

COMETARY COMPACT H II REGIONS ARE STELLAR-WIND BOW SHOCKS

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ABSTRACT

Comet-shaped H II regions, like G34.3+0.2, are easily explained as bow shocks created by wind-blowing massive stars moving supersonically through molecular clouds. The required velocities of the stars through dense clumps are less than $\sim 10 \text{ km s}^{-1}$, comparable to the velocity dispersion of stars in OB associations. An analytic model of bow shocks matches the gross characteristics seen in the radio continuum and the velocity structure inferred from hydrogen recombination and molecular line observations. The champagne flow model cannot account for these structures. VLBI observations of masers associated with the shells of cometary compact H II regions should reveal tailward proper motions predominantly parallel to the shell, rather than perpendicular. We predict that over a decade baseline, high signal-to-noise VLA observations of this class of objects will show headward pattern motion in the direction of the symmetry axis, but not expansion. Finally, shock-generated and coronal infrared lines are also predicted. In a companion paper we present detailed numerical models.

Subject headings: masers — nebulae: H II regions — nebulae: internal motions — shock waves — stars: winds

I. INTRODUCTION

Comet-shaped ultracompact H II regions are a relatively new class of ionized region found around what are presumed to be hot luminous stars deeply embedded in molecular clouds. Since the discovery of the prototype, G34.3+0.2 (Reid and Ho 1985; shown here as Fig. 1, reproduced from Wood, Mac Low, and Van Buren [1991; see § Va]), an additional four were found by Gaume and Mutel (1987), and the survey by Wood and Churchwell (1989; hereafter WC) turned up another 15. In addition, we consider it likely that G70.7+1.2 (Bally *et al.* 1989 and references therein) is also a member of this class on the basis of its energetics and morphology. The statistics of WC show that somewhat greater than 20% of all ultracompact H II regions have cometary shapes.

Well-ordered ultracompact H II regions with cometary and shell-like morphologies present a theoretical problem. Dense ionized structures on these scales are expected to evolve in a sound-crossing time ($\approx 10^3 \text{ yr}$) if there is no confining mechanism. That such a mechanism must exist is apparent by comparing the number of observed ultracompact H II regions with the formation rate of massive stars in the galaxy. Far too many exist. In the remainder of this paper we argue that the confinement mechanism is ram pressure caused by the relative motion of a wind-blowing ionizing star through a molecular cloud. In doing so we apply and extend the theory developed earlier by Van Buren and McCray (1988, hereafter VM) for stellar-wind bow shocks seen around runaway O stars in the diffuse inter-

stellar medium to the higher densities and lower velocities appropriate to molecular clouds.

In § II we present a summary of the observations to date and discuss the various models proposed for these objects. We then present in § III a simple analytic model for bow shocks in molecular clouds. We compare the model with observations in § IV, where we show that our picture reasonably accounts for the data. Finally, in § V we propose a set of experiments involving the detection of infrared cooling lines and proper motions of the ionized gas and associated OH masers to confirm or reject our hypothesis.

In a companion paper, Mac Low *et al.* (1990; hereafter Paper II), we present detailed numerical models of bow shocks in molecular clouds to compare directly with radio continuum images and spectral line observations. We find that good fits are always found without resorting to extreme parameters, supporting the view presented here that cometary compact H II regions are indeed bow shocks.

II. SUMMARY OF THE OBSERVATIONS AND PROPOSED INTERPRETATIONS

The prototypical cometary compact H II region is G34.3+0.2, first pointed out as such by Reid and Ho (1985). Its morphology is that of a limb-brightened paraboloidal shell with a bright coma, or head, and a swept-back tail of much lower surface brightness. At 15 GHz this source has a flux of 4 Jy within an area $3''.4$ across. At a distance of 3.7 kpc (WC), the flux corresponds to a star earlier than O6.5 ZAMS ionizing an H II region with a linear scale of 0.06 pc. The H II region is embedded in a dense NH_3 core and has NH_3 emission associ-

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FIG. 1.—A deep 8.4 GHz VLA image of G34.3+0.2, the prototypical compact cometary H II region. Note the parabolic coma at left, onto which attaches an apparent limb-brightened cylinder. Two unresolved sources appear at the far left, which we are using as positional references to measure the proper motion of the ionization front. Details of the observations will be presented in a future paper.

ated with the shell itself (Anderson and Garay 1986). Absorption by NH_3 and HCN is seen against the continuum source (Wood, Churchwell, and Bieging 1990; Anderson and Garay 1986), with a relative velocity between the ionized and molecular gas of *approach*. Radio recombination lines of hydrogen (Garay, Rodriguez, and van Gorkom 1986) show significant broadening and a systematic velocity gradient across the source. OH masers are associated with the “upstream” edge of the shell (Garay, Reid, and Moran 1985; Gaume and Mutel 1987) and have a net velocity toward the ionized gas. Other similar objects are W49 S, G31.4+0.3, Sgr B2 I, G348.5–1.0, and NGC 6334 F (Gaume and Mutel 1987), as well as the large number found by WC.

Many physical models have been proposed to explain these objects. Below we discuss the shortcomings of the various ideas and propose instead that cometary compact H II regions are the result of the interaction of winds from moderate-velocity O stars with stationary molecular cloud material to produce stellar-wind bow shocks.

Supersonic Strömgren regions.—The Strömgren problem for an O star moving with respect to a neutral medium was solved by Raga (1986) to provide an explanation for cometary compact H II regions. No effort was made to include the dynamics so there are no shocks or flows explicitly calculated. Consequently, the model shows no limb brightening and predicts that the velocity of the neutral (molecular) and ionized gas will be the same, in contradiction to their observed relative velocities of approach. The length scale in this problem, the Strömgren radius, matches the observed sizes. The shape of the

calculated Strömgren zone resembles a teardrop, but the observed objects do not pinch back down to the symmetry axis. An extension of this model to include a shock on the upstream side of the H II region is given by Choe (1988). The model ignores stellar winds, but they are likely to be dominant in exercising structural control near O stars in molecular clouds (Shull 1980; McKee, Van Buren, and Lazareff 1984).

Distorted magnetic H II regions.—Gaume and Mutel (1987) suggest that the magnetic pressure of an H II region balances the ram pressure of an interstellar wind. This accounts nicely for the observed relative velocities of approach for the ionized and neutral gas, but it does not reproduce the limb brightening seen. The required magnetic field strength of the H II region is unreasonably high since flux freezing during the expansion of the molecular gas when it becomes ionized results in a large decrease in field strength. Finally, the inertia of the ionized gas is small with respect to a similar volume of molecular material so the magnetized gas will be stopped while the star continues onward. This would be a very transitory phenomenon, not a common one as observed. The reason Earth’s magnetosphere morphologically resembles a cometary H II region is that the field is anchored to the planet; an H II region’s magnetic field is unlikely to be so well coupled to the star.

Champagne flows.—In a series of papers, Tenorio-Tagle and coworkers have presented numerical simulations of the evolution of H II regions in media with strong density gradients (Tenorio-Tagle 1979; Bodenheimer, Tenorio-Tagle, and Yorke 1979; Tenorio-Tagle, Yorke, and Bodenheimer 1979; Bedijn and Tenorio-Tagle 1981; and Yorke, Tenorio-Tagle, and

Bodenheimer 1983). In the latter paper they present model radio continuum maps of champagne flows seen from various aspect angles and at different phases of their evolution. Although the simulations have problems such as spurious symmetry breaking, we will discuss their results at face value. The most serious disagreement between the observations and the champagne model is that real cometary compact H II regions are limb-brightened, and, that near the symmetry axis of the object, there is little emission except at the coma, i.e., their centers have low emissivity. The champagne flow simulations rarely show limb brightening because the model H II regions are isobaric, so the density (except for the flow past the nozzle) is fairly uniform. Cylindrical symmetry then requires that the model's emission measure through the middle be higher than that through the edges because of the longer path length. The observed approach of molecular gas to the ionized gas is also not explainable in the context of the champagne model: ionized gas should be streaming away from the molecular gas, toward the low-density regions. In addition, the velocities achieved should only approach the sound speed of the ionized gas in the vicinity of the nozzle, not near the coma as observed. The champagne model is an evolutionary model and predicts that more objects should be at later stages, while the most closely matching simulations correspond to young objects—where are the older ones? As an evolutionary model, it also suffers from the statistical problem discussed in the Introduction. Finally, the champagne model ignores the effects of stellar winds, which for O stars can be dramatic.

Distorted stellar wind bubbles.—A stellar wind bubble (Pikel'ner 1968; Dyson 1973; Castor, McCray, and Weaver 1975) blown in a medium with a density gradient becomes distorted when its radius is of order the density scale length (Tomisaka and Ikeuchi 1985; Mac Low and McCray 1988). Eventually its shape resembles a plume. During its evolution, the swept-up shell of material remains fairly isobaric, so motions parallel to the shell are significantly less than the sound speed. This contradicts the radio recombination line observations near the coma of G34.3+0.2, which indicate supersonic relative motions tangent to the shell (Garay, Rodriguez, and van Gorkom 1986). Distorted bubbles reproduce the limb brightening seen and account for the relative approach of molecular and ionized gas. In this picture the shell-like compact H II regions like W3(OH) are closely related objects; they just happen to be located in regions of no large-scale density gradients and form ordinary stellar wind bubbles. But again, there are too many cometary compact H II regions to be explained by this model given the birthrate of OB stars.

Moving cocoons.—In the latter stages of accretion onto an O star, after the optical photosphere has reached temperatures capable of ionizing interstellar hydrogen, radiation pressure on dust grains can balance the ram pressure of accreting material (Kahn 1974; Cochran and Ostriker 1977; Yorke and Krügel 1977). Dreher and Welch (1981) originally proposed this as a model to explain the radio continuum shell around W3(OH) to avoid the age problem. If the star is moving with respect to the gas, an elongated structure will form. Balancing the ram pressure of the oncoming interstellar gas with the radiation pressure of the stellar photon field gives a size scale

$$l_{rp} = 4.00 \times 10^{16} L_{38}^{1/2} \mu_H^{-1/2} n_{H,5}^{-1/2} v_{*,6}^{-1} \text{ cm}. \quad (1)$$

Here the stellar luminosity is $L_{38} \times 10^{38} \text{ ergs s}^{-1}$, its velocity with respect to the ambient medium of density $n_{H,5} \times 10^5 \text{ H-nuclei cm}^{-3}$ is $v_{*,6} \times 10^6 \text{ cm s}^{-1}$, and the dust is assumed to

be optically thick to stellar radiation; μ_H is the mean mass per hydrogen nucleus. The size scale is similar to that of stellar wind bow shocks (discussed below) because the momentum flux in the winds and radiation field are similar; winds of O stars are in fact driven by radiation pressure (Castor, Abbott, and Klein 1975). In the absence of a strong wind, this would be a viable model, but we know that strongly ionizing massive stars always have powerful winds.

Stellar-wind bow shocks.—The shortcomings of the above models lead us to suggest that the compact cometary H II regions are produced by stars with powerful winds moving supersonically through molecular cloud clumps, i.e., stellar-wind bow shocks. Hughes and Viner (1976) argue that the velocity structure of W3(OH) could be explained by using qualitative results from a shock model of Earth's magnetopause, and Reid and Ho (1985) suggest a bow-shock model to explain G34.3+0.2 so this idea is not new. Neither of these groups, however, provide any quantitative interpretation.

A synopsis of the bow-shock scenario, which we develop further below, is in order (Fig. 2). A UV-bright, mass-losing star moves through a molecular cloud with a velocity exceeding the local sound speed. The interaction of its wind with cloud material produces a structure which qualitatively resembles the flow around a supersonic body in the atmosphere. This "bow shock" is illuminated by stellar ionizing photons, giving rise to an H II region shaped like a thin paraboloidal shell.

One attraction of the bow-shock hypothesis is that it neatly solves the puzzle of compact H II region ages in two ways.

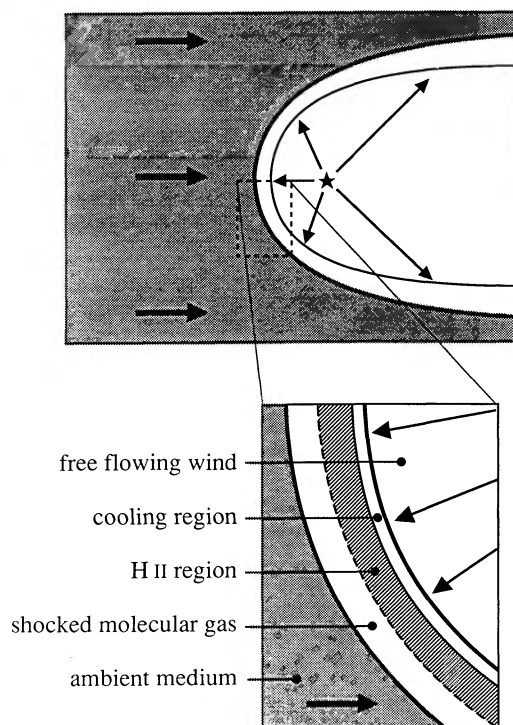


FIG. 2.—A schematic view of the stellar wind bow-shock model for compact cometary H II regions. The figure is drawn in the frame of the O star, which sees a supersonic interstellar wind from the left, distorting what would have been a stellar wind bubble into a bow shock. Moving from the inside out, the various labeled regions are separated from each other by an adiabatic shock, a conductive interface, an ionization front, and a strong radiative shock. Typical physical parameters are: density 10^5 cm^{-3} , star's velocity through cloud 10 km s^{-1} , and stand-off distance of shock $5 \times 10^{16} \text{ cm}$.

First, the H II region is confined by the ram pressure of (in its frame) oncoming molecular cloud material, so it does not expand away. Consequently the sound-crossing time of the H II region is not a measure of true age. Second, a given star may traverse a number of cloud clumps with a density of 10^5 cm^{-3} before it leaves the cloud and becomes optically visible, giving rise to several cometary H II regions over its lifetime, each of which lasts $\approx 10^4 \text{ yr}$.

The estimated size scales of stellar-wind bow shocks compare well with the observed sizes of cometary compact H II regions (below). Limb brightening occurs in the bow-shock picture since the stellar wind compresses the ionized molecular cloud material into a thin shell which subsequently flows around the stellar-wind cavity. Because the bow shock accelerates swept-up gas, at least part of which becomes photoionized, there is a net approach of the ambient molecular material toward the ionized gas. Finally, large velocity gradients are expected in the shell because gas near the head is stationary with respect to the star, while gas along the trailing edge is nearly at rest with respect to the molecular cloud. Velocity widths of order v_* are to be expected when the bow shock is observed from perpendicular to the star's direction of motion. Below, in § III we discuss the bow-shock model quantitatively in order to prepare for confronting the observations in § IV.

III. THE BOW-SHOCK MODEL

Figure 2 shows a schematic view of a stellar-wind bow shock. A star is traveling supersonically through a molecular cloud with H nucleus density $n_H = 2n_{H_2} + n_{H I} + n_{H II}$ at a velocity v_* . The stellar wind has a mass loss rate $\dot{m}_*$ and terminal velocity v_w . It flows freely from the star in all directions until it encounters a terminal shock. In the forward direction this shock occurs at a distance l from the star where the momentum flux in the wind equals the ram pressure of the ambient medium (Weaver *et al.* 1977),

$$l = 5.50 \times 10^{16} \dot{m}_{*, -6}^{1/2} v_{w, 8}^{1/2} \mu_H^{-1/2} n_{H, 5}^{-1/2} v_{*, 6}^{-1} \text{ cm}. \quad (2)$$

This scale agrees well with the sizes of compact H II regions observed by WC. In equation (2) $\dot{m}_{*, -6}$ is the stellar wind mass-loss rate in units of $10^{-6} M_\odot \text{ yr}^{-1}$, and $v_{w, 8} \times 10^8 \text{ cm s}^{-1}$ is the wind's terminal velocity.

In the absence of efficient cooling the thermalized wind would fill a region of roughly constant pressure whose thickness was estimated in VM to be near $0.5l$. We show below that cooling via conduction and subsequent copious radiation is so rapid that this layer is thin enough that the dynamics is governed by momentum conservation, considerably simplifying the calculations in Paper II. As the star moves through the ambient molecular medium, its wind sweeps material up into a thin paraboloidal shell exterior to the shocked wind region. The flow through the shell resembles the flow around a blunt body in a supersonic stream (Rusanov 1976 and references therein). Depending on the emission measure $\int n_e^2 dl$, all or part of the star's ionizing luminosity is trapped in the shell.

In the context of molecular clouds, the stand-off distance of the terminal wind shock from the star l is given in equation (2), while the column density of the swept-up shell in the neighborhood of the stagnation point N , the thickness of the shell δ assuming a radiative shock, and the velocity of swept-up material tangential to the shell are easily estimated once the shape of the bow shock is assumed. In VM, the shape was taken to be $y = x^2/l$ where x and y are the coordinates perpendicular and parallel to the direction of motion. Explicit

numerical work, presented in Paper II, shows that near the vertex the shape is instead closely described by the curve $y = x^2/3l$. This is similar to the results of Baranov, Krasno-baev, and Kulikovskii (1971) and Dyson (1975) for $x \ll l$.

It is easily shown that cooling via H_2 emission in the post-shock molecular cloud gas is rapid, so the outer shock can be considered radiative. We describe the flow through the shell in the neighborhood of the stagnation point using mass and momentum conservation. The rate at which mass is swept up by the bow shock interior to a point x is

$$\dot{m}_{\text{shock}}(x) = \pi x^2 \mu_H m_H n_H v_*. \quad (3)$$

Mass flows through the shell at a rate

$$\dot{m}_{\text{shell}}(x) = 2\pi x \sigma v_{\text{shell}, x}(x), \quad (4)$$

where the surface density of the shell is σ and $v_{\text{shell}, x}$ is the x -component of the flow velocity. In the steady state, these two mass fluxes are equal. We estimate $v_{\text{shell}, x}$ near the stagnation point by noting that the parallel component of momentum is preserved across the shock. We have

$$v_{\text{shell}, x}(x) \dot{m}_{\text{shell}}(x) = \int v_{\text{shock}, x}(x) d\dot{m}_{\text{shock}}(x). \quad (5)$$

We neglect the contribution from the stellar wind, which is of higher order in (x/l) since we are only interested in the region near $x = 0$. Equating $\dot{m}_{\text{shock}}(x)$ to $\dot{m}_{\text{shell}}(x)$ and solving simultaneously with equation (5), we find that the tangential flow velocity in the shell and surface density are

$$v_{\text{shell}, x}(x) = \frac{4}{9} v_* \frac{x}{l} \\ \sigma = \frac{9}{8} \mu_H m_H n_H l. \quad (6)$$

These values are different from the results of VM and should be considered to supersede those calculations. The difference comes about by the assumption made in VM about the shape of the shock in the vicinity of the stagnation point. The column density, $N = \sigma/(\mu_H m_H)$ and thickness $\delta = 9/8 \mathcal{M}^{-2} l$ ($\mathcal{M}$ is the Mach number) are then

$$N = 6.18 \times 10^{21} \dot{m}_{*, -6}^{1/2} v_{w, 8}^{1/2} \mu_H^{-1/2} n_{H, 5}^{1/2} \text{ cm}^{-2}, \quad (7)$$

$$\delta = 5.14 \times 10^{14} \gamma_I T_{I, 2} \dot{m}_{*, -6}^{1/2} v_{w, 8}^{1/2} v_{*, 6}^{-3} \mu_H^{-3/2} n_{H, 5}^{-1/2} \text{ cm}, \quad (8)$$

where γ_I is the ratio of specific heats in the neutral gas and $T_{I, 2}$ is its temperature in units of 100 K. It is instructive to compare the ionizing flux of the star $S_{49} \times 10^{49} \text{ photons s}^{-1}$ with the potential emission measure in the shell. Considering only the stagnation point, their ratio $\mathcal{T}$ (for trapped ionization front) is

$$\mathcal{T} = 4\pi l^2 n_e^2 \alpha \delta (T_{II}/T_I)/S \\ = 2.82 \times 10^{-1} \dot{m}_{*, -6}^{3/2} v_{w, 8}^{3/2} v_{*, 6}^{-1} n_{H, 5}^{1/2} T_{II, 4}^{-1} S_{49} \\ \times \mu_H^{-1/2} \gamma_I^{-1} \alpha_{-13}, \quad (9)$$

where $T_{II, 4} \times 10^4 \text{ K}$ is the temperature of the ionized gas and the hydrogen recombination rate to all levels but the ground state, α , is given in units of $10^{-13} \text{ cm}^3 \text{ s}^{-1}$. We neglected the effects of dust, which is an efficient absorber of stellar ionizing radiation. A normal dust-to-gas ratio causes UV absorption $A_{UV} = 2.2 \times 10^{-21} N \approx 9$ magnitudes across the shell. Including the dust is therefore very important in estimating $\mathcal{T}$. To approximate the effect of dust in models, S should be replaced with $0.927S/A_{UV}$ in all equations unless a radio continuum

measurement of S is available, otherwise the observed S should be used since it directly counts the ionization rate. The size of this effect corresponds roughly to the difference between radio continuum and *IRAS* determined values of S presented in WC. Conditions where $\mathcal{T}$ exceeds unity and the ionization front is trapped in the bow shock are high-density, low stellar velocity, and early spectral type. The latter holds because the scaling of stellar mass loss $\dot{m}_*$ with ionizing photon flux S is steeper than the first power for O stars.

The thickness of the ionized part of the shell can be estimated by taking a fraction $\mathcal{T}^{-1}$ of the material and determining how thick a layer that would form in pressure balance with the wind:

$$\begin{aligned}\delta_{HII} &= \delta \min(1, \mathcal{T}^{-1})(T_{II}/T_I) \\ &= 1.82 \times 10^{17} \dot{m}_{*, -6}^{-1} v_{w,8}^{-1} n_{H,5}^{-1} \mu_H^{-1} v_{*,6}^{-2} S_{49} \\ &\quad \times T_{II,4}^2 \alpha_{-13}^{-1} \gamma_I^2 \text{ cm},\end{aligned}\quad (10)$$

when $\mathcal{T} \geq 1$.

Expressed as a fractional thickness, and assuming a thin ionized layer,

$$\begin{aligned}\Delta_{HII} &= \delta_{HII}/l \\ &= 3.31 \gamma_I^2 \dot{m}_{*, -6}^{-3/2} v_{w,8}^{-3/2} n_{H,5}^{-1/2} v_{*,6}^{-1} \\ &\quad \times T_{II,4}^2 S_{49} \mu_H^{-1/2} \alpha_{-13}^{-1}.\end{aligned}\quad (11)$$

The structure will thus appear limb brightened when $\dot{m}$, v_w , n_H , or v_* are large, or when S is small or much dust is present. When the stellar wind passes through the terminal shock to come into pressure balance with the ram pressure of the molecular cloud, it thermalizes at a temperature

$$\begin{aligned}T_{\text{shock}} &= \frac{3}{16} \frac{\mu_H m_H}{k} v_w^2 \\ &= 2.255 \times 10^7 \mu_H v_{w,8}^2 \text{ K}.\end{aligned}\quad (12)$$

We now show that while radiative cooling alone is negligible, thermal conduction is effective in causing energy loss from the hot shocked wind by lowering its temperature to the point where radiative losses are rapid.

From 10^5 K to several times 10^8 K the dust-free radiative cooling function can be approximated to within a factor of 2 by (Gaetz and Salpeter 1983; Mac Low and McCray 1988),

$$\Lambda_R \approx \begin{cases} 1.58 \times 10^{-18} T^{-7/10} \text{ ergs cm}^3 \text{ s}^{-1}, & 10^5 \leq T \leq 10^7; \\ 3.79 \times 10^{-27} T^{1/2} \text{ ergs cm}^3 \text{ s}^{-1} & T > 10^7. \end{cases}\quad (13)$$

The cooling time at constant pressure (this region is isobaric) is then

$$\begin{aligned}t_{\text{cool}, R} &= \int_0^T \frac{3}{2} nk/n^2 \Lambda_R dT \\ &= 7.93 \times 10^{12} n_{H,5}^{-1} v_{*,6}^{-2} v_{w,8}^3 \mu_H^{1/2} \text{ s},\end{aligned}\quad (14)$$

where we take the number of particles per h-nucleus to be 23/9. This is considerably longer than the crossing time

$$\begin{aligned}t_{\text{cross}} &= l/v_* \\ &= 5.50 \times 10^{10} \dot{m}_{*, -6}^{1/2} v_{w,8}^{1/2} n_{H,5}^{-1/2} \mu_H^{-1/2} v_{*,6}^{-2} \text{ s},\end{aligned}\quad (15)$$

so the shocked wind would not have time to cool radiatively in the absence of other effects. Note that the cooling time given holds only for stars with wind shocks that thermalize above 10^7 K, i.e., wind velocities over 700 km s^{-1} .

Because a dense H II region at only 10^4 K bounds the hot shocked wind, there will be large temperature gradients so thermal conduction can lower the temperature of the hot gas. This does not in itself lead to an energy loss from the region. There is a flow of material from the swept-up shell into the hot shocked wind in such a way that nT remains nearly constant. We refer to this process as thermal evaporation of the shell; it dilutes the temperature of the hot gas. Provided enough thermal evaporation takes place, the temperature is reduced to a point near the peak of the cooling function and the thermal energy is quickly radiated away. We call this *poisoning* of the hot gas.

If the temperature gradient is steep enough, the heat transport saturates because the thermal energy of the gas can be transported by at most the sound velocity. Cowie and McKee (1977) define a saturation parameter,³ which if greater than unity, means that the conduction is saturated,

$$\sigma_0 = \frac{\kappa T_e}{5 \phi_s \mu_H m_H n_H c_s^3 L_T}.\quad (16)$$

In this equation κ is the thermal conductivity, $5.94 \times 10^{-7} T^{5/2} (31/\Lambda) \text{ ergs s}^{-1} \text{ K}^{-1} \text{ cm}^{-1}$ (Spitzer 1982), Λ is the Coulomb logarithm, $T_e = T_{\text{shock}}$ is the electron temperature, ϕ_s is a dimensionless factor of order unity, and L_T is the temperature scale length. We choose for the latter the upper limit given by the ratio of the thickness of the shocked wind given by VM to the number of e -foldings of the temperature profile,

$$\begin{aligned}L_T &= 0.47l/\ln(T_e/T_{II}) \\ &= 3.33 \times 10^{15} \dot{m}_{*, -6}^{1/2} v_{w,8}^{1/2} \mu_H^{-1/2} n_{H,5}^{-1/2} v_{*,6}^{-1} \lambda,\end{aligned}\quad (17)$$

where for convenience we define

$$\lambda = [0.1295 \ln(\mu_H v_{w,8}^2 T_{II,4}) + 1]^{-1},\quad (18)$$

which usually takes on values between 0.7 and 1.0. In the bow shocks of O stars in molecular clouds the temperature scale length is so short that thermal conduction is strongly saturated:

$$\sigma_0 = 2.70 \times 10^2 \mu_H^3 n_{H,5}^{-1/2} v_{w,8}^{11/2} v_{*,6}^{-1} \lambda^{-1} \phi_s^{-1}.\quad (19)$$

If we approximate the bow shock's shape as a hemisphere for convenience, then thermal evaporation of the shell proceeds at a rate

$$\begin{aligned}\dot{m}_{\text{evap}} &= 2\pi l^2 \mu_H m_H n_H c_s \phi_s F(\sigma_0) \\ &= 1.62 \times 10^{21} \dot{m}_{*, -6}^{13/16} v_{w,8}^{33/16} \phi_s^{5/8} v_{*,6}^{-3/8} \\ &\quad \times n_{H,5}^{+3/16} \lambda^{-3/8} \text{ g s}^{-1},\end{aligned}\quad (20)$$

where $F(\sigma_0) = 2.73 \sigma_0^{3/8}$ (Cowie and McKee 1977). This rate is one to two orders of magnitude greater than the rate at which hot gas is produced at the wind terminal shock, so we expect that thermal dilution will be very important. To estimate the temperature of the conduction zone, we simply mix the two gases

$$\begin{aligned}T_{\text{cool}} &= T_{\text{shock}} \frac{\dot{m}_*}{(2\dot{m}_{\text{evap}} + \dot{m}_*)} \approx T_{\text{shock}} \frac{\dot{m}_*}{z\dot{m}_{\text{evap}}} \\ &\approx 4.38 \times 10^5 \dot{m}_{*, -6}^{3/16} v_{w,8}^{-1/16} \phi_s^{-5/8} v_{*,6}^{3/8} n_{H,5}^{3/16} \lambda^{+3/8} \mu_H^{-1/8} \gamma^{1/16} \text{ K}.\end{aligned}\quad (21)$$

³ While the theory of Cowie and McKee does not properly describe the nonlocal nature of electron heat condition in steep thermal gradients, it does provide a good estimate of the global heat flux.

The cooling function corresponding to this temperature is equation (13):

$$\begin{aligned}\Lambda_{\text{cool}} &= 1.58 \times 10^{-18} T_{\text{cool}}^{-7/10} \text{ ergs cm}^3 \text{ s}^{-1} \\ &= 1.78 \times 10^{-22} \dot{m}_{*, -6}^{-21/160} n_{\text{H}, 5}^{-21/160} v_{*, 6}^{-21/80} \\ &\quad \times \mu_{\text{H}}^{7/80} v_{\text{w}, 8}^{7/160} \phi_s^{35/80} \lambda^{-21/80} \gamma^{-7/160} \text{ ergs cm}^3 \text{ s}^{-1}, \quad (22)\end{aligned}$$

so the mixed layer cools in a very short time

$$\begin{aligned}t_{\text{cool}} &= \int_0^T \frac{3}{2} nk/n^2 \Lambda \, dT \\ &= 3.35 \times 10^8 \dot{m}_{*, -6}^{81/160} n_{\text{H}, 5}^{-79/160} \phi_s^{135/80} v_{*, 6}^{-79/80} \\ &\quad \times v_{\text{w}, 8}^{-27/160} \lambda^{81/80} \mu_{\text{H}}^{-107/80} \gamma^{27/80} \text{ s}, \quad (23)\end{aligned}$$

where values appropriate for the cooling zone are used. The cooling layer has a thickness sufficient to radiate the kinetic energy deposited by the wind,

$$\begin{aligned}\delta_{\text{cool}} &= \frac{\dot{m}_* v_w^2}{4\pi l^2 n_{\text{cool}}^2 \Lambda_{\text{cool}}} \\ &= 1.24 \times 10^{16} \dot{m}_{*, -6}^{81/160} v_{\text{w}, 8}^{133/160} n_{\text{H}, 5}^{-79/80} v_{*, 6}^{-81/160} \\ &\quad \times \phi_s^{-135/80} \mu_{\text{H}}^{-17/8} \lambda^{-81/80} \gamma^{27/160} \text{ cm}, \quad (24)\end{aligned}$$

The conduction/cooling layer is thinner than the equivalent nonradiative shocked wind region. Since it is a relatively narrow shell, any optically thin emission arising in this zone will be limb brightened. Note that the above estimate of δ_{cool} is an upper limit since a lower bound on σ_0 was used and radiative cooling during the mixing was neglected.

Cooling rates in a hot plasma are strongly dependent on the presence of dust grains. Energy is lost from the gas via inelastic collisions with dust grains, which then radiate in the infrared. An enhancement in the cooling by as much as a factor of 100 is possible for temperatures of 10^7 – 10^8 K (Dwek 1987). Dust which enters the cooling zone in the evaporative flow will survive for a fairly long time since the sputtering yields are quite low. The reason for this is that energy is deposited all along the path of a fast particle through a grain and rapidly redistributed into a large number of internal excitations, no one of which is energetic enough to eject a grain atom from the surface. For temperatures greater than 10^6 K the grain survival time is $10^6 a_{\text{um}}/n$ yr (Draine and Salpeter 1979; Seab 1986), comparable to the cooling time without dust given by equation (14). Figure 5 of Dwek (1987) shows that the presence of dust enhances the cooling by a factor of approximately $2(T/10^6 \text{ K})^{4/3}$ for temperatures above 10^6 K. Consequently, several times as much energy is radiated by regions hotter than T_{cool} than those colder, reducing the cooling length by a factor of a few from that given in equation (24). A detailed calculation of the temperature profile is required to assess this effect accurately.

Velocity shears are present between the shocked wind and the H II region, allowing Kelvin-Helmholtz instabilities to grow. If these become nonlinear and evolve into “cat’s eyes” that entrain 10^4 K gas in the hot material, further mixing and enhanced cooling will occur. Acoustic instabilities, discussed by Breitschwerdt and Kahn (1988), grow on a sound crossing time for the conduction layer and can also lead to enhanced mixing. On the other hand, the presence of magnetic fields will tend to inhibit thermal conduction, and hence the mixing of dust into the hot gas. Because the saturation parameter is so large, this may be only a minor effect.

Before concluding this section, we remind the reader that the theory presented here applies only to the neighborhood of the stagnation point in the flow.

IV. CONFRONTING THE MODEL WITH OBSERVATIONS

a) Detailed Description of G34.3+0.2

We now turn to the observed properties of the prototypical object, G34.3+0.2. A deep 8.4 GHz VLA map is shown in Figure 1. Important characteristics of cometary compact H II regions which G34.3+0.2 shares are

- (1) an ionization front well described by a parabola,
- (2) surface brightness greatest at the coma (head),
- (3) limb brightening,
- (4) an extremely sharp “upstream” edge,
- (5) radio recombination lines which are unusually broad.

Churchwell, Wood, and Bieging (1990, hereafter CWB) model the high-resolution radio continuum flux distribution and radio recombination line emission to find an electron temperature $T_e = 9000 \pm 500$ K, and rms electron density of $\langle n_e \rangle = 8.75 \pm 0.1 \times 10^4 \text{ cm}^{-3}$ and a maximum emission measure of $5.6 \times 10^8 \text{ pc cm}^{-6}$. The electron density deduced from recombination lines is $2 \times 10^5 \text{ cm}^{-3}$, implying that the ionized material has density gradients. Line widths are large, H76 α is $\approx 49 \text{ km s}^{-1}$ wide at any given point in the object (Garay, Rodriguez, and van Gorkom 1986). The H42 α has the same width integrated over a $3''$ beam (CWB). Agreement between these two line widths argues against impact broadening so they likely indicate supersonic relative motions in the gas.

Ammonia has been observed in emission toward G34.3+0.2 by Heaton *et al.* (1985), Anderson and Garay (1986), Henkel, Wilson, and Mauersberger (1987), CWB, and Keto, Ho, and Reid (1990). HCN has also been observed in absorption against the H II region by CWB who found a striking difference between the NH₃ and HCN absorption. The HCN hyperfine components overlap to produce a very broad $\approx 20 \text{ km s}^{-1}$ and flat-bottomed profile, while the NH₃ line is narrow, $\approx 3 \text{ km s}^{-1}$ wide. Because the HCN is very optically thick this tells us that there is molecular material spanning a velocity range of greater than 9 km s^{-1} , the approximate intrinsic HCN line width. The NH₃ absorption, on the other hand, is optically thick only in the very densest material so it traces only the dense swept-up shell.

The systemic radial velocities of the ionized gas and the molecular gas are 45 km s^{-1} and 60 km s^{-1} respectively (CWB), indicating a relative motion of these components toward each other since the absorption line gas lies in front of the H II region. Where the two meet, there must be a very strong shock because the $\approx 15 \text{ km s}^{-1}$ velocity difference is highly supersonic in the molecular gas. It is unlikely that gravity could induce such motions, since the interior mass required is in excess of $600 M_\odot$, much greater than the stellar mass ($\approx 30 M_\odot$) and several times larger than that of a cluster with a normal initial mass function whose most massive star is $30 M_\odot$. The mass of the H II region gas is insignificant in comparison. Another possibility is that the relative motion is due to the expansion of the H II region or stellar wind bubble surrounding the star. In this case the structure would only have an age of 600 yr, much too short to account for the population statistics of compact cometary H II regions. We are thus led to argue that the relative motions represent a steady-state situation in which the ionized gas really is traveling through the molecular material at supersonic speeds.

A beautiful piece of evidence that supports the idea of a strong shock at the sharp leading edge of the ionization front is presented by CWB. The OH masers observed by Gaume and Mutel (1987) are located just upstream of the ionization front and lie along its whole leading edge. Densities greater than 10^7 cm^{-3} are required to produce these masers. Our interpretation is that the OH masers mark the location of the strong radiative shock which separates the cooler quiescent molecular cloud core from the warm, very dense NH_3 shell comprising the outer part of the bow shock.

b) Applying the Simple Theory to G34.3+0.2

We adopt the following values of parameters for G34.3+0.2, where the total recombination rate suggests an O6.5 V star: $\dot{m}_{*, -6} = 1$, $v_{w, 8} = 3$, $S_{4.9} = 0.55$ (WC, Panagia 1973), $n_{H, 5} = 3$, $\mu_H = 1.4$, $T_{II, 4} = 1$, and a distance of 3.7 kpc. Furthermore, we assume that the star is located at $4/3$ the vertex to focus distance of a parabola defining the outer edge of the H II region and that we are viewing the structure perpendicular to v_* , i.e., the direction of the star's motion is in the plane of the sky. The measured distance l is then $3.1 \times 10^{16} \text{ cm}$, which implies via equation (2) that $v_{*, 6} < 1.5$, where the upper limit arises from projection effects. This is to be compared with the total range of velocities seen in the molecular gas of $\sim 9\text{--}10 \text{ km s}^{-1}$ and the velocity difference of 15 km s^{-1} between the ionized and molecular gas (CWB). Both these velocities should be of order v_* .

There is no measurement of the total column density of the swept-up shell to compare directly with the estimate of equation (7), but we can use the NH_3 observations of CWB if a fractional abundance of NH_3 is assumed. Calculations of molecular cloud chemistry which include instantaneous heating of grains once the density reaches 10^7 cm^{-3} were presented by Brown, Charnley, and Millar (1988). These conditions correspond very well to the passage of molecular material through the bow shock. Nitrogen becomes almost completely hydrogenated as NH_3 , so the NH_3 abundance relative to hydrogen nuclei is equal to abundance of nitrogen nuclei, taken to be solar (1.18×10^{-4}). We do not expect significant dissociation of NH_3 behind a less than 15 km s^{-1} shock because radiative depopulation of collisionally excited states is very rapid (Dalgarno and Roberge 1979). Consequently, the NH_3 column of the bow shock along the direction of motion is

$$N_{\text{NH}_3} = 7.3 \times 10^{17} \dot{m}_{*, -6}^{1/2} v_{w, 8}^{1/2} v_{*, 6}^{-1} \mu_H^{-1/2} n_{H, 5}^{1/2} \text{ cm}^{-2}. \quad (25)$$

The predicted column for G34.3+0.2 is $7.1 \times 10^{17} \text{ cm}^{-2}$, to be compared with the observation of CWB, 6.3×10^{17} (but along a different line of sight). We assume here that $\mathcal{T}$ is greater than a few. The gas-phase abundance of NH_3 in the preshock gas is likely to be very low, $\sim 10^{-8}$ (Brown, Charnley, and Millar 1988), due to freeze out onto grains. The NH_3 column contributed by the quiescent cloud, whose total H_2 column density is $4 \times 10^{23} \text{ cm}^{-2}$ is therefore negligible.

By applying equation (9) to find the trapping parameter $\mathcal{T}$, we find that the ionization front is trapped, $\mathcal{T} = 6.2$, where the radio-derived value of $S_{4.9}$ is used. Radio continuum emission is expected to have a sharp edge on the upstream side, as is observed. The not very definitive value means that there could be bow shocks in molecular clouds for which $\mathcal{T}$ is less than unity. In that case the appearance will be amorphous, corresponding to (possibly clumpy) ionized quiescent cloud material surrounding the completely ionized bow shock. The shock could then be quite weak since the sound speed in the ionized

gas is $\sim 10 \text{ km s}^{-1}$. A typical member of this class may be W51 D.

In G34.3+0.2 stellar ionizing photons are able to maintain an H II region of fractional thickness $\Delta_{HII} = 0.061$ of l , corresponding to $0''.026$. The ionized shell's thickness is unresolvable with the VLA at 2 cm, and indeed, appears unresolved on the maps of WC, which have a resolution of $\approx 0''.4$.

Finally, to conclude our comparison of the analytic model with G34.3+0.2, we find that the velocity dispersion, δv , of ionized gas in the leading hemisphere of a bow shock is

$$\delta v \approx [458(T_{II, 4} + 1.25v_{*, 6}^2) + (31.4T^{2/10}n_{II, 5})^2]^{1/2} \text{ km s}^{-1}, \quad (26)$$

where use has been made of equation (5) and equations (10–11) of Garay, Rodriguez, and van Gorkom (1986). In deriving this equation we assume that Kelvin-Helmholtz instabilities have produced turbulence with a velocity $\approx (4/9)v_*$. Taking $v_* = 15 \text{ km s}^{-1}$ and $n_{II, 5} = 2$, we find $\delta v \approx 43 \text{ km s}^{-1}$, in reasonable agreement with the observed values $\approx 50 \text{ km s}^{-1}$ (Garay, Rodriguez, and van Gorkom 1986).

V. PROPOSED CONFIRMING OBSERVATIONS

The bow-shock model makes specific unique predictions regarding the proper motions of ionized and molecular material. The H II region moves along the symmetry axis in the direction of the coma. As discussed below, this pattern speed will be detectable with observations over a multiyear baseline using the VLA A-array at 2 cm. Associated OH masers, which are located in the shock-compressed flow around the ionized region, move away from the symmetry axis at velocities given by equation (5). VLBI observations taken only a few years apart will easily show this motion. Finally, the region of the conductive interface between the hot shocked stellar wind and inner ionized portion of the swept-up shell will radiate of order $10^3 L_\odot$ in cooling lines, some of which are in the infrared and hence observable through the large dust columns that are present. Infrared lines are also expected from the dense shocked shell comprising the outer portion of the bowshock.

a) Pattern Motion

The entire structure moves at a velocity v_* in the direction of the coma. Let the motion be inclined at an angle ι to the line of sight, then the proper motion of the structure is

$$\dot{\theta} = 2.1 \times 10^{-3} v_{*, 6} D_{\text{kpc}}^{-1} \sin \iota \text{ arcsec yr}^{-1}, \quad (27)$$

where D_{kpc} is the distance to the object in kiloparsecs. The VLA in A-array, observing at 2 cm, can detect angular motions of order $0.005 \text{ arcsec yr}^{-1}$ in high dynamic range images, extrapolating from the experience of Masson (1986). Thus a typical compact cometary H II region at a distance of 5 kpc should have a detectable pattern motion over a 10 yr baseline. In order to plan VLA observations of these objects, we have produced simulated images (including noise) of a moving step function at different epochs. These were then subtracted and the peak signal, which indicates the motion, was measured relative to the thermal noise in the simulated maps. These results are fitted approximately by the formula

$$s = 0.54 \text{ SNR } (\omega/\omega_{\text{Beam}})^{1/2} (\dot{\theta}t/\omega_{\text{Beam}}) \quad (28)$$

where SNR is the mean signal-to-noise ratio over the leading edge of the coma, ω is its angular scale, ω_{Beam} is the beam width, $\dot{\theta}t$ is the angular displacement between the two epochs,

and s is the significance of the result in units of the rms noise. A typical time baseline to get a 10σ result in the A-array at 2 cm with a signal-to-noise ratio of 1000 in the epoch maps is roughly $v_{*,6}^{-1} D_{\text{kpc}}$ years. This is probably an underestimate because the step-function model is unrealistic for true sources. In actuality, there will be some structure and the edge will not be perfectly sharp. Doubling or tripling the time baseline is probably safer. One last point on this experiment. Because differential astrometry with the VLA is limited by the accuracy to which reference positions are stable, this experiment can only be done if a strong enough background source is located within the primary beam. The background source then provides a reference direction. It should be strong enough so that its detection is at a signal-to-noise ratio comparable to that of the H II region.

b) Maser Proper Motions

If the masers are clumps in the shocked shell, and not in the preshock material, then they will trace the streamlines of the flow. Figure 3 shows the predicted velocity field of the G34.3+0.2 masers assuming $v_* = 10 \text{ km s}^{-1}$, our line of sight is inclined 90° from the direction of motion and the masers we observe are all located on the near side of the shell. The amplitudes of the motions (a few milliarcseconds per yr) and the strengths of the masers (0.1 Jy) make their mapping a very interesting and straightforward VLBI experiment. In the best case, there will be a strong continuum source located within the primary beam to serve as a reference position. Otherwise only the relative motions of the masers, up to an arbitrary vector in the plane of the sky are measurable. Observations separated with baselines of $\approx$ one year over a multiyear period would convincingly confirm or rule out our hypothesis.

c) Coronal Infrared Fine-structure Lines

In the conduction zone between the hot shocked wind and the H II shell some of the cooling occurs via line emission. Most

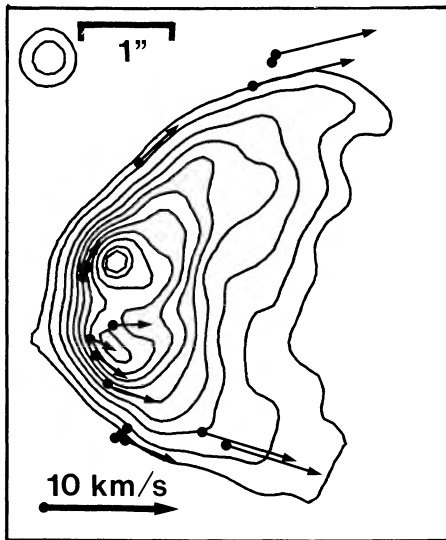


FIG. 3.—Predicted OH maser velocity field for G34.3+0.2 assuming the bow-shock hypothesis is true. Dots show maser positions given by Gaume and Mutel (1987); arrows are the predicted proper motions for $v_{*,6} = 1$. All masers are assumed to be on the near side of the ionized structure and are known to be approaching the observer at velocities intermediate between the recombination line velocity and molecular cloud velocity (most positive).

of the line luminosity arises from transitions with $\lambda \approx hc/kT$, at which wavelengths the molecular cloud is opaque. A small amount of the emission will be in infrared lines longward of several microns where the cloud becomes transparent. As an example, consider the infrared fine structure lines of Fe VII at $7.81 \mu\text{m}$ and $9.51 \mu\text{m}$, which arise in a plasma with $T \approx 5 \times 10^5 \text{ K}$. The critical density for these lines is $n_{\text{crit}} \approx 2 \times 10^7 \text{ cm}^{-3}$ according to the atomic data given by Nussbaumer and Storey (1982), while the density in the cooling region, $n_{\text{cool}} = n_{\text{H}} \mu_{\text{H}} m_{\text{H}} v_w^2 / kT_{\text{cool}} \approx 3 \times 10^3 \text{ cm}^{-3}$. Virtually all the Fe VII ions are in the ground state, so every collision into an upper state results in an emitted photon. In these circumstances the line luminosity is

$$L = h\nu n_e n_{\text{FeVII}} \gamma 2\pi l^2 \delta_{\text{cool}}, \quad (29)$$

where $h\nu$ is the energy of an emitted photon, n_e is the electron density, n_{FeVII} is the number density of Fe VII, and γ is now the collisional excitation rate to conform with standard notation, which in terms of the collision strength Ω and statistical weight g is

$$\gamma = 8.63 \times 10^{-6} \Omega / g T^{1/2} \text{ cm}^3 \text{ s}^{-1} \quad (30)$$

for a Maxwellian particle distribution. Adopting fiducial values for all the parameters of the problem, we find

$$\begin{aligned} L(9.51 \mu\text{m}) &\approx 5 \times 10^{34} \text{ ergs s}^{-1} \\ L(7.81 \mu\text{m}) &\approx 8 \times 10^{34} \text{ ergs s}^{-1}. \end{aligned} \quad (31)$$

At a distance of 5 kpc, a typical object will have a line flux of $\sim 10^{-14} \text{ W m}^{-2}$, below the IRAS Low Resolution Spectrometer detection limit. A spectral resolution of $\lambda/\Delta\lambda = 3000$ or greater will match the thermal line widths, yielding a peak flux density of order 100 Jy. These line strengths are likely to be somewhat overestimated since only a fraction of the material is at the temperature T_{cool} . An additional uncertainty arises from nonequilibrium ionization effects, which can alter line strengths by an order of magnitude from equilibrium values.

These fluxes are high enough that a ground-based instrument (for lines in the 8–13 μm window) with a large aperture will readily yield a detection in only a few second exposure with a spectral resolution that matches the thermal line width. There are a large number of coronal lines, characteristic of a 10^5 – 10^6 K plasma, available. Their detection will confirm the presence of strong stellar winds interacting with dense surroundings.

d) Infrared Emission from the Outer Shock

In the cooling region behind the outer shock, infrared emission from H_2 and O I is expected. The results of Draine, Roberge, and Dalgarno (1983) show that both the $9.665 \mu\text{m}$ H_2 line and the $63.2 \mu\text{m}$ O I line should have surface brightnesses $\sim 10^{-3} \text{ ergs cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$. In terms of luminosities, this is several $10^{32} \text{ ergs s}^{-1}$, considerably weaker than the coronal lines, but still observable with current instrumentation. Detection of these lines will confirm the presence of a $\sim 10 \text{ km s}^{-1}$ shock in the molecular material, i.e., the bowshock itself. We note here that the $2.1 \mu\text{m}$ H_2 luminosity of ultracompact H II regions, when detected, is of order $0.1 L_\odot$ (Olivia and Moorwood 1986), corresponding to $v_{*,6} \approx 2$. This is perhaps a large value, but the presence of density fluctuations in the molecular cloud could easily lead to such shock velocities as the leading edge of the bow shock flutters about.

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