FROM ULTRACOMPACT TO EXTENDED H \textsc{ii} REGIONS. II. CLOUD GRAVITY AND STELLAR MOTION

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ABSTRACT

The dynamical evolution of H \textsc{ii} regions with and without stellar motion in dense, structured molecular clouds is studied. Clouds are modeled in hydrostatic equilibrium, with Gaussian central cores and external halos that obey $\rho \propto r^{-2}$ and $\rho \propto r^{-3}$ power laws. Cloud gravity is included as a time-independent, external force. Stellar velocities of 0, 2, 8, and 12 km s$^{-1}$ are considered, permitting stars to move from the central core toward the edge of the cloud. Ultracompact H \textsc{ii} regions are seen to evolve into extended H \textsc{ii} regions as the stars move toward lower density regions. Our main conclusion is that ultracompact H \textsc{ii} regions are pressure-confined entities while they remain embedded within dense cores. The confinement comes from either ram or ambient pressures, or a combination of both. The survival of the ultracompact regions depends on the position of the star with respect to the core center, the stellar lifetime, and the crossing time of the cloud core. Stars with velocities less than the cloud dispersion velocity can produce cometary ultracompact H \textsc{ii} regions for $2 \times 10^4$ yr or more, in statistical agreement with observations. The sequence ultracompact H \textsc{ii} → compact H \textsc{ii} → extended H \textsc{ii} shows a variety of structures induced by various instabilities. Some ultracompact H \textsc{ii} regions with a core-halo morphology could be explained by self-blocking effects, when stars overtake and ionize leading, piled-up clumps of neutral gas.

Subject headings: hydrodynamics — ISM: clouds — H \textsc{ii} regions

1. INTRODUCTION

The understanding of H \textsc{ii} region dynamics has evolved substantially since the seminal works of Strömgren (1939), Kahn (1954), and Savedoff & Greene (1955). Both analytical and numerical treatments have grown increasingly sophisticated, taking into account the effects of density stratifications (e.g., Tenorio-Tagle 1979; Yorke 1986; Franco et al. 1990; Shu et al. 2002; Henney et al. 2005), turbulent density distributions (Li et al. 2004; Meléndez et al. 2006), gravity (Keto 2002; González-Áviles et al. 2005), stellar winds (Comeron 1997; Arthur & Hoare 2006), moving stars (Rafwala 1969; van Buren et al. 1990; Arthur & Hoare 2006), magnetic fields (Krumholz et al. 2007) and instabilities of various classes (Giuliani 1979; García-Segura & Franco 1996; Williams 2002). A wealth of new information has been produced, which can be used to interpret the observations, none of which conform to the textbook picture of spherical, uniformly expanding H \textsc{ii} regions. See Henney (2006) for a recent review of H \textsc{ii} region dynamics.

Along the way, observational studies have identified new classes of H \textsc{ii} regions, smaller and denser than the large, optical H \textsc{ii} regions that motivated earlier studies. These compact, ultracompact, and hypercompact H \textsc{ii} regions represent earlier evolutionary stages and are generally still embedded in the natal molecular cloud, hence observable only at infrared and radio frequencies (e.g., Wood & Churchwell 1989; Kurtz et al. 1994; Walsh et al. 1998; Hanson et al. 2002; Sewilo et al. 2004; de Pree et al. 2005). The relative youth of these compact nebulae suggests that the star formation process and the circumprotostellar environment should be important for their dynamics, unlike the case of evolved H \textsc{ii} regions, for which the dynamics are dominated by the structure of the interstellar medium at much larger scales.

Low angular resolution radio observations indicate that many ultracompact (UC) H \textsc{ii} regions are surrounded by low-density ionized gas (Mezger et al. 1967; Garay et al. 1993; Koo et al. 1996; Kurtz et al. 1999; Kim & Koo 2001). The physical relationship between the extended and the compact components is unclear; the evidence for a relationship is morphological in most cases. Kim & Koo report radio recombination line observations that support the hypothesis of a relation based on close agreement of radio recombination line velocities between the ultracompact and extended components.

Further work, both observational and theoretical, is needed to understand the relationship between the compact, high-density gas and the extended, low-density gas. The Orion Nebula, the archetypal H \textsc{ii} region, has ionized gas ranging from $n_e$ of about 400 cm$^{-3}$ to over $10^5$ cm$^{-3}$ on scales from parsecs to tens of milliparsecs (Felli et al. 1993); realistic models should explain the coexistence of such different densities on different size scales. The location of a young massive star near the edge of its natal dense core—which itself may be located near the edge of a lower density clump—may provide the setting to produce ionized gas at a range of densities on a range of size scales. A general scheme based on this idea was suggested by Franco et al. (2000) and is shown pictorially in Figure 8 of Kim & Koo (2001). A basic question to explore is whether the diffuse and compact components owe their origin to the ionization of an intrinsically density-structured medium, or whether they arise from the dynamical evolution of the
ionized gas. Indeed, both origins are possible, and reality may be a combination of the two.

Using the free-free radio-continuum spectral index of three UC H II regions, Franco et al. (2000) found power-law density distributions with exponents between $-1.5$ and $-4$. These gradients are steeper than those found by molecular line studies of star-forming clouds (which show exponents between $-1.5$ and $-2$, typical of equilibrium isothermal clouds). The reasons for these steeper exponents are unclear. However, Franco et al. point out that the extreme value derived (close to $-4$) may be an overestimate of the actual density gradient. The observed density stratifications in these star-forming cores are due to the gravitational fields of the cores themselves. On even more compact scales the stellar gravity will cause a more local density stratification (Keto 2003).

The location of a massive star with respect to these density structures can affect the dynamical evolution of its H II region. Franco et al. (1990) showed that for stars forming in the center of spherical clouds with a halo density distribution steeper than $r^{-3/2}$, then once the ionization front reaches the edge of the core the whole cloud becomes ionized, i.e., no static equilibrium solution exists for the position of the ionization front. For stars off-center from the center of steep density distributions, a champagne or blister-type H II region can result (Tenorio-Tagle 1979; Arthur & Hoare 2006). However, once formed, a massive star will move within the gravitational potential of the natal cloud, i.e., it will not remain in a fixed position with respect to the density gradient. For a deeper understanding of the early evolution of H II regions in such environments it is therefore necessary to include both gravity and stellar motion in the numerical simulations.

In this paper we present numerical simulations of the formation and evolution of H II regions assuming power-law density gradients in the ambient cloud material and including stellar motion. We consider power-law exponents of $-2$ and $-3$, consistent with observations of star-forming cores. An imposed spherical gravitational field and hydrostatic pressure distribution give the cloud its initial structure and the ionizing source is allowed to move up or down the density gradient. Such models allow us to investigate the combined effects of ram and thermal (or other) pressures on the H II region evolution. García-Segura & Franco (1996, hereafter Paper I) used thermal pressure as an illustration of the impact that any isotropic pressure may have on the dynamical evolution of H II regions (e.g., turbulent pressure), and we adopt the same ersatz pressure in the present paper, where we focus particularly on cases where the total ambient pressure is similar to the ionized gas pressure. The inclusion of stellar motion introduces a preferential direction with respect to the ambient pressure gradient and gives rise to a variety of different situations as the star leaves its parental cloud core. A wide range of morphologies is encountered for extended H II regions, and we confirm the appearance of sawtooth ionization-shock fronts discussed in Paper I and Williams (2002).

This paper also presents the results of calculations of H II regions forming off-center in spherical clouds with high core densities ($\sim 10^7$ cm$^{-3}$). We find that a large number of ionizing photons is required for complete core ionization and our simulations admit the possibility that the outer core material either remains neutral or recombines at a late evolutionary stage.

The structure of the paper is as follows: § 2 describes the models and the assumptions made, § 3 explains and shows the gas dynamical simulations, and § 4 discusses the results and summarizes our conclusions.

2. INITIAL CONSIDERATIONS

We assume a cloud consisting of an internal core of radius $r_c$ and a halo in hydrostatic equilibrium, with a constant total velocity dispersion ($P = \rho c_s^2$, and $c_0$ = constant),

$$\nabla P = -\rho g \Rightarrow \nabla \rho = -\frac{\rho}{c_s^2} g_c.$$  \hspace{1cm} (1)

For simplicity, due to computational restrictions, we assume that the gas is always described by a single equation of state and the total velocity dispersion is set equal to the isothermal sound speed, $c_s$. Thus, the resulting effective temperature is representative of the kinetic energy required to provide support against gravity, and is well above any actual molecular cloud temperature (i.e., the contribution from turbulent pressure is included in an implicit manner). Assuming that the cloud is spherically symmetric ($\nabla \rightarrow d/dr$ and $g \rightarrow g_c$), we solve only along the radial coordinate:

$$\frac{d\rho}{dr} = -\frac{\rho}{c_s^2 g_c}. \hspace{1cm} (2)$$

For the first set of simulations (designated model A) we assume that the halo density falls off as an $r^{-2}$ power law, as in a self-gravitating isothermal cloud:

$$\rho(r) = \rho_c (r/r_c)^{-2} \hspace{1cm} \text{for} \hspace{0.2cm} r \geq r_c. \hspace{1cm} (3)$$

Solving for $g_c$, with equations (2) and (3), we find

$$g_c = \frac{2c_s^2}{r_c} \hspace{1cm} \text{for} \hspace{0.2cm} r \geq r_c. \hspace{1cm} (4)$$

Inside the core, $g_c$ must grow from zero up to $2c_s^2/r_c$. A simple way to achieve this is to let $g_c$ grow linearly within the core, as $g_c = A(r/r_c)$. In order to join the halo solution (eq. [3]), we require that $A = 2c_s^2/r_c$.

The density distribution inside the core is found by integrating equation (2) with the new $g_c$, giving

$$\rho(r) = \rho_0 \exp\left[-(r/r_c)^2\right] \hspace{1cm} \text{for} \hspace{0.2cm} r \leq r_c, \hspace{1cm} (5)$$

where $\rho_0$ is the central density at $r = 0$, and $\rho_c = \rho_0/e$.

In summary, the density distribution of the cloud is given by

$$\rho(r) = \begin{cases} 
\rho_0 \exp\left[-(r/r_c)^2\right] & \text{for} \hspace{0.2cm} r \leq r_c, \\
\rho_0/e(r/r_c)^{-2} & \text{for} \hspace{0.2cm} r \geq r_c,
\end{cases} \hspace{1cm} (6)$$

and the gravitational acceleration by

$$g_c = \begin{cases} 
2c_s^2/(r_c^2) & \text{for} \hspace{0.2cm} r \leq r_c, \\
2c_s^2/(r_c^2) & \text{for} \hspace{0.2cm} r \geq r_c.
\end{cases} \hspace{1cm} (7)$$

For the second set of simulations (designated model B) we assume that the halo density falls off as an $r^{-3}$ power law, and using the above approach we find the density distribution

$$\rho(r) = \begin{cases} 
\rho_0 \exp\left[-3/2(r/r_c)^2\right] & \text{for} \hspace{0.2cm} r \leq r_c, \\
\rho_0/e^{3/2}(r/r_c)^{-3} & \text{for} \hspace{0.2cm} r \geq r_c.
\end{cases} \hspace{1cm} (8)$$
and the gravitational acceleration

\[ g_r = \begin{cases} 
3c_s^2/r_c(r/r_c) & \text{for } r \leq r_c, \\
3c_s^2/r_c(r/r_c)^{-1} & \text{for } r \geq r_c.
\end{cases} \tag{9} \]

For both sets of simulations we assume that the star was born in situ, inside the core of its parental cloud (modeling of “runaway” stars is not considered). The largest expected stellar velocity that such a cloud can produce is given by the dispersion velocity in hydrostatic equilibrium. Noting that (see § 2.1 in Paper I) \( P = 16/9 \pi \rho_c r_c^2 \), and solving for \( c_s \), we find

\[ c_s = 4.07r_{0.1}n_{0.2}^{1/2} \text{ km s}^{-1} \tag{10} \]

where \( r_{0.1} \) is the core radius in units of 0.1 pc and \( n_0 \) the core density in units of \( 10^6 \text{ cm}^{-3} \). Thus, for core densities of \( 10^7 \text{ cm}^{-3} \) and radii of 0.1 pc, stellar velocities up to \( \sim 13 \text{ km s}^{-1} \) can be considered. The actual value of \( c_s \) used in the simulations corresponds to the sound speed for 100 K molecular gas.

3. TWO-DIMENSIONAL GAS DYNAMICAL SIMULATIONS

3.1. Overview of the Simulations

The numerical simulations are performed with the gas dynamical MHD code ZEUS-3D version 3.4 (Stone & Norman 1992; Clarke 1996; see Paper I for details). We use Cartesian coordinates with the \( y \)-axis being the symmetry axis (i.e., a slab geometry). Thus, the star can be placed anywhere in the \( x-z \) plane. This is a particularly safe choice because it does not introduce any axis artifact in the computations (as would cylindrical coordinates).

The setup is similar for all models: the star is fixed on the computational mesh and remains at the same location during the computation. The stellar motion is simulated by setting the gas to a single speed throughout the mesh. As the simulation proceeds, the outer boundary of the \( z \)-axis is updated with the incoming gas at the stellar velocity. This update depends on the nature of the problem and accounts for the change in the cloud density distribution as a function of position. In reality, the situation is far more complicated because the velocity is also a function of time: the stellar velocity changes as it moves through the gravitational potential of the cloud, but unfortunately, this variation cannot be included here in a self-consistent manner. All models have the same numerical resolution of 250 zones along both the \( x \) and \( z \) axes.

We use the same approach as Paper I to model the \( \text{H} \) \( \Pi \) region (see Bodenheimer et al. 1979), i.e., we solve a radial integral to find the position of the ionization front (Strömgren 1939). The temperature inside the \( \text{H} \) \( \Pi \) region is set to \( 10^4 \text{ K} \), the approximate photoionization equilibrium temperature.

The modeling of a “realistic” UC \( \text{H} \) \( \Pi \) region includes a hypersonic stellar wind (i.e., an ultracompact bubble) and is very expensive in computational time, owing to the fact that the wind speeds from massive main sequence stars are of order \( 10^3 \text{ km s}^{-1} \) when their bubble sizes are of order \( 10^{-2} \text{ pc} \). This forces the Courant condition to calculate very small time steps during the simulations. Thus, it is possible to include stellar winds in two-dimensional simulations that cover physical times of order \( 10^3 \text{ yr} \), but this is very time consuming when the required physical times are of order \( 10^5 \text{ yr} \). Arthur & Hoare (2006) included stellar winds in their simulations, but restricted the evolution to about \( 4 \times 10^4 \text{ yr} \). Hence, the long computational time required and the spatial resolution needed to resolve the swept-up wind shell. Although we do not include the effects of the stellar wind, Paper I showed

![Gas density snapshots of models A0 (left) and B0 (right). The star is located at the edge of the core, 0.1 pc from the core center: \( (x, z) = (0.2, 0.1) \).](image-url)
that at these high densities the size of the UC H II region is primarily determined by the gas pressure. Thus, our results can be considered qualitatively correct.

3.2. Stationary Stars

To begin our study we computed several cases (not shown here) in which the star was located at the center of the core (given by eq. [8]) and had zero velocity. For a core radius of 0.1 pc, a central density of 10^6 cm^{-3} and an ionizing photon flux of F_\nu = 10^{48} s^{-1}, the models followed similar tracks as those described in the one-dimensional solutions of Paper I. This is understandable, because the Gaussian core described by equation (8) produces a central plateau which resembles the constant density medium used in the first part of Paper I. Thus, pressure equilibrium is achieved on timescales matching those of Paper I. The novel aspect here is the inclusion of gravity, which affects the final density structure. This is illustrated in the next paragraph.

A second set of models, also without stellar motion (models A0 and B0; see Table 1), are calculated with the star at the edge of the core (Fig. 1). The purpose of these models is to study the expansion and dynamics of the H II regions in the two different density ramps given by equations (6) and (8) for r \geq r_c. Model A0 is shown in detail in Figure 2, while model B0 is shown in Figure 3. The ionized gas has a flat density distribution during the first forty thousand years (first two curves) in both cases. But as the evolution proceeds, the ionized gas—subject to the cloud gravity—adjusts its density distribution to come into hydrostatic equilibrium. This is clearly seen in model A0, where pressure equilibrium is achieved (curves 7–10) and the final density distribution follows the initial density ramp. The final density is a factor of 200 lower, however, because the gas was originally at 10^5 K but hydrostatic equilibrium is achieved at 10^4 K. In contrast, model B0 does not reach pressure equilibrium in the computational domain, and a blister-like region or champagne flow (Tenorio-Tagle 1979) is produced. Gas is permanently photoevaporated from the cloud, but the final stellar position is 0.65 pc from the center. This motion is represented in the model by the molecular core moving leftward, out of the figure. Log mass densities of -17, -17.5, and -18 g cm^{-3} correspond to number densities of 3 x 10^6, 9.5 x 10^5, and 3 x 10^4 cm^{-3}, respectively, for the molecular gas.
resulting in a steady state mass loss. As a result, the final solution is closer to an $r^{-2}$ profile than to the original $r^{-3}$ density ramp.

3.3. Moving Stars

A third set of models (A2, A8, A12, B2, B8, and B12) include stellar motion at 2, 8, and 12 km s$^{-1}$ (Table 1). In these cases the star moves from the core center toward the edge, eventually leaving the core and entering the density ramp described in equations (6) and (8). While the star is still inside the core, these models create cometary shapes, with sizes typical of UC H II regions. When the star enters the density ramp, the H II region grows in size along the density gradient and drops in density. If the stellar velocity is high compared to the H II expansion velocity, the cometary shape is maintained. If the stellar velocity is low compared to the H II expansion velocity, the region quickly evolves into a blister morphology. The results of this third set of models, including a constant stellar motion, are shown graphically in Figures 4–9. We note that the density ramps for these models are off-center from the stellar position; i.e., the origin of the power law does not coincide with the stellar coordinates. As such, the spherically symmetric solutions given by Franco et al. (1990) are not directly applicable in all cases.

Figures 4 and 5 show models A2 and B2, respectively. Owing to the lower stellar velocity, four time frames are adequate to show the evolution. In these two simulations the star is initially at the core center and moves toward $+z$. Our computational method leaves the star fixed at the position $[(x, z) = (0.25, 0.1)]$, while the gas moves toward $-z$ at 2 km s$^{-1}$. At the time of the first frame (0.8 x 10$^5$ yr) the star has already reached the core edge and entered into the density ramp. In subsequent frames, the star has moved outward in the density ramp, leaving the cloud progressively further behind (toward $-z$). The behavior of models A2 and B2 is qualitatively similar. However, owing to the steeper density gradient, model B2 grows more rapidly, showing greater expansion.

Figure 6 shows the final shape of the ionized region of models A0, B0, A2, and B2 (the final times shown for each model are different, as described in the figure legend). All four models present a cometary morphology, but in all cases the gas dynamics correspond to a champagne flow, not a bow shock. The sizes and shapes of the ionized regions are different (except for the B0 case) than those of the final cavities shown in Figures 1, 4, and 5. This is because the density of the ionized region tends to a uniform value, washing away the initial gradient, and as a consequence the ionization front recedes (i.e., it becomes a recombination front). The gas in the outer part of the expanding cavity cools and recedes. The pressure drops quickly in these regions, forming bubbles and filaments due to thermal and dynamical instabilities. The effects are clearly noticeable in the B2 model (last frames of Figs. 5 and 6).

Figures 7 and 8 show the total (Fig. 7) and ionized (Fig. 8) gas densities for models A8 and A12. A cometary morphology is clearly evident during the first half of the evolution for both models.

During the second half of the evolution, as the star moves into much lower density gas, the cometary arc opens substantially. Also, instabilities in both the ionization and shock fronts, as those described in Paper I and Williams (2002), create complex structures
with finger-like condensations that resemble elephant trunks (see also Williams et al. 2001). Indeed, it is not clear that at late times these ionized regions would even be recognized as being cometary. Corresponding time frames between the A8 and A12 models show that the H II region has evolved to a larger size for the A12 model than the A8 model. This is because the higher stellar velocity of the A12 model has carried the star to a lower density region of the cloud, where the H II region can expand more freely.

Figures 9 and 10 show the total (Fig. 9) and ionized (Fig. 10) gas densities for models B8 and B12. Compared to the A8 and A12 models, the B8 and B12 models evolve more rapidly, expanding beyond the ultracompact stage in about half the time (the time sequence shown in Figs. 9 and 10 is shorter than that of Figs. 7 and 8). As stated in the case of Figure 6, the ionization front is always inside the expansion cavity, and the gas recombines at several locations. Perhaps most intriguing about the B8 and B12 models is that both pass through a phase (at 52,000 yr and 40,000 yr, respectively) when distinct high-density components are seen within the H II region. These high-density components have sizes and densities typical of UC H II regions, but they are embedded in larger, lower density ionized gas more typical of compact H II regions. This is discussed further in § 4.6.

The sequence UC H II → compact H II → extended H II shows, as stated above, a rich variety of unpredictable structures due to
the ionization-shock front instability (see Paper I). Figures 7–11 show examples of how rich the development of the ionization-shock front instability can be. The small-scale structures found in the computations (Fig. 11) are always transient. The lifetime of these structures, or elephant trunks, is proportional to their mass and the size of the $\text{H}\alpha$ region, and inversely proportional to the stellar speed. This is because their destruction depends on the photoevaporation rate, which in turn depends on the position and orientation of the ionizing source. The model does not purport to explain the M16 pillars, which are 50 times larger, but does suggest, as discussed in more detail by Williams et al. (2001), how their formation might occur.

4. DISCUSSION AND CONCLUSIONS

The models presented here extend previous work (e.g., Franco et al. 1990; Paper I; Arthur & Hoare 2006) by the adoption of an off-center position for the star in the density gradient, the inclusion of the core gravity, and stellar motion both up and down the density gradient. As mentioned in § 3, we have not included a stellar wind in the calculations and so all the photoionized regions presented in this paper are filled, rather than confined to a swept-up wind shell (see Arthur & Hoare (2006) for examples including stellar winds).

The results presented here show a number of important features that may be useful in understanding some of the observational puzzles of $\text{H}\alpha$ region evolution. Our simulations are performed with simplified models, however, and further studies with additional refinements are needed. In particular, the gravitational acceleration is treated here as a time-independent radial function, and the stellar velocity is simply set to a constant value. Obviously, both quantities are continually modified as the cloud is photoevaporated and the star changes location in the gravitational field.
Thus, a self-consistent treatment of the problem will provide better details of the evolutionary features. Also, we do not include preexisting clumps in the cloud matter, nor large-scale vorticity in the gas velocity, which will add complexity to the resulting structures. Despite these restrictions, our results provide an adequate qualitative guideline to the effects resulting from the motion of recently formed massive stars in centrally condensed cloud cores.

4.1. Bow Shock Models and Champagne Flows

Cometary H II regions are generally classified as “champagne flows” or “bow shocks,” although observations often show attributes corresponding to a mixture of the two (Lumsden & Hoare 1999; Garay & Lizano 1999). The bow shock model (e.g., van Buren et al. 1990; van Buren & Mac Low 1992) consists of a star moving at supersonic speeds through a (uniform) medium. The star possesses a strong stellar wind, and the H II region is confined to the shell swept up by the stellar wind. A champagne flow, on the other hand, occurs when photoionized gas from the ionization front in a medium with a strong density gradient accelerates in the direction of decreasing density. The bow shock model can be distinguished from the champagne model both morphologically and kinematically. The bow shock model will have a limb-brightened morphology in radio-continuum images, while...
a champagne flow will have a more centrally condensed morphology. Kinematically, the bow shock model has the highest ionized gas velocities at the head of the cometary shape, and these velocities are roughly equal to the stellar velocity. In the tail, the velocities are lower. For a champagne flow, the highest velocities are in the tail region, because the gas evaporating from the ionization front accelerates down the density ramp, owing to the strong pressure gradient.

By not including a stellar wind in the present calculations, we force the H\textsc{ii} regions generated to be of the “champagne flow” type. Arthur & Hoare (2006) have recently made detailed numerical simulations of champagne flows, both with and without stellar winds, and also bow shock models where the star moves in a uniform medium and up a density gradient. The effect of the stellar wind is to confine the champagne flow to an open shell around the stellar wind bubble, thereby giving a limb-brightened morphology while preserving the overall kinematics of a champagne flow model, i.e., the highest velocities in the tail region. The present work uses a slab geometry, so a direct comparison with observations (such as that done by Arthur & Hoare 2006) is not possible. Although (supersonic) stellar motion is included in the present models, the structures produced are not bow shocks in the van Buren & Mac Low sense. Because no fast stellar wind is included, the only shocks are those in the neutral gas ahead of the ionization front. These shocks are not pressure driven in the same way as a stellar wind bubble and are not properly described as bow shocks. Because the H\textsc{ii} region is not swept up into a thin shell ahead of the star, the expected kinematics from our stellar motion models

Fig. 10.—Photoionized gas density counterpart of Fig. 9. Frames 4 (bottom of first column) and 11 (third frame of third column) illustrate the effect of the star overtaking a leading, piled-up clump of gas. This could result in either a core-halo morphology or an UC H\textsc{ii} region with extended emission. See § 4.5 and Fig. 12.
is more similar to a champagne flow model; i.e., the photo-evaporating gas accelerates away from the ionization front in the dense medium and down the pressure gradient into the less dense medium. Note that the radio-continuum morphologies of the H II regions presented in Figures 8 and 10 do not show limb-brightening, as would occur for bow shocks.

The fact that the star is supersonic with respect to the ionized gas (models A12 and B12) is not particularly important for the H II region dynamics—it merely results in an elongated (filled) H II region such as those described by Raga (1986) for stellar motion in a uniform medium. The internal motions of the H II region could be marginally affected since the internal pressure and density gradients dissipate on a sound crossing time. However, for models A12 and B12, the fact that the star is moving faster than both the ionized gas sound speed and the ionization front, means that in the densest regions it can overtake the main ionization front and form a new, smaller, elongated H II region around itself.

4.2. Density Gradients

Franco et al. (1990) described the formation and evolution of H II regions in spherically symmetric power-law density distributions. Their main result was that if the initial Strömgren radius
is larger than the radius of the central, uniform core, then for power-law indices steeper than $-3/2$ the whole cloud is immediately photoionized; i.e., the ionization front is at infinity. In the present paper, power-law indices of $-2$ and $-3$ are used and yet the initial Strömgren region is confined—even when the star is located at the core edge, as in models A0 and B0. Moreover, during the expansion stage, only for the power law $-3$ case does the ionization front break out and overtake the preceding shock wave. At first sight this appears to contradict Franco et al. (1990). In fact, the nature of the problem is changed by the star and $H\,\Pi$ region being off-center from the power-law density distribution.

The density distribution is given by

$$n = n_c \left( \frac{r}{r_c} \right)^{-\alpha}, \quad (11)$$

with the star positioned at $r_c$. In a reference system centered on the star, we can rewrite the density distribution using polar coordinates $(R, \theta)$ as

$$r^2 = R^2 + r_c^2 - 2r_c R \cos (\pi - \theta), \quad (12)$$

where $\theta$ is measured from the symmetry axis and the core center is located at $\theta = \pi$. The density distribution of equation (11) then becomes

$$n(R, \theta) = n_c \left( \frac{R^2 + r_c^2 + 2r_c R \cos \theta}{r_c^2} \right)^{-\alpha/2}. \quad (13)$$

Along the symmetry axis, $\theta = 0$, this is

$$n(y) = n_c (1 + y)^{-\alpha}, \quad (14)$$

where we have defined $y = R/r_c$.

For $y < 1$ we can write $n(y) \approx n_c (1 - \alpha y)$, which is no longer a power law. For $y > 1$ we find $n(y) \approx n_c y^{-\alpha}$; that is, the offset density distribution resembles the original power-law density distribution. Consequently, if the initial Strömgren radius $R_\text{S}$ is smaller than the core radius of the initial density distribution, then we expect the initial $H\,\Pi$ region to remain confined, even for power laws steeper than $\alpha > 3/2$ and for stars located at the core edge.

In fact, for $\alpha > 3/2$, and considering ionization balance along the $\theta = 0$ direction, we obtain the general criterion

$$\frac{1}{3} y_{sc}^3 < \frac{-2}{(1-2\alpha)(2-2\alpha)(3-2\alpha)} \quad (15)$$

for the $H\,\Pi$ region to remain confined. Here $y_{sc} = R_\text{S}/r_c$ is the nondimensional Strömgren radius for a uniform medium of number density $n_c$ (S. J. Arthur 2007, in preparation). For example, when $\alpha = 2$, we require $y_{sc} < 1$, while for $\alpha = 3$, the condition is $y_{sc} < 0.1$.

In the models of this paper, $r_c = 0.1$ pc, the density $n_c = 10^{7}/\alpha^{2}, Q_{H\Pi} = 10^{48}$ s$^{-1}$, and the nondimensional Strömgren radius is thus $y_{sc} = 9.5 \times 10^{-3}$ for the $\alpha = 2$ case and $y_{sc} = 1.1 \times 10^{-2}$ for the $\alpha = 3$ case. Thus, in both cases the initial Strömgren region will remain compact.

During the expansion stage, the $H\,\Pi$ region will remain bounded so long as the ionization front does not accelerate to supersonic velocities before pressure balance is reached with the ambient medium. Once the ionization front is supersonic with respect to the ionized gas it quickly races off. This occurs first for the $\theta = 0$ direction where the density gradient is steepest in our off-center reference frame. This problem is studied in detail in S. J. Arthur (2007, in preparation). Initially, for small $y_{sc}$, the expansion is similar to the classical (Spitzer 1954) expansion because the $H\,\Pi$ region does not feel the density gradient until its radius is greater than $r_c$, hence the ionization front velocity is decreasing. Once the density gradient becomes important, the ionization front will accelerate. For model A0 of this paper ($\alpha = 2$), the ionization front velocity does not become supersonic before pressure balance is reached, while in model B0 ($\alpha = 3$) it does—and as a consequence the $H\,\Pi$ region quickly ionizes out to the edge of the grid in the directions where this condition is true.

The results for model A0 are also interesting because if pressure equilibrium is achieved, the density gradient in the $H\,\Pi$ region will match the original density gradient of the parental cloud (see Franco et al. 2000). However, model B0 shows that molecular cloud density gradients above $\rho \propto r^{-2}$ cannot lead to pressure equilibrium, indicating that $H\,\Pi$ region gradients obtained in these cases probably do not reflect the initial cloud conditions.

Another aspect that is important to stress is that champagne flows are transient, and in some cases they may be short-lived. Figures 1–6 show that, except for model B0, the initial champagne flow disappears either because the $H\,\Pi$ region reaches pressure equilibrium or the ionized gas recombines. The latter case leads to neutral outflows generated by a recombination front, as originally discussed for disklike clouds by Franco et al. (1989). Three scenarios give rise to more-or-less steady champagne flow solutions: density gradients steeper than $\rho \propto r^{-2}$ at the core edge; ionizing photon rates high enough to ionize the whole core; and core densities much smaller than the value used here ($n_0 = 10^7$ cm$^{-3}$). The latter possibility appears to be ruled out for hot molecular cores (which have densities equal to or even larger than $10^7$ cm$^{-3}$).

4.3. Instabilities

Evidence of instabilities forming at the ionization front is very clear in Figures 7–10, for models A8, A12, B8, and B12, with stellar velocities of 8 and 12 km s$^{-1}$. For these models, the shape of the photoionized region is very elongated owing to the high stellar speed. Shear flows between the shocked neutral gas and the photoionized gas just inside the ionization front generate Kelvin-Helmholtz instabilities. The eddies that form create enhancements in the gas density on both sides of the ionization front. These small density enhancements cause a shadowing instability to form at the ionization front. This instability has been extensively discussed by Williams (1999) and was seen in the simulations of Arthur & Hoare (2006). The density fluctuations cause variations in the optical depth, which leads to the shadowing instability.

At later times in the evolution of models A8, A12, B8, and B12, the distinctive spokes of ionized gas characteristic of this instability are seen. These spokes represent the propagation of $R$-type ionization fronts in collimated beams through regions of lower density at the ionization front. The bases of the spokes are generally wider than the tips because the ionized gas at the base overpressurizes first and the gas can start to expand laterally. The shadowing instability can lead to very dense neutral clumps forming at the base of the ionized spokes, and, indeed, their density continues to increase with time as more gas is channeled into the clumps (Williams 1999). This is very evident in the final panels of Figures 7–10.

4.4. The Star Formation Rate

The results for model A0 indicate that solutions exist in which clouds are not completely disrupted by ionization fronts; hence, star formation could proceed for longer times in those clouds.
Because stellar motion is able to confine H\textsc{ii} regions by ram pressure, the ionized gas is confined within the cloud and the star formation rate of the cloud is relatively unaffected. The results are quite different for model B0, with a steep gradient, or any of the models with larger stellar velocities, (8 km s\(^{-1}\) or above) in which the rapid expansion of the ionized gas in the lower density regions of the cloud would quickly lead to cloud destruction, thus ending the star formation phase (Franco et al. 1994; Díaz-Miller et al. 1998).

4.5. The Lifetime Problem and UC H\textsc{ii} Morphologies

Any proposed solution to the lifetime problem for UC H\textsc{ii} regions must explain how dense, compact ionized gas can remain in that state for timescales of order 10\(^5\) yr. We find that model A0 (the stationary model with \(\rho \propto r^{-2}\), and also the stationary one-dimensional model of Paper I) and models A2 and B2, (the low stellar velocity models) are able to produce UC H\textsc{ii} regions smaller than 0.1 pc at evolutionary times larger than 10\(^5\) yr. UV absorption by dust grains inside the cloud cores substantially reduces the size of H\textsc{ii} regions in dense places (Díaz-Miller et al. 1998; Arthur et al. 2004). We do not include the effects of dust, hence our results provide only upper limits for the sizes of the photoionized regions. This makes even easier the confinement of UC H\textsc{ii} regions during the required timescales.

The stationary models remain compact because the H\textsc{ii} region reaches pressure equilibrium with the high-density surroundings. In the case of model A0, the H\textsc{ii} region is at the edge of the exponential core; for the one-dimensional model of Paper I, the H\textsc{ii} region is in a uniform density gradient. UC H\textsc{ii} regions remain ultracompact (\(<0.1\) pc) as long as they remain inside the cloud core. In this case, the lifetime as an ultracompact region would be defined by the core-crossing time. In the event that the stellar orbit never reaches the core edge, the ultracompact lifetime would be the stellar lifetime, as discussed in Paper I.

A star moving at 1 km s\(^{-1}\) will leave a 0.1 pc core in 10\(^5\) yr (starting from the core center) and enter the density ramp. The reason why solutions of models A2 and B2, with 2 km s\(^{-1}\) stellar velocities, are still within the range of UC H\textsc{ii} parameters is because of ram pressure confinement.

Wood & Churchwell (1989) found that the cometary and core-halo morphologies were the most common (20% and 17%, respectively, in their sample), suggesting that these morphologies have longer lifetimes than the others, thus making their detection more probable. For comparison, we find the evolutionary time for each model when the solution (still dominated by ram pressure) reaches a size of 0.1 pc. These times (see Table 1) are A2–B2 \(\sim 8 \times 10^4\) yr, A8–A12 \(\sim 6.4 \times 10^4\) yr, and B8–B12 \(\sim 4 \times 10^4\) yr. The times for the higher velocity models are roughly one-third of the nominal 10\(^5\) yr UC H\textsc{ii} lifetime, consistent with the 37% of the Wood & Churchwell sample showing a cometary or core-halo morphology.

As mentioned above, preexisting, neighboring clumps within the cloud are not included in our model. This excludes the possibility of a star entering a molecular core after it has previously exited its natal core. This might occur if a nearby molecular core exists, and if the stellar velocity is sufficient to reach the core within the stellar lifetime. Such an event might occur even for relatively low velocity stars, because core separations could be of the order of tenths of parsecs—easily attainable during a stellar lifetime, even at 1 km s\(^{-1}\). Such a scenario would not be a mere time-reversal of the present simulations. Rather, the stellar velocity would move the star up the density gradient, leading to distinct dynamical effects. Although the details of such an occurrence require further modeling, the general result would be the eventual recombination of the lower density gas, and the choking off of the H\textsc{ii} region as the star moves into denser gas.

4.6. The Coexistence of Compact and Extended Emission

An intriguing self-blocking effect, in which the star overtakes, piled-up clump of gas and ionizes it, is seen in models B8 and B12. The effect occurs at 52,000 and 40,000 yr, respectively, shown in frames 4 and 11 of Figure 10. Such regions could appear in radio maps as having a core-halo morphology. Moreover, if the linear extent of the ionized gas is several tenths of a parsec or more when the self-blocking occurs, this phenomenon could appear as an UC H\textsc{ii} region coexisting with extended emission. The H\textsc{ii} density contrast in frames 4 and 11 is about 20. For comparison, inferred values for the density contrast in the Orion H\textsc{ii} region range from about 40 (Wilson & Jäger 1987) to nearly 600 (Felli et al. 1993), with the larger value resulting from sensitivity to small, dense clumps.

The self-blocking transients seen in models B8 and B12 depend on the resolution used; a resolution convergence study is needed to confirm the effect. This effect is not seen in models A8 or A12. This could be a result of the shallower density gradient or it could also be a result of the resolution employed. We caution that a more complete treatment of the radiative transfer could modify these results, hence we do not make any quantitative claim about the frequency of occurrence or lifetime of this morphology.
We merely suggest that the core-halo or extended emission morphology might arise from self-blocking effects. This is illustrated in Figure 12, where we show a contour plot of $n_e^2$ for model B12 together with a contour image of G31.3–0.2, a UC H II region with extended emission. The model does not purport to explain the G31.3–0.2 region, which is about 30 times larger, but it does suggest how its formation might occur.

4.7. Summary

To summarize, we have modeled the evolution of H II regions in molecular clouds in hydrostatic equilibrium and including the effects of cloud gravity. The cloud structure was modeled as Gaussian central cores and external halos that obey $\rho \propto r^{-2}$ and $\rho \propto r^{-3}$ power laws. The models indicate that UC H II regions are long-lived objects while they are inside their parental cloud cores. Provided that the stellar velocity is not too great, H II regions with $\rho \propto r^{-2}$ density gradients can be confined for long periods, because of their off-center position. When stars escape from these cores, transitions to larger structures, possibly UC H II regions with extended emission, are expected to occur.

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