ON ACCRETION COMPONENT OF THE FLARE ACTIVITY IN ALGOL

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Abstract. A critical assessment of the observational data on flare activity in Algol-type binaries is given. Two sites of the flare activity have been identified in Algol through observations in the X-ray and microwave regions. One of them is apparently the accretion shock region on the B8-type primary component. Another, the predominant one, is associated with coronal mass ejection from the Roche lobe filling late-type secondary. We still do not know the morphology of the circumbinary gas with sufficient precision to determine the relative contribution from both these sources of activity. Future observations should concentrate both on systems like U Cep with mass-loss $M \simeq 10^{-6} M_\odot/y$ and on Z Vul-type objects where secondaries are earlier than F2-type stars, i.e. presumably without an extensive convection zone. A simple model of radially expanding stellar wind in a binary system is presented. The effect of anisotropy which appears due to displacement of a sonic point, caused by the gravitational field of the companion star, is investigated. We estimate the X-ray flux for Algol, caused by the accretion to the primary component and find satisfactory agreement with the X-ray Ginga data for a mass transfer rate $M = (0.4 - 2.0) \cdot 10^{-11} M_\odot/y$.

Key words: stars: close binaries, eclipsing binaries, activity, mass-loss, circumstellar matter

1. Observational data on flare activity in Algol-type binaries

Systematic astrophysical research of Algol-type binaries has been underway for nearly a century. However, reliable data suggesting the occurrence of flares in these objects have started to accumulate in
the last 10–15 years due to systematic radio and X-ray observations. Unfortunately, the bulk of the data is still limited to the prototype, Algol itself. The earliest results concerning flare activity in the optical region have been sketchy and sparse. Thus, Svetchnikov (1986) quotes observations of RZ Cas by Kuznetsova et al. (1982), who found evidence for an abrupt appearance in 1977 of intensive emission lines of CII, OIII, NII and some other ions along with a considerable ultraviolet excess suggesting the sporadic powerful mass ejections from the system. Similarly, Hall (1989) reports evidence for the chromospheric activity (Ca II emission lines) found in U Cep by Baldwin (1973) and in S Vel by Bond (1972), in both cases during the phase of total eclipse. Only quite recently, ultraviolet spectra of U Cep, taken during the total eclipse phase, revealed the presence of emission lines typical for the chromospherically active RS CVn-type stars (McCluskey 1993).

Richards (1990), who has analyzed numerous light curves of Algol at 1.2 \( \mu\)m, finds evidence for the enhanced activity of the secondary component, a cool subgiant (variations in depths of both minima amounting to 0.1 mag, a peculiar shape of the secondary minimum, etc.). However, the results of this analysis should be taken with some caution: the author has earlier indicated (Pustylnik 1990) that the recombination emission from the circumbinary gas with characteristic masses of \( M_{\text{env}} \simeq 10^{-9} \text{--} 10^{-7} M_\odot \) should produce a similar effect. The contribution of the envelope to the total luminosity of a binary increases with the wavelength and it simulates the effects similar to the reflection and tidal distortion effects, whereas the corresponding color excesses in the visible and the near infrared regions are very modest amounting only to several hundredths of mag. Although just quoted estimates of \( M_{\text{env}} \) are much higher than the mass of the accretion bulge given by Richards (1992), they are in a good accord with the measurements of the linear polarization made by Koch et al. (1989) for a sample of 35 Algol-type binaries.

The Algol-type binaries are weak X-ray and radio sources (Elias and Mutel 1993, Lestrade et al. 1988). X-ray studies of several nearby Algol-type systems indicate that they have quiescent X-ray luminosities in the range of \( L_x \simeq 10^{30} \text{--} 10^{31} \text{erg s}^{-1} \) (White and Marshall 1983). These authors have found that \( L_x \) luminosity is in excellent agreement with the relation \( L_x = 10^{27} \cdot v_{\text{rot}}^2 \) obtained by Pallavicini et al. (1981) for single late-type stars, \( v_{\text{rot}} \) being the equatorial velocity for the synchronized rotation. As in the case of
X-ray studies, only a small number of the Algol-type binaries have been investigated in the radio wavelength range with the VLA at 3 cm and 6 cm. The radio properties of Algol-type systems are very similar to those of chromospherically active RS CVn-type stars. The median radio luminosity at 6 cm found for the Algol systems, nearly $L \simeq 10^{16}$ erg s$^{-1}$Hz$^{-1}$, is very close to the value found for a large sample of the RS CVn-type objects. Very high brightness temperatures of the order $T_B \simeq 10^8 - 10^9$ K suggest a non-thermal origin of the radio emission.

Although both X-ray and radiofluxes are lower by many orders of magnitude than the optical luminosities of the Algol-type stars, these spectral windows are more advantageous for studying the flare activity, because flare fluxes are comparable to the flux in a quiescent state and because they probe physical conditions in the circumstellar and circumbinary gas. Multi-wavelength observations are badly needed for constructing reliable models and separating the flare radiation component from the radiation in the quiescent phase. Up to now this has been done only for Algol itself, because of its proximity to us.

Since the radio quiescent flux is by at least ten times lower than the flare component, it enables one to determine in principle the morphology of the ambient gas (see, for instance, Woodworth and Hughes 1976) or at least to estimate the upper limit of the mass-loss rate. These estimates set up upper limits of $\dot{M} \lesssim 10^{-8} M_\odot$ per year for typical Algol-type binaries.

An alternative way to estimate $\dot{M}$ comes from the analysis of light curves in the visual region. Mass transfer from the secondary component should result in small variations of the depth of a primary minimum. It has been found that $\dot{M}$ is related to the optical depth of gas $\tau$ through quite a simple relation $\dot{M} = \tau \cdot f(i, r_1, r_2)$, where $f$ is an elementary function depending on orbital elements of a binary, which are known with sufficient accuracy (Pustylnik and Einasto 1990). The value of $\tau$ can be estimated from the fluctuations of the depths of minima on the light curve. Both estimates of $\dot{M}$ are in a satisfactory agreement.

On the other hand, the accurate measurement of $\dot{M}$ is a prerequisite for reliable determination of the accretion luminosity. However, the fitting of X-ray quiescent spectra by two-temperature thermal bremsstrahlung models yield the temperatures of $(25 - 35) \cdot 10^6$ K
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which are higher by nearly an order of magnitude than the conservative estimates of the accretion temperatures (Stern et al. 1992).

Because of a possible contamination from the Perseus cluster in the Ginga data, it is still difficult to estimate the contribution from the accretion shock region into the total X-ray luminosity. The derived emission level is such that nearly one per cent of the total mass loss would be sufficient to ensure the observed X-ray luminosity. According to the IUE ultraviolet data (Peters and Polidan 1984), a moderately hot \(10^5\) K and low density \(n_e \approx 10^9\) cm\(^{-3}\) plasma has been discovered, apparently surrounding the hot component. Typical dimensions of the high turbulent accretion region (HTAR) identified by Peters and Polidan in AU Mon, U CrB, CX Dra and TX UMa are in a reasonably good agreement with the scale height for the X-ray source \(H_x \approx 1.2 \cdot 10^{11} \cdot T_7^2\) cm (Harnden et al. 1977), where \(T_7\) is the temperature in units of \(10^7\) K. On the other hand, the decay time of the HTAR, comparable to the orbital period, agrees more favorably with the idea of coronal mass ejections (CME) following from the Ginga observations of the X-ray flares on Algol (Stern et al. 1992).

One possible solution of some controversy concerning the contribution from the HTAR would be to concentrate in future on the X-ray observations of such systems like U Her and Z Vul, where the secondary component is possibly earlier than F0 (thus, the latter is not expected to possess an extensive convection zone) and, on the other hand, like U Cep, where during the periods of high activity mass-loss is \(\dot{M} \approx 10^{-6} M_\odot\) per year. Thus, the accretion luminosity must be higher by at least two orders of magnitude than in Algol itself.

A large flare, with \(L_x \approx 10^{31}\) erg s\(^{-1}\), lasting for over 12 hours, has been recorded by Ginga observations in January 1989 (Stern et al. 1992) with the temperature in the flare peak exceeding the temperature in the quiescent state by a factor of two. Similar results have been obtained earlier in the EXOSAT observations (White et al. 1986). The presence of the Ca K line in both cases confirms the thermal character of the emission.

The Ginga observations set up the upper limits of about 20 % and 10 %, respectively, for the eclipsed flux fraction during the primary and the secondary eclipses. This fact, along with the high flare luminosities (in comparison with the quiescent state), plus the long duration of the flares suggest that the flare activity is associated with the extensive emitting volumes comparable to the size of the orbit and with a complicated morphology of the emitting region.
One can recall in this connection the earlier paper by Florkowski (1980) who, interpreting the radioflares in Algol, proposed a picture in which radio flares have been generated due to interaction effects of random mass ejections with the gas cloud surrounding the whole binary. At present, the analysis of X-ray flares is based upon a solar-type CME with evolved loop-like structures, assuming different scenarios of radiative and (or) conductive cooling and involving the respective time-scales (see Stern et al. 1992 and references therein). This analysis suggests loop heights comparable with the radius of the secondary K-type component (between $5 \cdot 10^{10}$ cm and $2 \cdot 10^{11}$ cm) and $n_e \simeq 6 \cdot 10^{10}$ cm$^{-3}$. Despite the high degree of sophistication of current models of CME, they still do not take into account the binary nature of the flaring object. In what follows, we shall consider a crude model of the radially expanding wind in a detached binary system, in particular, we shall estimate explicitly the angular dependence of the mass flux and the results will be applied for evaluation of the accretion luminosity and the respective X-ray flux in the Algol system.

2. Anisotropic stellar wind in a binary system

Next we shall consider a somewhat simplified model of anisotropic stellar wind in a moderately close binary system where stellar wind from one component overwhelms the wind from the companion star. Our purpose is to estimate the effects of anisotropy due to the displacement of a sonic point caused by the gravitational field of the companion star.

For a stationary flow ignoring the Coriolis forces due to stellar rotation, the equation of motion of gas and the equation of continuity are

\[(\vec{u} \nabla) \vec{u} + \nabla \Phi + \frac{1}{\rho} \nabla P = 0,\]  
\[\nabla (\rho \vec{u}) = 0.\]  

Here $\vec{u}$ is the velocity of gas, $\rho$ and $P$ are, respectively, the local density of matter and the pressure exerted by both gas and radiation, $\Phi$ is the gravitational potential. Equation (1) is a good approximation, provided that $u_s/u_{orb} \ll 1$, where $u_s$ is the local sound velocity and $u_{orb}$ is the velocity of the orbital motion. We shall treat here an
idealized case of the radial expansion of gas from one component of a close binary.

In this case, Equation (2) reduces to the conservation of mass flux (mass flow rate per unit solid angle) along the stream-line, i.e.

\[ J = \rho u R^2 = \text{const} \]  

(3)

and Equation (1) reduces to

\[ u \frac{du}{dR} \quad \text{and} \quad \frac{dP}{dR} = \frac{1}{\rho} \frac{dP}{dR} \quad \text{and} \quad \frac{d\Phi}{dR}. \]  

(4)

Introducing the gravitational potential of the Roche model, using Equation (3) and the relation between the gas pressure and the density of matter, valid for ideal gas \( P = \frac{1}{2} u^2 n \), one obtains

\[ \frac{1}{2} \left( 1 - \frac{u_s^2}{u^2} \right) \frac{du^2}{dR} = \frac{2}{R} u_s^2 - \frac{du_s^2}{dR} - \frac{d\Phi}{dR}, \]  

(5)

where \( u_s \) is the local sound velocity.

At the sonic point the following set of equations should be fulfilled

\[ R = R_s, \quad u = u_s, \]  

(6a)

\[ \frac{2}{R_s} u_s^2 - \frac{2k}{m} \frac{dT}{dR} \left. \right|_s - \frac{d\Phi}{dR} \left. \right|_s = 0. \]  

(6b)

Introducing, for the sake of simplicity, the gravitational potential for a two-dimensional case (valid for the equatorial plane)

\[ \Phi = \frac{GM_1}{R} + \frac{GM_2}{\sqrt{a^2 - 2aR\mu + R^2}} + \frac{G(M_1 + M_2)}{2a^3} \left[ R^2 - \frac{2M_2}{M_1 + M_2} a R\mu + \left( \frac{M_2}{M_1 + M_2} \right)^2 a^2 \right], \]  

(7)

we find consequently from (6b)
\[
\frac{dT}{dR_s} \bigg|_s = \frac{2T_s}{R_s} - \frac{m}{2k} \left[ \frac{GM_1(1 - \Gamma_s)}{R_s^2} - \frac{GM_2(a\mu - R_s)}{(a^2 - 2a\mu R_s + R_s)^{3/2}} \right. \\
\left. - \frac{G(M_1 + M_2)}{a^3} \left( R_s - \frac{M_2a\mu}{M_1 + M_2} \right) \right]. \tag{8}
\]

Here \( G \) is the gravitation constant, \( M_1 \) and \( M_2 \) are masses of the components, \( \mu \) is the cosine of the angle between the line joining the centers of two stars and the direction considered in the equatorial plane from the center of a mass-losing component, \( k \) is the Boltzmann constant, \( m \) is the mass of the hydrogen atom, \( T_s \) is the temperature at the sonic point and the factor \( \Gamma_s \) takes into account the radiation pressure. To facilitate derivation of the subsequent formulae, we express the radial velocity in dimensionless units \( v = u/u_e \), where \( u_e = \sqrt{2G(M_1 + M_2)/a} \). Also, we introduce a dimensionless potential \( C = 2\Phi a/G(M_1 + M_2) \) (Plavec and Kratochvil 1964) and express \( r \) in the units of \( a \). In Equation (8) \( \mu \) is treated as a parameter. Actually the absence of another equation involving \( du/d\mu \), in addition to Equation (5), implies that we assume simply \( u_\mu \ll u_r \), contrary to the case of the accretion disc theory, where usually it is supposed just the opposite, i.e. \( u_\mu \gg u_r \).

Turning back now to the condition of conservation of mass flux \( J \), we should expect (Hadrava 1987)

\[
J \sim \exp \frac{-\Phi_s}{kT_s}. \tag{9}
\]

As we see from Equation (8), \( T_s \) must be direction dependent, or in other words, Equation (9) quantitatively describes the anisotropy which appears due to the influence of the companion star. We may estimate the rate of anisotropy by assuming that

\[
T_s = T_{s0} + \left( \frac{dT}{dr} \right)_s \Delta r_s, \tag{10}
\]

where the displacement of a sonic point \( \Delta r_s \) can be found with the aid of the subsequent formulae (see also Pustylnik 1994):

\[
\Delta r_s \simeq \frac{q}{m} \frac{u_e^2}{u_s^2} f(q, r_s), \tag{11}
\]
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\[ u_e^2 = 2G(M_1 + M_2)/a , \quad q = M_2/M_1 \]

and

\[ f(q, r_s) = 2(1 + q)^{-1}(1 - r_s)^{-2}(1 + r_s)^{-2} - 1. \]

Relation (11) can be readily found from Equation (8) if we assume that \( T = T_s(r_s/r)^m \) for \( r_s \leq r \leq r_s + \Delta r_s \) and set \( \Delta r_s = r_{s1} - r_{s2} \), where \( r_{s1} \) and \( r_{s2} \) are the coordinates of the sonic points for \( \mu = 1 \) and \( \mu = -1 \), respectively. In fact \( \Delta r_s \) is nothing else than the scale of the thermal drop-off. By \( T_{s0} \) in Equation (10) we denote an average temperature which is fixed by simply postulating that the kinetic energy of a gas particle plus the specific enthalpy should be equal to the difference between the potentials at the surface of the mass losing component and the Roche critical potential, thus,

\[ T_{s0} = T_{s0}^0 \left[ \frac{(1 - \Gamma_s)}{r_s^0} + \frac{q}{1 - r_s^0} + \frac{1}{2}(1 + q) \left[ \frac{r_s^0}{1 + q} - C_1(q) \right] \right], \]

where

\[ T_{s0}^0 = 7.64 \cdot 10^6 K, \quad C_1(q) = \frac{2\Phi a}{G(M_1 + M_2)}. \]

Combining Equation (3) and Equations (8)–(12), we have the condition of the mass flux conservation in the form

\[ J = \rho v r^2 = B J_0 F(v_e/v_s, q, \mu, r_s^0), \]

where \( J_0 = \rho_0 v_0 r_s^2 \) and the total mass-loss is

\[ \dot{M} = 2\pi \int_{-1}^{+1} J d\mu = 4\pi J_0. \]

The function \( F \) is given by

\[ F = \exp \left( -\frac{\Phi_s}{kT_{s0}} \left[ 1 + \frac{d\ln T}{dr}|_{s0} \Delta r_s^0 \right]^{-1} \right), \]

and it can be easily calculated by formulae (8) and (10)–(12). Normalization constant \( B \) is the following:

\[ B = \frac{1}{2\pi} \frac{1}{\int_{-1}^{+1} F d\mu}. \]
Fig. 1 demonstrates the angular dependence of the mass flux $J(\mu)$ for several values of the mass ratio $q$. Note that the rate of anisotropy increases with increasing $q$ and that maximum of $J(\mu)$ falls on the value of $\mu = 0.5$, due to the influence of the term responsible for the centrifugal force. Since our calculations are based on the assumption of a radially expanding wind, the position of maximum of the $J(\mu)$ distribution is relevant only to our specific model. The rate of anisotropy is sensitive to the ratio $v_s/v_e$ and the scale height of the thermal drop-off (or to the displacement $\Delta r_s$ of a sonic point). We mention here in passing that different ad hoc flux-tube approximations have been introduced in a number of papers, among others by Modisette and Kondo (1980), Haisch, Linsky and Basri (1980), Kopp and Holzer (1976). From all these considerations it is evident that Equation (14) may be regarded as an equation of a flux-tube based on very simple but physically justified assumptions.

Let us now consider the case of a plane-parallel thin conduction zone. The equation of heat conduction with the energy loss by radiation is

$$\frac{d}{dr} \left( K_0 T^{5/2} \frac{dT}{dr} \right) = j_0 n_e^2 ,$$

where $j_0$ is the emission coefficient and $K_0$ is rather insensitive to the density and the temperature (Hearn 1975). Taking this into account and neglecting the change of the gravitational field within the conduction zone, we may insert Equation (8) for the temperature gradient into Equation (16) and estimate the angular dependence of the emission measure $j_0 n_e^2$. The results are shown in Fig. 2 for some values of $r_s$ and $q$. Despite the approximate nature of our treatment (a radially expanding wind is not a satisfactory approximation for a close binary), it is noteworthy to point out that the effect of anisotropy is most pronounced just for the case, when the sonic point is located close to the first Lagrangian point and the mass of the perturbing component is comparable to or larger than the mass of the companion losing the matter through the wind. Both these conditions are met for the mass losing secondary components of Algol-type binaries. Thus, the anisotropy of the wind may influence the flares. It may affect the standard picture of the shock region, i.e. the depth of the region where the ram pressure becomes equal to the pressure of gas in the atmosphere. On the other hand, the phase variations of $dT/dr$, stemming up from the binary nature of the source, should be taken into account in the future analysis.
Fig. 1. Angular dependence of the mass flux $J(\mu)$. The direction towards the mass accreting component of the binary corresponds to $\mu = 1$. Identification of the curves: 1 is for $q = 1.0$ and $r_s = 0.35$, 2 is for $q = 0.5$ and $r_s = 0.30$, 3 is for $q = 0.3$ and $r_s = 0.26$. $r_s$ is the coordinate of the sonic point.

Fig. 2. An example of the angular dependence of the emission measure for several values of the mass ratio $q$. The direction towards the mass accreting component of the binary corresponds to $\mu = 1$. For identification of the curves see Fig. 1.
3. Evaluation of the mass transfer rate in Algol

Now we may use Equation (8) and Equations (11) – (15) of the preceding section together with the known observed parameters of Algol to estimate the mass transfer rate due to accretion on the B8 primary component. The thermal (free-free) radioflux for a single star losing matter through the isotropic stellar wind is given by a subsequent expression (Panagia and Felli 1975)

\[
S_\nu = 5.12 \left( \frac{\nu}{10 \text{ GHz}} \right)^{0.6} \left( \frac{T_e}{10^4 \text{ K}} \right)^{0.1} \left( \frac{\dot{M}}{10^{-5} \text{ } M_\odot/\text{y}} \right)^{4/3} \\
\times \left( \frac{\bar{m}}{1.2} \right)^{-4/3} \left( \frac{u_{\text{esc}}}{10^3 \text{ km/s}} \right)^{-4/3} \bar{Z}^{2/3} \left( \frac{d}{\text{kpc}} \right)^{-2} \text{ mJy.}
\]

(17)

Here \(\bar{m}\) is the mean molecular weight, \(\bar{Z}\) is the average ionic charge, and the remaining notations are self-explanatory. Using the recent very precise determination of Algol’s parallax, 0."0343 ± 0."00085 (Gatewood, Jonge and Heintz 1995), assuming \(S_\nu\) equal to radioflux of the non-flaring component, \(S_\nu = 50 \text{ mJy} (\nu = 2695 \text{ MHz})\) (Woodworth and Hughes 1976), and taking the escape velocity equal to \(u_e = \sqrt{2GM_1/M}\) which gives \(u_e = 350 \text{ km/s}\) for the Algol system, we obtain \(\dot{M} = 2 \cdot 10^{-7} \text{ } M_\odot \bar{m} \sqrt{\bar{Z}}/\text{y}\). This is certainly the upper limit estimate, because it is now generally believed that the predominant source of radiation, measured in the cm wavelength range, is of synchrotron nature. Richards et al. (1990) and Richards (1992) find that \(10^{-11} \text{ } M_\odot \leq \dot{M} \leq 10^{-10} \text{ } M_\odot\) per year. This estimate follows from the mass of a localized HII region identified by spectroscopic observations by Richards, and also from the considerations which presume equipartition between the magnetic field strength and the kinetic energy of the gas stream at the inner Lagrangian point (Bolton 1989). Another independent estimate of \(\dot{M}\) comes from the analysis of the optical light curves of Algol. The following relation between \(\dot{M}\) and the radial optical depth, for the Thompson scattering in the radially expanding stellar wind, has been derived by Pustylnik and Einasto (1990):

\[
\dot{M} = 10 \pi \mu u_s r_2 f(a, i, r_2, R_{\text{accr}})^{-1},
\]

(18)

where
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\[ f(a, i, r_2, R_{accr}) = \arctan \sqrt{\frac{R_{accr}^2 - a^2 \cos^2 i}{a \cos i}} + \arctan \frac{a \sin i - r_2}{a \cos i}. \]

Here \( \dot{M} \) is the mass loss in the units of \( 10^{18} \) g/s, \( u_s \) is the velocity at the sonic point in the units of 100 km/s, \( r_2 \) is the radial optical depth at the centre of the secondary minimum, \( R_{accr} \) is the accretion radius (according to the spectroscopic data of Richards (1992), \( R_{accr} = (1.8 \pm 0.2) \cdot r_1 \), \( r_1 \) being the radius of the primary, mass accreting component. The value of \( r_2 \) can be estimated using the data on fluctuations near the bottom of the secondary minimum of Algol (see, for instance, the compilation of data by Sjöderhjelm 1980). In the case of Algol, the bottom of the secondary minimum gives us the best chance to estimate the optical depth of gas along the line of sight, because the inclination angle is \( i \simeq 89^\circ \), and in the secondary minimum we observe a partial eclipse. We assume that \( r_2 \leq (1 - 5) \cdot 10^{-3} \), because small fluctuations (up to 0.01–0.02 mag) have been observed at the respective phases. Thus, from Equation (18) we find \( 10^{-10} M_\odot \leq \dot{M} \leq 5 \cdot 10^{-10} M_\odot \). Again, this is only an upper limit estimate.

The third and actually the most straightforward way to estimate the mass-loss rate \( \dot{M} \) is to analyze observations of the period changes which for Algol have been recorded over 300 years. Paradoxically, this classical approach is the least reliable method for evaluation of \( \dot{M} \). If we assume a simple mass-loss (no mass transfer, no angular momentum loss), then for \( \Delta P/P \simeq 10^{-9} \) (Sjöderhjelm 1980) this results in \( \dot{M} \simeq 2 \cdot 10^{-8} M_\odot/y \). However, this seems to be a gross overestimation, because it is generally believed that, due to the presence of magnetic field (\( B \simeq 100 \) G) the period variations connected with the Roche-lobe-filling secondary component, are caused basically by the angular momentum loss from the orbital motion (Tout and Hall 1991, Yungelson, Tutukov and Fedorova 1989, Bolton 1989).

Summing up the above-given arguments, we take \( \dot{M} \simeq (1 - 5) \cdot 10^{-10} M_\odot/y \) as the most reliable upper limit estimate of the mass-loss rate for Algol. Then, integrating Equation (13) for the mass flux \( J(\mu) \) over \( \mu \), we find (using the data of Richards (1992) for \( R_{accr} \)) the following estimate for the accretion rate: \( \dot{M}_{accr} \simeq (0.4 - 2.00) \cdot 10^{-11} M_\odot/y \). By approximating the average temperature in the accretion region by Equation (12), we find for Algol \( T_{accr} \simeq 2.8 \cdot 10^6 \) K. According to the calculations of Lubow and Shu (1975) (see also Ulrich and Burger (1976)), the effective cross-section of the
gas stream can be approximated as

$$S_{str} = 6.9 \cdot 10^{-5} \frac{T_4 r_2^3}{\bar{m} M_2} ,$$

(19)

where $S_{str}$ is expressed in solar units, $T_4$ is the temperature in the units of $10^4 K$, $\bar{m}$ is again the mean molecular weight and $r_2, M_2$ are also expressed in solar units. Thus, we find from Equation (19) $S_{str} \simeq (1 - 3) \cdot 10^{20} \text{cm}^2$. Actually, this estimate practically coincides with a simple evaluation of a thermal drop-off scale of the stream width $H/R \simeq u_s^2/u_e^2$ in our notation (see, for instance, Pringle 1985). If we take this estimate close enough to the first Lagrangian point, then the value of $H/R$ is fully determined by the ratio $G(M_1 + M_2)/kT$ and by geometry of the Roche model (see also Meyer and Hoffmeister 1983). We intentionally emphasize this point since both the emission measure and the X-ray luminosity are proportional to $H^3$. Taking now our derived value of $\dot{M}_{accr}$ and taking into account that gas is fully ionized in the stream, we find for the average electron density $n_e \simeq (1 - 3) \cdot 10^{10} \text{cm}^{-3}$. Although this estimate is only of an order of accuracy, it agrees satisfactorily with the value found from the analysis of the Ginga data on X-ray flares. This last figure is translated into the emission measure $EM$ for the accretion region $\int n_e^2 dV \simeq (1 - 5) \cdot 10^{52} \text{cm}^{-3}$.

We take the emissivity function for optically thin gas for the X-ray range considered ($1.2 - 18$ keV) from Mewe, Lemen and van den Oord (1986)

$$P_c(\lambda, T) = 2.051 \cdot 10^{-22} \cdot n_e^2 \cdot G_c \lambda^{-2} \cdot T^{-1/2} \exp [-143.9/(\lambda T)] ,$$

(20)

where $P_c(\lambda, T)$ is in the units of erg·cm$^{-3}$·s$^{-1}$·Å$^{-1}$, $\lambda$ is in Å and $T$ in $10^6$ K. According to Mewe, Lemen and van den Oord (1986), the total Gaunt factor (the sum from free-free, bound-free and two photon emission contributions) is practically independent of $\lambda$ for the energy range considered. Thus, we have assumed an average value of $G_c$ taken from the graphical data of the just quoted paper. Now combining this with our estimate of the emission measure for an accretion region, we finally obtain for the X-ray luminosity $L_x \simeq (0.4 - 2) \cdot 10^{30} \text{erg s}^{-1}$ (which is nearly by an order lower than the earlier estimate of Harnden et al. 1977). Our upper limit estimate relates to a higher value of $T_{accr} = 6 \cdot 10^6$ K and it reflects an uncertainty in the impact velocity for the stream. We arrive at a
lower value for $L_x$, because our mass transfer rate is roughly lower by an order of magnitude than that used in the paper of Harnden et al. (1977).

To summarize, our results do not contradict the Stern et al. (1992) analysis of the X-ray observations of Algol by the *Ginga* satellite, which set the 10–20 % upper limit for the eclipsed flux fraction during the two eclipses. As mentioned above, the effective area of the accretion shock region must be by at least two orders smaller than the eclipsed region.

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