together.

Second, for some orientations and some of the time, the X-ray emission could be observed through the tenuous secondary wind instead of through the bulk of the dense primary wind, which would imply a lower column density. These orientations are represented in Figure 6 by the narrow arrows. For observer $A$ in Figure 6, this occurs after the X-ray minimum, while for observer $B$ this occurs before the X-ray minimum. The winds collision model implies significant absorption of the X-ray emission during the X-ray maximum just before the X-ray minimum. This rules out the orientation of observer $B$ in Figure 6. (We note that Falceta-Goncalves et al. [2005] and Abraham et al. [2005] argue to the contrary; however, they require the line of sight to be through the orbital plane, which is in contradiction with the structure of the Homunculus, from which the orbital plane is deduced to be tilted by $\sim 48^\circ$ from an edge-on view [Smith 2002].)

At other phases, the X-ray emission is partially absorbed by the much higher column density of the primary wind (wide arrows, Fig. 6). To estimate the importance of the geometry, we assume that our line of sight is at $\sim 45^\circ$ to the orbital plane (Smith 2002).

Let $\phi_\alpha$ be the asymptotic half opening angle of the contact discontinuity (Fig. 3). At an angle of $45^\circ$, the projection of the half opening angle on the orbital plane, $\phi_{45}$, is given by $\tan \phi_{45} = (\tan^2 \phi_\alpha - 1)^{1/2}$. For $\phi_\alpha = 45^\circ$, $60^\circ$, and $70^\circ$, we find $\phi_{45} = 0, 54.7$ and $68.7^\circ$, respectively. Namely, if $\phi_\alpha \leq 45^\circ$ we will always observe the X-ray emitting gas through the primary wind. For the typical parameters used here $\phi_\alpha \approx 60^\circ$. Hence, $\phi_{45} \approx 55^\circ$. Therefore, during a substantial fraction of the orbital orbit we might be observing the X-ray emission through the fast secondary wind (i.e., low column density). We do not elaborate on this effect in this paper, because there are too many unconstrained parameters to consider: (1) there is no consensus yet on the orbital orientation (e.g., Ishibashi 2001; Smith et al. 2004); (2) the nonspherical mass-loss geometry from the primary, with denser material in the polar direction (Smith et al. 2003); (3) the cone depicted in Figure 6 actually has a spiral structure due to the orbital motion. Namely, the winds interaction region is winding around the binary system as it flows outward, like the dust in the interacting winds binary system WR 98a (Monnier et al. 1999). This includes the cool, dense postshocked primary wind material. Hence, some absorption is expected. Each one of these effects introduces at least one free parameter. At this stage of the development of the model, there is no point in exploring this immense parameter space. Instead, we notice that most of the absorption occurs by the primary wind in regions close to the stagnation point and use this to build a simple geometry that includes the main features of the primary wind near the stagnation point. For comparison, we also consider the absorption for the orbital orientation proposed by Smith et al. (2004).

Following the discussion in the previous paragraph, in our calculations here we take, as our basic geometrical structure, the momentary column density through the primary wind to be that to the stagnation point, along a line of sight perpendicular to the line connecting the two stars. Taking $y$ to be the coordinate perpendicular to the line connecting the two stars and using $M_1 = n_1 \mu m_H 4 \pi r_1^2 v_1$, where $r_1 = D_1 + y^2$, this hydrogen column number density can be expressed as:

$$N_{\text{HII}} = 0.43k_g \int_0^\infty \frac{\dot{M}_1}{4\pi(D_1^2 + y^2)\mu m_H} dy,$$

where $\mu m_H$ is the mean mass per particle in a fully ionized gas, $n$ is the total number density, 43% of which are hydrogen nuclei, and $k_g$ is a factor that depends on the geometry and orientation. It is used here to compensate for the many unknown parameters. Performing the integral under the assumption of constant $v_1$, and substituting typical values near periastron gives

$$N_{\text{HII}} = 1.3 \times 10^{24} \frac{\dot{M}_1}{3 \times 10^{-4} M_0 \text{ yr}^{-1}} \left( \frac{v_1}{500 \text{ km s}^{-1}} \right)^{-1} \times \left( \frac{D_1}{1 \text{ AU}} \right)^{-1} k_g \text{ cm}^{-2}.$$
At apastron the binary separation is 31.6 AU, and the stagnation point is a distance \( D_1 = r - D_{1/2} \approx 22 \) AU from the primary, which from equation (10) yields a column density at apastron of \( N_{H1p} \approx 6 \times 10^{22} \) cm\(^{-2}\). Observations indicate a column density of \( N_{H1} \approx 3 \times 10^{22} \) cm\(^{-2}\) during most of the orbit (Ishibashi et al. 1999; Corcoran et al. 2001b). We therefore take \( \kappa_g = 0.5 \) in this work, although this value can be larger when the system emerges from the X-ray minimum.

In the second row of Figure 4 we plot the calculated X-ray transmission factor using the column density given by equation (10), with \( \kappa_g = 0.5 \). In the third row of Figure 4, we plot the X-ray luminosity (top row), but now including absorption. It can be seen that after absorption \( L_{X2} \) is even more dominant over \( L_{X1} \). The geometry used to calculate the column density in equations (9) and (10) was the simplest one, and we chose it because of the many uncertainties. We could use another geometry, such as that suggested by Smith et al. (2004), in which the inclination angle \( i \) of our line of sight to a norm to the orbital plane is 42°, and the secondary is on the far side of the primary before reaching periastron (the semimajor axis is perpendicular to our line of sight). In this geometry, the line of sight to the stagnation point practically always goes through the primary wind, unlike some orientations shown in Figure 6, which was plotted for the \( i = 90^\circ \) case. The reason is that after periastron passage, the bow shock of the secondary wind is fully recovered only at phase \( \approx 0.04 \) (see Fig. 9), when the orbital angle is \( \theta \approx 140^\circ \) (Fig. 2, top row).

It can be shown that when the semimajor axis is perpendicular to the line of sight, then the column density of the primary wind as a function of the binary azimuthal angle \( \theta \) and inclination angle \( i \) is

\[
N_{H1}(i) = \frac{0.43k_\beta n_1}{4\pi m_{H1}(r_D_1)} \left( \frac{1 + \sin^2 i}{\tan^2 2i} \right)^{1/2} \sin \theta \tan i \left( \frac{1 + \cos^2 \theta \tan^2 i}{\sin^2 2i} \right)^{1/2},
\]

where \( \theta \) here is negative before periastron, zero at periastron, and positive after periastron (see Fig. 1). The ratio \( N_{H1}(i = 42^\circ)/N_{H1} \) as a function of the orbital phase is plotted in the top row of Figure 5 for a fixed value of \( \kappa_g \). In the second row of Figure 5, the absorbed X-ray luminosity is plotted with \( \kappa_g = 0.5 \), but with the column density given by equation (11). The bottom row of Figure 5 shows the ratio of this X-ray luminosity to that shown in Figure 4 for the \( i = 90^\circ \) case. It can be seen that the only significant difference arises near periastron, where according to the proposed model, the secondary wind does not exist. We conclude that the difference between the two geometries is practically small and that the X-ray luminosity is not very sensitive to the inclination angle \( i \) (as long as \( i \) is not too large). The emission of the shocked primary wind is reduced substantially; but in our model most X-ray emission in the 2–10 keV bend is attributed to the shocked secondary wind anyway. To summarize, although the binary orientation assumed in calculating the absorbing column density used in equation (9) (or eq. [10]) is not the exact orientation of the binary system in \( \eta \) Car, it is adequate for the main goal of this paper, as it emphasizes the weak dependence of the results on the exact geometry. The orientation used in deriving equation (9) has the advantage that it does not depend on the orientation of the periastron \( (\omega) \) with respect to our line of sight. The small difference between the results obtained for the two geometries used here demonstrates that our model is sensitive neither to the orientation of the periastron nor to the inclination of the orbital plane (as long as \( i \) is not too large).

The X-ray minimum lasts \( \approx 70 \) days, the time from the beginning of the decline to the end of the rise from the minimum is \( \approx 120 \) days. According to our proposed model, the decline starts \( \approx 40 \) days before periastron passage \( (r \approx 5.3 \) AU\), when the absorbing column density increases and blobs from the postshock primary wind start to be accreted. Bondi-Hoyle--type accretion starts \( \approx 20 \) days before periastron \( (r \approx 3.3 \) AU\). The minimum starts \( \approx 10 \) days before periastron and ends \( \approx 60 \) days after periastron \( (r \approx 7.2 \) AU\), taking another \( \approx 20 \) days to fully rise back \( (r \approx 8.7 \) AU and \( D_1 \approx 6 \) AU). At the beginning of the decline \( (40 \) days before periastron), where the absorption effect is most significant, the column density is \( N_{H1} \approx 1.5 \times 10^{23} \) cm\(^{-2}\) (with \( \kappa_g = 0.5 \)). At this column density, 29%, 83%, and 71% of the flux is transmitted through the primary wind in the 2–5, 7–10, and 2–10 keV bands, respectively. The rest is absorbed by the primary wind. This accretion flow is of the Bondi-Hoyle type (Fig. 7). The very high Mach number creates an accretion shock. Because the cooling time of the postshocked primary wind material is very short near periastron, about 1% of the time (e.g., Pittard & Corcoran 2002; Soker 2003), the accretion flow can be treated as isothermal. Such a flow has been simulated (e.g., by Ruffert 1996). In the following we distinguish between segments of the primary wind impacting the secondary directly and segments that go around the secondary toward the accretion column (behind the secondary) along a curved trajectory. The basic flow structure is as follows.

The accretion column—the elongated region behind the secondary enclosed by the shocked accreted matter—is narrow (see Fig. 7), and the Mach number of the preshocked primary wind is...
high. As an approximation, we take the velocity of the primary wind just before it hits the shock wave as the velocity it would have reached on the accretion line—the symmetry line behind the accreting body. Using the definition of $R_{\text{acc}2}$

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

(12)

this velocity can be found from energy conservation to be

$$v_{c1} = v_{\text{wind}1} \left( 1 + \frac{R_{\text{acc}2}}{z} \right)^{1/2},$$

(13)

where $z$ is the distance from the secondary along the accretion line. The component in the perpendicular direction with respect to the accretion line of the primary-wind velocity before hitting the accretion-column shock $v_{c1}(z)$ is found from angular momentum conservation. Assuming the stream lines of the primary wind to be parallel at large distances from the secondary, we can write

$$v_{c1} = v_{\text{wind}1} \left( \frac{R_{\text{acc}2}}{z} \right)^{1/2}. \quad (14)$$

Simulations of high Mach number, nearly isothermal flows (e.g., Ruffert 1996) show that the flow structure is very similar to that of the classical Bondi & Hoyle (1944) accretion flow. The total accreted mass can be approximated for our purposes as follows:

$$\dot{M}_{\text{acc}2} \simeq \pi \rho_0 R_{\text{acc}2}^2 v_0 = 0.01 M_1 \left( \frac{r}{1 \text{ AU}} \right)^{-2} \left( \frac{R_{\text{acc}2}}{0.2 \text{ AU}} \right)^2 \times \left( \frac{v_1}{500 \text{ km s}^{-1}} \right)^{-1} \left( \frac{v_{\text{wind}1}}{500 \text{ km s}^{-1}} \right), \quad (15)$$

where we take the density far upstream to be equal to the density at the location of the secondary $\rho = \rho_1 \left( r / (4\pi r^2 v_1) \right)$ and the upstream speed is $v_0 = v_{\text{wind}1}$. The mass-accretion rate as a function of the orbital phase is plotted in the top panel of Figure 8, only during the orbital phases in which, according to our model, accretion occurs. Part of the accreted mass hits the accretion line in the range $z_{\text{min}} \leq z \leq R_{\text{acc}2}$, while another fraction of the incoming mass, $\sim (z_{\text{min}} / R_{\text{acc}2}) \dot{M}_{\text{acc}2}$, hits the accretion shock close to the accreting body. The rate of mass hitting the accretion column per unit length is (Bondi & Hoyle 1944)

$$\frac{d\dot{m}_z}{dz} = \pi \rho_1 (r) v_{\text{wind}1} R_{\text{acc}2}. \quad (16)$$

4.2. X-Ray Emission due to Accretion

The material with impact parameter $b \leq R_{\text{acc}2}$ will be attracted by the secondary’s gravity and flow toward the accretion line, passing through the shock almost perpendicular to the shock front. This gas is subsequently heated to a postshock temperature appropriate for its incoming velocity, which is given by equation (13). We assume this gas to be completely thermalized. The thermal energy gained per unit length is

$$\frac{d\dot{E}_z}{dz} = \frac{1}{2} \frac{d\dot{m}_z}{dz} v_{c1}^2. \quad (17)$$

![Figure 8](image_url)

**FIG. 8.—**Quantities related to the X-rays emitted by primary wind material accreted onto the secondary during the X-ray minimum. **Top:** Accretion rate $\dot{M}_{\text{acc}2}$ (eq. [15]). **Second panel:** The thick line shows the X-ray luminosity in the 2–10 keV band of material accreted by the secondary through the accretion column $L_{\text{X}2}$ (eq. [18]). The thin line shows the X-ray luminosity in the 2–10 keV band of material accreted directly by the secondary $L_{\text{X}1}$, eq. (20). **Third panel:** Total luminosity of the shocked accreted gas, namely, taking $F_{210} = 1$ in eqs. (18) and (20). **Bottom:** Luminosities as shown in the second panel, but with absorption included. Note that $L_{\text{X}2}$ is completely absorbed.

We further assume that the cooling time is very short and that all this energy is radiated away, giving rise to an X-ray luminosity in the 2–10 keV range of

$$L_{\text{X}1} = \frac{1}{2} \int_{z_{\text{min}}}^{R_{\text{acc}2}} F_{210}(z) \frac{d\dot{m}_z}{dz} v_{c1}^2 dz, \quad (18)$$

where $F_{210}$ depends on $z$ through the dependence of the postshock temperature $T_s$ on $v_{c1}$.

The material impacting the secondary directly is also assumed to pass through the accretion shock perpendicular to the shock front at a distance $z_{\text{min}}$ from the secondary with a speed given by equation (13), with $z = z_{\text{min}}$. We take $z_{\text{min}}$ to be approximately the radius of the secondary star $R_2 \simeq 22 R_\odot \simeq 0.1 \text{ AU}$ (Verner et al. 2005). This contribution to the X-ray luminosity is therefore

$$L_{\text{X}2} = \frac{1}{2} \frac{z_{\text{min}}}{R_{\text{acc}2}} \dot{M}_{\text{acc}2} (F_{210} v_{c1}^2)_{z=z_{\text{min}}}. \quad (19)$$

For $z_{\text{min}} = 0.1 \text{ AU}$ and near periastron we find for the preshock speed $v_{c1} \simeq 900 \text{ km s}^{-1}$, and for the postshock temperature $T_s \simeq 10^7 \text{ K}$ and $F_{210} \simeq 0.08$. We scale equation (19) to find

$$L_{\text{X}2} \simeq 4 \times 10^{34} \left( \frac{z_{\text{min}}}{R_{\text{acc}2}} \right) \left( \frac{0.5}{0.08} \right) \left( \frac{v_{c1}}{1000 \text{ km s}^{-1}} \right)^2 \text{ergs s}^{-1}. \quad (20)$$

The two contributions to the intrinsic luminosity in the 2–10 keV band are plotted in the second panel of Figure 8. The third panel shows the total luminosity of the two contributions, namely, in the 0–$\infty$ keV band. This intrinsic luminosity is much larger than the luminosity observed during the minimum. However, during the minimum, the column density through the primary
wind is considerable \((N_{\text{HI}} > 10^{23} \text{ cm}^{-2})\) and it totally absorbs the X-ray accretion source. We find that accretion cannot even account for the weak X-ray flux during the minimum as demonstrated below.

The mass hitting the accretion column at \( z \gtrsim R_{\text{acc}} \) is not accreted by the secondary. It passes through an oblique shock at an angle of \( \theta \) between the inflow direction and the shock front with \( \tan \theta \approx v_{\text{wind1}} / v_{\text{wind}} \). This postshock temperature is less than the postshock temperature of the primary wind when accretion does not occur. Hence this gas emits soft X-rays, which are totally absorbed by the primary wind. We do not consider its contribution in this paper.

As the primary wind flows toward the accreting body its density increases. Therefore, the column density to the X-ray emitting regions here is somewhat larger than that in equation (10). In addition, the density in the accretion column is very large, making it completely opaque. Therefore, only about half of the intrinsic luminosity can be observed. For the emerging radiation, we take the column density of equation (10), but with \( r \approx 2 \text{ AU} \) replacing \( D_1 \) and a lower value for \( v_1 \):

\[
N_{\text{HI}} = 10^{24} \frac{\dot{M}_1}{3 \times 10^{-4} M_\odot \, \text{yr}^{-1}} \left( \frac{v_1}{400 \text{ km s}^{-1}} \right)^{-1} \left( \frac{2 \text{ AU}}{r} \right)^{-1} \text{ cm}^{-2}.
\]

(21)

The expected X-ray emission given by equations (18) and (20), including absorption (eq. [21]), is plotted in the bottom panel of Figure 8. This radiation is soft, with most of the contribution to the 2–10 keV band coming from radiation in the 2–3.5 keV energy range. These simplified estimates clearly show that the soft emission due to accretion, together with the high column density it needs to traverse, rules out this component as a candidate for the weak X-ray emission during the minimum. A different alternative is pursued in the next subsection.

4.3. X-Ray Emission during the Minimum

The main goal of this paper is to show that the X-ray behavior up to the minimum can be explained by the assumed orbital parameters and by the variation of the radiative cooling timescale of the secondary’s wind relative to its flow timescale. This subsection, on the other hand, is more speculative. We put forward a suggestion for the nature of the very weak X-ray emission during the minimum. We already showed that X-ray emission due to accretion cannot explain even the weak X-rays observed during the minimum. Because of the speculative nature of this subsection, the treatment is more qualitative. We suggest that the weak X-ray emission during the minimum, with its variation in hardness and intensity (see property 3 in § 2.1), could come from two different components: (1) an almost constant weak contribution from old shocked secondary wind, termed here the residual component; and (2) a possible temporary collimated outflow resulting from accretion of mass with relatively high specific angular momentum. These two components are treated in the next two subsections.

We do not attempt to refute other contributions, e.g., other stars in the field and, more importantly, the possibility that most of the X-ray emission during the X-ray minimum comes from reflected light by the nebula (Hamaguchi et al. 2005); time delay ensures that the X-ray emission is observed during the minimum (Corcoran et al. 2004). We only point out that other processes should be considered as well, before further observations and calculations reveal the most significant process contributing to the X-ray emission and its variability during the minimum.

4.3.1. Residual Emission

We propose that some fraction of the soft and weak X-ray emission observed even during the minimum results from previously shocked secondary wind segments. This residual emission exists basically along the entire orbit, but it can be noticed only after the main X-ray source has been shut down.

During the winds collision phase, the shocked secondary wind forms a hot bubble flowing away from the secondary in what we term the secondary hot tail, winding around the binary system as it flows outward, like the dust in the interacting-winds binary system WR 98a (Monnier et al. 1999). The shocked secondary wind is engulfed eventually by the slower primary wind and expands with it to large distances at a speed of \( v_1 \). We wish to crudely estimate the contribution of such a bubble to the X-ray emission. This crude estimate is by no means a replacement for a full three-dimensional numerical simulation of this flow, but it does give a rough idea of its expected X-ray luminosity.

Let the main contribution to the residual emission come from gas shocked during an average time \( t_r \) before present, where \( t_r \) is a substantial fraction of the orbital period. Let a fraction \( q_m \) of the mass blown by the secondary during this time, \( q_m \dot{M}_2 t_r \), be enclosed in a fraction \( q_f \) of the volume \( 4\pi r(t_r) \). Because of the expansion, this gas is much cooler than the post-shock temperature of \( \sim 10^8 \text{ K} \). Say \( T(t_r) \sim 10^7 \text{ K} \), for which the contribution of the cooling gas to the emission in the 2–10 keV band is only \( F_{210} \sim 0.06 \) (eq. [5]). Because of the expansion to large distances, this emission is hardly absorbed in comparison to the case for the accretion X-rays, although some absorption still exists. Taking gas at \( \sim 10^7 \text{ K} \), we find the X-ray emission to be

\[
L_{\text{X-res}} \sim \frac{3 \times 10^{32} q_m^2 F_{210}^2 t_r}{q_f^2 0.05 (t_r)^{-1} (1 \text{ yr})^2} \left( \frac{\dot{M}_2}{10^{-5} M_\odot \, \text{yr}^{-1}} \right)^2 \left( \frac{v_1}{500 \text{ km s}^{-1}} \right)^{-3} \text{ ergs s}^{-1}.
\]

(22)

This suggests that old shocked secondary wind might in principle account for a significant fraction of the soft \( \sim 3 \times 10^{33} \) ergs s\(^{-1}\) emission during the X-ray minimum. In a future paper the spectrum of this radiation will be compared with observation.

4.3.2. Collimated Outflow during the Accretion Phase

We consider the specific angular momentum of the accreted matter \( j_a \) and compare it to \( j_f = (GM_2 R_2)^{1/2} \), the specific angular momentum of a particle in a Keplerian orbit at the equator of the accreting star of radius \( R_2 \). When a compact secondary star moves in a circular orbit and accretes from the wind of a mass-losing star, such that the accretion flow reaches a steady state, the ratio of the specific angular momenta is (Soker 2001)

\[
1 < \frac{j_a}{j_f} \sim 0.1 \frac{\left( \frac{M_1 + M_2}{150 M_\odot} \right)^{1/2} \left( \frac{M_2}{30 M_\odot} \right)^{3/2} \left( \frac{R_2}{20 R_\odot} \right)^{-1/2}}{\left( \frac{2 \text{ AU}}{r} \right)^{3/2} \left( \frac{v_{\text{wind1}}}{400 \text{ km s}^{-1}} \right)^4},
\]

(23)

where \( \eta \) is the ratio of the accreted angular momentum to that entering the Bondi-Hoyle accretion cylinder. There is a net accreted angular momentum because more mass is accreted by the secondary from the denser region facing the primary than from the other side, behind the secondary. When a steady state accretion flow is reached, the accretion column (and the accretion
line) bends toward the lower density region, partially compensating for its lower density and reducing the net accreted angular momentum. For this case, numerical simulations show that $\eta \sim 0.2$ (e.g., Ruffert 1999). However, here no steady state is reached around phase $\sim 0.004$, when the accretion radius increases by a large factor in a short time (Fig. 2). This increases the mass-accretion rate in a short time (Fig. 8), and therefore the accretion column has no time to bend and reduce the specific angular momentum of the accreted gas. This is the reason for taking $\eta = 1$ here.

As seen, at phase $\sim 0.004$, which occurs $\sim 8$ days after periastron, the specific angular momentum is quite large, although still less than that required to form an accretion disk ($j_{\text{ac}}/j_{\text{2}} > 1$). This implies that the accreted mass will be concentrated near the equatorial plane, with lower density regions along the polar directions. Consequently, the strong wind blown by the secondary (when undisturbed) is not efficiently suppressed along the polar directions and, together with a disk wind from the accreted matter, leads to the formation of jets or a biconical collimated wind. This wind is then shocked by the impact of the ambient primary wind and emits the low-level X-rays observed in the middle of the minimum. This emission is highly obscured by the high column density of the primary wind, and hence it contributes mostly in the hard band. We show that the X-ray emission during minimum can be fitted by reasonable parameters. Let us take the polar outflow to be at a speed of $v_p = 2000$ km s$^{-1}$, with the mass outflow rate being a fraction of 0.05 of the accreted mass. Taking the accretion rate at phase 0.05 (10 days after periastron passage) from our calculations (Fig. 8, top) to be $\sim 10^{-6} M_\odot$ yr$^{-1}$, the total kinetic energy in this polar outflow is $L_{\text{p-kin}} = 6 \times 10^{34}$ ergs$^{-1}$. For a column density of $10^{24}$ cm$^{-2}$, all the soft X-ray is absorbed, while the hard X-rays, comprising $\sim 0.4$ of the flux, are reduced by a factor of $\sim 5$, giving a luminosity of $L_{\text{p-x}} \sim 5 \times 10^{33}$ ergs$^{-1}$. This hints that a short-duration collimated outflow might account for the increase in the middle of the minimum. On even shorter timescales and smaller amplitudes, stochastic variations in the accretion process—both the mass-accretion rate and angular momentum accretion rate—which are well documented in numerical simulations (Ruffert 1996), might lead to the small spikes observed during the minimum.

5. DISCUSSION AND SUMMARY

As was shown by Soker (2005b), the effect of the gravity of the secondary star on the primary wind becomes significant as the system approaches periastron. Cold and dense blobs are likely to form in the postshock primary-wind region, which becomes unstable. It seems plausible that close to periastron passage, when the secondary’s gravity becomes significant, these dense blobs will be accreted by the secondary. Adopting the hypothesis of Soker (2005b), we have assumed that these segments of the primary wind, which are accreted, vigorously disrupt the acceleration zone of the secondary wind, so that the secondary wind ceases to exist. The formation of cold blobs will be verified in a future paper via three-dimensional gasdynamical numerical simulations. The effect of binary accretion on the launch of stellar winds will have to be studied theoretically in the future by examining the sensitivity of the wind acceleration zone in O stars to accreted cold gas.

In this paper the accretion of dense blobs is assumed to start $\sim 40$–50 days before periastron passage (orbital phase $\sim 0.025$ to $\sim 0.02$), and the Bondi-Hoyle-type accretion flow (Fig. 7) is assumed to start $\sim 20$ days before periastron passage (orbital phase $\sim 0.01$); several days later the (almost flat) X-ray minimum starts. As can be seen from the top panel of Figure 8, $\sim 60$ days after periastron passage (orbital phase $\sim +0.03$) the accretion rate diminishes and it is assumed that the secondary star builds back its acceleration zone and its wind reappears. This is when the system starts to get out of the X-ray minimum, a process lasting $\sim 20$ days.

An interesting feature of the proposed model is the asymmetry around periastron in the relevant properties. The asymmetry results from the asymmetry of the relative velocity of the primary and secondary winds at the stagnation point. As the two stars approach each other, the relative velocity is higher than when they recede. This can be seen in the lower row of Figure 2. This effect causes the stagnation distance $D_{\text{g2}}$, the accretion radius $R_{\text{sec2}}$, the postshock primary wind temperature, and other quantities to acquire asymmetric values about periastron passage.

In the proposed model the steep decline in the X-ray emission is due to the collapse of the stagnation-point region onto the secondary, very close to periastron passage (Paper I). The collapse starts a few weeks before periastron passage, when the secondary’s gravity at the stagnation point becomes significant. This occurs when the ratio of the accretion radius to the stagnation distance $R_{\text{sec2}}/D_{\text{g2}}$ increases to about a few times 0.1. For the assumed wind parameters in the $\eta$ Car system, when this occurs, the ratio of the outflow time of the shocked primary wind, $\tau_{\text{1}}$, to the free-fall time $\tau_{\text{f2}}$, becomes roughly a few times 0.1 as well, further supporting the importance of accretion. At present, we cannot say what exact value of these ratios is needed for accretion to start. This would require three-dimensional hydrodynamical numerical simulations. Furthermore, there are large uncertainties in the binary and wind parameters of $\eta$ Car. Phenomenologically, what we can say is that the accretion model can account for the behavior around the X-ray minimum of $\eta$ Car if accretion starts when $R_{\text{sec2}}/D_{\text{g2}} \sim 0.2$ (Paper I). This occurs when the stagnation point distance from the secondary is $D_{\text{g2}} \approx 0.6$ AU. We therefore assume that accretion starts at that phase. The collapse occurs over a time period of a few times the free-fall time from the stagnation point to the secondary

$$\tau_{\text{f2}} = 5.5 \left( \frac{M_*}{30 M_\odot} \right)^{-1/2} \left( \frac{D_{\text{g2}}}{0.6 \text{ AU}} \right)^{3/2} \text{ day.}$$

The details of this process require three-dimensional numerical simulations, which are beyond the scope of this paper. We therefore make do with a simple phenomenological approach. The start of the minimum is assumed to be the time when $\tau_{\text{1}}/\tau_{\text{f2}}$ and $R_{\text{sec2}}/D_{\text{g2}}$ reach some critical value. From the X-ray light curve we assume the minimum starts when $\tau_{\text{1}}/\tau_{\text{f2}}$ and $R_{\text{sec2}}/D_{\text{g2}}$ reach $\sim 0.25$ and $\sim 0.2$, respectively. We assume that the minimum ends when these ratios fall back to these same values. Values of $\tau_{\text{1}}/\tau_{\text{f2}} \approx 0.25$ and $R_{\text{sec2}}/D_{\text{g2}} \approx 0.2$ yields a duration of $\sim 70$ days for the minimum (Paper I), in good agreement with observations. We assume a simple linear approximation for the decline to minimum and subsequent recovery. We take the significant contributions to the total X-ray luminosity, $L_{210}$ in our simple model, to be

$$L_{210}(p) = \begin{cases} L_{X1} + L_{X2}, & p < -0.015 \text{ or } p > 0.04 \\ (0.005 - p) & \end{cases}$$

$$L_{X-res} = \begin{cases} (L_{X1} + L_{X2})/0.01, & -0.015 \leq p < -0.005 \\ 0, & -0.005 \leq p < 0.03 \\ (p - 0.03) & \end{cases}$$

$$\times (L_{X1} + L_{X2})/0.01, \quad 0.03 \leq p \leq 0.04,$$
where the phase \( p \) is defined in the range \(-0.5 \leq p < 0.5\). We have included \( L_{X1} \) and \( L_{X2} \) from equations (6) and (8), but we ignore for our purposes here the small contributions of accretion and the residual emission \( L_{X\text{res}} \). A value of \( k_2 = 2 \) is assumed in eq. (7). The absorbing column density is according to eq. (10), with \( k_9 = 0.5 \) (thick line) and \( k_9 = 1 \) (thin line). Bottom row: Model absorbed X-ray luminosity with \( k_9 = 0.5 \) (thick line) and the observed X-ray luminosity from Corcoran (2005) in the 2–10 keV band for a period covering two X-ray minima (two thin lines). The dashed line in the bottom right panel shows the X-ray emission for \( \omega = 180^\circ \) and \( i = 42^\circ \) (see eq. [11]).

**Fig. 9**. Top row: Absorbed X-ray luminosity between 2–10 keV according to our model, namely, eq. (25), but with absorption included and not considering the residual emission \( L_{X\text{res}} \). A value of \( k_2 = 2 \) is assumed in eq. (7). The absorbing column density is according to eq. (10), with \( k_9 = 0.5 \) (thick line) and \( k_9 = 1 \) (thin line). Bottom row: Model absorbed X-ray luminosity with \( k_9 = 0.5 \) (thick line) and the observed X-ray luminosity from Corcoran (2005) in the 2–10 keV band for a period covering two X-ray minima (two thin lines). The dashed line in the bottom right panel shows the X-ray emission for \( \omega = 180^\circ \) and \( i = 42^\circ \) (see eq. [11]).
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