THE ASTROPHYSICAL JOURNAL, 427:384–387, 1994 May 20

A Q CONDITION FOR LONG-RANGE PROPAGATING STAR FORMATION

B. G. ELMEGREEN
IBM Research Division, T. J. Watson Research Center, P.O. Box 218, Yorktown Heights, NY 10598
Received 1993 September 13; accepted 1993 November 17

ABSTRACT

Collapse conditions for large expanding shells and rings in a disk galaxy are derived and shown to lead to a condition for star formation that is similar to the condition $Q < 1$ for spontaneous instabilities. This result implies that both spontaneous and stimulated star formation are sensitive to Q, and that the observation of a critical surface density for star formation that is based on Q does not necessarily imply that star formation results from large-scale quiescent instabilities. The results also suggest that in regions with high gas densities and high rotation rates, such as starburst galaxy nuclei, the normal balance between stimulated and spontaneous star formation mechanisms could shift to give a higher proportion of stars forming in shells and other swept-up debris, and less in giant cloud complexes containing the local Jeans mass.

Subject heading: stars: formation

1. INTRODUCTION

Star formation is often observed along the periphery of giant expanding shells or rings. One of the best-known examples is Constellation III in the Large Magellanic Clouds (Westerlund & Mathewson 1966). The Lindblad ring could be another example with the Orion, Perseus, and Sco-Cen molecular clouds and OB associations on the periphery (Elmegreen 1982; Olano 1982). Many other examples are reviewed in Elmegreen (1985a, 1987), Tenorio-Tagle & Bodenheimer (1988), and Elmegreen (1992). Such star formation is presumably initiated by the formation of new molecular clouds in the swept-up matter along the shell or ring. The clouds form as the material expands because of gravitational instabilities in the compressed gas. The first models of this type were by Ögelman & Maran (1976), Tenorio-Tagle (1981), Elmegreen (1985b), Tenorio-Tagle & Palouš (1987), McCray & Kafatos (1987), and Franco et al. (1988). A recent model is by Comeron & Torra (1994).

Here we derive approximate conditions for the collapse of expanding and decelerating shells and rings. The results illustrate the basic time and length scales for the process, and suggest some constraints on whether collapse will or will not occur. In particular, we suggest that on a large scale this mechanism of triggering should operate in an interstellar medium where the stability parameter Q for the gas is close to or less than the threshold value of 1 if the shells or rings are relatively thick.

2. DERIVATION

2.1. *Collapse in Decelerating, Expanding Shells*

Consider a three-dimensional shell expanding into a uniform medium. The unperturbed mass column density in the shell is σ_0, which is assumed to increase with radius as $(4\pi R^3 \rho_0/3)/(4\pi R^2) = \rho_0 R/3$ by mass conservation with a pre-shell density ρ_0, and the shell expansion speed and radius are V and R, which are assumed to vary as $V \propto t^{-0.4}$ and $R \propto t^{0.6}$ for continuous energy deposition into the cavity (Castor, McCray, & Weaver 1975).

The shell will in general contain perturbations in the column density and transverse velocity that may lead to instabilities of various types. The perturbed column density will be denoted by σ_1 and the transverse (perturbed) velocity in the shell is $v = R\,\partial\theta/\partial t$ for a small region with angular coordinate θ. We consider the angular velocity $\partial\theta/\partial t = \Omega$ of part of the shell, rather than the translational velocity $v = R\Omega$ because R changes with time and the collapse of the shell along the periphery is a collapse in angular coordinates. We also assume that the perturbed pressure times thickness in the shell is $\sigma_1 c^2$ for constant velocity dispersion c in the shell.

The perturbed equation of motion for transverse flows in the shell is

$$\sigma_0 R \frac{\partial \Omega}{\partial t} = -c^2 \nabla \sigma_1 + \sigma_0 g_1 - 4\sigma_0 \Omega V , \quad (1)$$

where the perturbed gravitational acceleration satisfies Poisson's law,

$$\nabla \cdot g_1 = -4\pi G \rho_{\rm shell} . \quad (2)$$

The last term in equation (1) is from the expansion of the shell and the accretion of new mass. It comes from the total inertial force on the perturbed region, which is the rate of change of total transverse momentum for perturbed mass M_1: $d(M_1 v)/dt = d(M_1 R\Omega)/dt = (dM_1/dt)R\Omega + M_1 V\Omega + M_1 R(d\Omega/dt)$. We set $M_1 = 4\pi\alpha\rho_0 R^3/3$ for fractional solid angle of the perturbation α, so $dM_1/dt = 4\pi\alpha\rho_0 R^2 V$, and then divide all of these terms by the instantaneous perturbation area $A_1 = 4\pi\alpha R^2$ with $\sigma_0 = \rho_0 R/3$. The last term in equation (1) is then a combination of the term $M_1 V\Omega/A_1 = \sigma_0 V\Omega$ from the expansion of the shell and $(dM_1/dt)R\Omega/A_1 = 3\sigma_0 V\Omega$ from the accretion of new matter. These are the same as the first two terms in equation (2.15c) in Vishniac (1983).

An additional term in Vishniac's equation (2.15c) is from the action of the shock-driving pressure on perturbations in the trailing surface of the shell. This term leads to transverse motions in the shell on the dynamical timescale R/V, but these motions are apparently not disruptive as suggested by Vishniac (1983), nor should they prevent the eventual gravitational instability on the longer timescale $(G\rho_{\rm shell}\mathscr{M})^{-1/2}$ for $\mathscr{M} = V/c$. The transverse flows appear to be regular and confined to the shell in the nonlinear regime (Mac Low & Norman 1993). They probably affect the gravitational instability by adding kinetic energy to the collapsing gas (Vishniac 1994). At the same time, some of the energy in the shell will be lost by various nonlinear dissipation mechanisms related to turbulence and motions in the shock front. To simplify all of these matters, we assume that the total internal kinetic energy in the swept-up gas can be represented by the velocity dispersion c, which enters into the pressure term. The swept-up magnetic field in the shell can also be represented to some extent by an additional pressure contributing to c. The value of c is uncertain without a detailed calculation of the sources and sinks of energy and magnetic pressure in the shell. Nevertheless, we parameterize our results with a dimensionless quantity $\mathscr{M} = V/c$, which is a measure of the shell's compression or thickness ($\mathscr{M}$ is not the Mach number of the shock, which would be V divided by the sound speed ahead of the front). When $\mathscr{M}$ is large, the shell is thin and the collapse is rapid because of the high density. Strong dynamical instabilities in the shell tend to make $\mathscr{M}$ small if the corresponding energy does not dissipate, and this delays the collapse by making the shell thick and the average shell density low. More complete discussions of the dynamical instability are in Hunter et al. (1986), Vishniac & Ryu (1989), Nishi (1992), and Yoshida & Habe (1992).

The perturbed equation of mass continuity in the transverse direction is

$$\frac{\partial \sigma_1}{\partial t} = -\sigma_0 R \nabla_T \cdot \Omega - 2\sigma_1 \frac{V}{R}, \tag{3}$$

where the first term is from the convergence of the perturbed flow in the shell and the second term is from the stretching of the region of perturbed surface density with the expansion. This second term comes from the time derivative for the perturbed mass M_1 per unit area A_1 for constant M_1, that is, $\partial(M_1/4\pi\alpha R^2)/\partial t = -2(M_1/4\pi\alpha R^2)(V/R) = -2\sigma_1 V/R$.

Now we approximate the perturbation by sine functions in the small part of the shell under consideration. We write $\sigma_1 = \hat{\sigma}_1 \cos \eta\theta$ for angular coordinate θ along the shell and constant angular wavenumber $\eta = kR$; here k is the spatial wavenumber, equal to 2π divided by the wavelength (k varies as $1/R$). Then $\nabla\sigma_1 = -k\hat{\sigma}_1 \sin \eta\theta$. We also write the perturbation's angular velocity $\Omega = \hat{\Omega} \sin \eta\theta$. Finally, we write ω for the time derivatives in analogy with the exponential growth rates found in other instability problems, even though there is no real exponential growth in this problem because of the expansion and time dependence of σ_0. This latter approximation would be replaced, in a more exact solution, by a numerical integration of the equations over time (e.g., Yoshida & Habe 1992). (If gravity and internal pressure were ignored, the solution would be a power law, as found by Vishniac 1983).

With these approximations, and with $g_1 = -2\pi G\hat{\sigma}_1 \sin \eta\theta$, we can rewrite the equations of motion and continuity as

$$\sigma_0 R\omega\hat{\Omega} = c^2 k\hat{\sigma}_1 - 2\pi G\hat{\sigma}_1\sigma_0 - 4\sigma_0\hat{\Omega}V, \tag{4}$$

$$\omega\hat{\sigma}_1 = -\sigma_0\eta\hat{\Omega} - 2\hat{\sigma}_1 V/R, \tag{5}$$

and then eliminate $\hat{\Omega}$ and $\hat{\sigma}_1$ to get the instantaneous growth rate ω as a function of the angular wavenumber η:

$$\omega = -\frac{3V}{R} + \left(\frac{V^2}{R^2} + \frac{2\pi G\rho_0 \eta}{3} - \frac{c^2\eta^2}{R^2}\right)^{1/2}. \tag{6}$$

The peak growth rate occurs at the wavenumber where $d\omega/d\eta = 0$, which is

$$\eta_{\text{peak}} = \frac{\pi G\rho_0 R^2}{3c^2}, \tag{7}$$

and the peak growth rate is

$$\omega_{\text{peak}} = -\frac{3V}{R} + \left[\frac{V^2}{R^2} + \left(\frac{\pi G\rho_0 R}{3c}\right)^2\right]^{1/2}. \tag{8}$$

The peak growth rate is smaller than what it would have been for a nondivergent shell because of the dilution of the density perturbation with accretion and the stretching of the density perturbation with expansion. These effects stabilize the shell when R is small (Ostriker & Cowie 1981). Instability occurs only for $\omega_{\text{peak}} > 0$, which requires

$$\frac{\pi G\rho_0}{3c} > \frac{8^{1/2}V}{R^2}. \tag{9}$$

Writing $V/R = 0.6/t$, this constraint becomes

$$t > \frac{0.986}{(G\rho_0 \mathscr{M})^{1/2}}, \tag{10}$$

where $\mathscr{M} = V/c$ and c is the rms dispersion in the shell. For an adiabatic shock, $\mathscr{M} = (16/5)^{1/2} = 1.8$, which gives a long collapse time and a thick shell with large R at the time of collapse. Approximately the same value of $\mathscr{M}$ applies to cloud agglomeration fronts that have traveled less than or equal to several cloud mean free paths (Elmegreen 1988), or to shocks in which dynamical instabilities are active and clumps form with an rms speed of around the shock speed, or magnetic shocks in which the magnetic pressure dominates the thermal pressure, or very old shells for which V is small. In all of these cases, the shell is thick and the density relatively low, so the collapse takes a long time. If most of the shock energy dissipates and the shell is thin, then $\mathscr{M}$ can be large and the collapse rapid. The value of $\mathscr{M}$ for a particular region can be inferred from observations of the relative shell thickness ΔR because $\mathscr{M} = (3\Delta R/R)^{-1/2}$ for a shell-to-preshell density ratio equal to $\mathscr{M}^2$.

There is another constraint for the calculation to be valid and that is that the minimum wavelength for instabilities must fit inside a fraction of the shell circumference. The minimum wavelength comes from the equation $\omega(\eta) > 0$, which gives

$$\eta < \frac{\pi G\rho_0 R^2}{3c^2}[1 + (1 - \xi^2)^{1/2}], \tag{11}$$

where the dimensionless parameter

$$\xi = \frac{8^{1/2}V/R}{\pi G\rho_0 R/3c} \tag{12}$$

was implicitly used for the condition that $\omega_{\text{peak}} > 0$, i.e., $\xi < 1$, as discussed above. For the minimum wavelength less than the shell radius, we write $2\pi < \eta$, so we have

$$2\pi < \frac{\pi G\rho_0 R^2}{3c^2}[1 + (1 - \xi^2)^{1/2}], \tag{13}$$

which is the same as

$$\frac{\pi\xi}{2^{1/2}\mathscr{M}} < 1 + (1 - \xi^2)^{1/2} \tag{14}$$

or $\xi < 0.428$, 0.748, and 0.994 for $\mathscr{M} = 0.5$, 1, and 2. Larger $\mathscr{M}$ makes equation (14) satisfied automatically when $\omega_{\text{peak}} > 0$; because $\mathscr{M}$ has to be fairly large anyway for the approximations involving a shell geometry to be valid, this size constraint, $2\pi < \eta$, is not likely to be important.

The collapse begins when $\omega_{peak} > 0$, but the collapse becomes well developed so that clouds form only at a later time, defined by $t \approx 1/\omega_{peak}(t)$. To evaluate this condition, we write ω_{peak} from above in terms of ξ:

$$\omega_{peak} = \frac{V}{R}\left[-3 + \left(1 + \frac{8}{\xi^2}\right)^{1/2}\right] \quad (15)$$

and set this equal to $1/t = (5/3)(V/R)$. Then we find that $\xi = 0.62$ when $t = 1/\omega(t)$, independent of $\mathcal{M}$. This gives a time and distance for significant collapse and cloud formation,

$$t_{clouds\ form} = \frac{1.25}{(G\rho_0 \mathcal{M})^{1/2}} = 103\left(\frac{n_0 \mathcal{M}}{cm^{-3}}\right)^{-1/2} \text{Myr}, \quad (16)$$

$$R_{clouds\ form} = 176\mathcal{M}^{1/2}\left(\frac{c}{km\ s^{-1}}\right)\left(\frac{n_0}{cm^{-3}}\right)^{-1/2} \text{pc}. \quad (17)$$

The angular wavenumber of peak growth at this time, $\pi G\rho_0 R^2/3c^2$, gives the wavelength relation

$$\frac{1}{\eta} = \frac{\lambda}{2\pi R} = \frac{0.354\xi}{\mathcal{M}} = \frac{0.22c}{V} \quad (18)$$

so that $\sim(4V/c)^2$ big condensations form in the shell; recall that V is the instantaneous expansion speed and c^2 is the average ratio of total pressure to density in the swept-up-gas.

2.2. *Collapse of Expanding Rings*

The equations are very similar for the collapse of a ring in the disk of a galaxy. The scenario we have in mind now is the expansion of a giant shell in which most of the accumulated material originally in the interior is forced to move outward. Some of this material may go into the halo, but most of it will stay in the galactic plane, expanding away from the pressure source as a ring. We imagine that this ring has a radius R, half-thickness r, and half-height equal to about r also. Thus, the shape of a small perturbation is approximately a curved cylinder. Then the gravity term in the dispersion relation considered above, $2\pi G\sigma_1 k = 2\pi G\rho_0 \eta/3$, which was for part of a shell, should be replaced by $4\pi G\rho_1(1 - krK_1[kr]) \sim 2G\mu_1 k^2 \ln(2/[kr])$ for Bessel function K_1; the latter approximation is for $kr \ll 2$ with $k = \eta/R$. The mass per unit length in the ring, $\mu_0 = \rho_0 Rr$, replaces the mass per unit area, σ_0, in the shell. Also for a ring, the coefficients of the terms in the equations of motion and continuity that result from the accretion and expansion change because of the different geometry, giving for these equations,

$$\mu_0 R \frac{\partial\Omega}{\partial t} = -c^2\nabla\mu_1 + \mu_0 g_1 - 3\mu_0\Omega V, \quad (19)$$

$$\frac{\partial\mu_1}{\partial t} = -\mu_0 R\nabla_T \cdot \Omega - \mu_1\frac{V}{R}. \quad (20)$$

The dispersion relation for an expanding ring now becomes

$$\omega = -\frac{2V}{R} + \left[\frac{V^2}{R^2} + \frac{2G\mu_0\eta^2\ln(2R/\eta r)}{R^2} - \frac{c^2\eta^2}{R^2}\right]^{1/2}. \quad (21)$$

From this dispersion relation we can determine that the peak growth occurs at the wavenumber given by

$$\frac{\eta_{peak} r}{2R} = \exp\left[-0.5\left(1 + \frac{c^2}{GRr\rho_0}\right)\right], \quad (22)$$

and that the peak growth rate is

$$\omega_{peak} = -\frac{2V}{R} + \left(\frac{V^2}{R^2} + \frac{G\rho_0 r\eta_{peak}^2}{R}\right)^{1/2}. \quad (23)$$

Now we set this peak growth rate equal to $1/t = (5/3)V/R$ and find the time for significant collapse; the result is an equation for t,

$$t \exp\left[-0.5\left(1 + \frac{c^2}{GRr\rho_0}\right)\right] = \frac{0.60}{(G\rho_0)^{1/2}\mathcal{M}}, \quad (24)$$

where we have used the condition for pressure balance, $\rho = \rho_0\mathcal{M}^2$, for ring density $\rho = \mu_0/(\pi r^2)$. This equation has to be solved numerically. We write for simplicity

$$t = \frac{T_r}{(G\rho_0)^{1/2}\mathcal{M}}, \quad (25)$$

and introduce the scale height

$$H^2 = \frac{c_0^2(1 + \alpha + \beta)}{2\pi G\rho_{0T}} \quad \text{for} \quad \alpha = \frac{B^2}{8\pi P} \quad \text{and} \quad \beta = \frac{P_{CR}}{P} \quad (26)$$

with ambient magnetic field strength B, cosmic ray pressure P_{CR}, turbulent pressure P, velocity dispersion c_0, and total gas + star density ρ_{0T} in the gas layer. We also write $\mu_0 = \pi R^2 2H\rho_0/(2\pi R) = R\rho_0 H$, or $r = H$. With these substitutions,

$$\frac{c^2}{GRr\rho_0} = \frac{0.6(2\pi)^{1/2}c/c_0}{T_r([1 + \alpha + \beta]\rho_0/\rho_{0T})^{1/2}}. \quad (27)$$

Typically, $(1 + \alpha + \beta)\rho_0/\rho_{0T} \sim 1$, so $c^2/(GRr\rho_0) \sim (1.5/T_r)(c/c_0)$.

With these substitutions, we find that the parameter T_r in the collapse time satisfies the equation

$$T_r \exp\left[-0.5\left(1 + \frac{1.5c}{T_r c_0}\right)\right] = 0.60, \quad (28)$$

from which we obtain $T_r = 1.26$, 1.59, 2.05, and 2.84 if $c/c_0 = 0.5$, 1, 2, and 4, respectively. Using a typical value of $T_r \approx 1.5$, the resulting timescale for significant collapse and molecular cloud formation is

$$t_{clouds\ form} \approx \frac{1.5}{(G\rho_0)^{1/2}\mathcal{M}} = 124\left(\frac{n_0\mathcal{M}^2}{cm^{-3}}\right)^{-1/2} \text{Myr}, \quad (29)$$

which is comparable to the final result for the shell except that now $\mathcal{M}$ appears in the denominator and before it was $\mathcal{M}^{1/2}$. The reason for the $1/\mathcal{M}$ dependence now is that ω_{peak}^2 scales with $G\rho_0 r\eta_{peak}^2/R$ from equation (23), but $\eta_{peak} \sim 2R/r$ from equation (22), and $R/r = \pi\mathcal{M}^2$ from the pressure relation $\rho/\rho_0 = \mathcal{M}^2$, so $G\rho_0 r\eta_{peak}^2/R \sim 4\pi G\rho_0\mathcal{M}^2$, and this is proportional to $1/t^2$ at the time of significant collapse.

This collapse time corresponds to a radius $R = 5tV/3$ which is independent of $\mathcal{M}$:

$$R_{clouds\ form} \approx \frac{2.5c}{(G\rho_0)^{1/2}} = 211\left(\frac{c}{km\ s^{-1}}\right)\left(\frac{n_0}{cm^{-3}}\right)^{-1/2} \text{pc}. \quad (30)$$

3. A Q CONDITION

In galaxies with rotation and shear, an expanding shell or ring will be distorted into an ellipsoid or ellipse by Coriolis forces. The time when a ring has its maximum extent on the minor axis, just before the Coriolis force begins to decrease the minor axis, is $\sim 2.5/\kappa$ for epicyclic frequency κ (Palouš, Franco, & Tenorio-Tagle 1990; Fig. 3*a*). In the solar neighborhood, $2.5/\kappa = 80$ Myr. This limiting time for Coriolis forces is a constraint on the collapse of the ring into giant molecular clouds. If the collapse time is much longer than $2.5/\kappa$, then the shell or ring will deform so severely before new giant molecular clouds form that the central cavity will begin to close up and the compressed gas will get stretched in the azimuthal direction of the galaxy (Palouš et al. 1990). Thus we require $t_{\text{clouds form}} < 2.5/\kappa$, which for the case of the expanding ring, gives approximately

$$\frac{1.5}{(G\rho_0)^{1/2}\mathscr{M}} < \frac{2.5}{\kappa} . \tag{31}$$

Note that a ring is most appropriate for very large scales, once R exceeds the scale height of the disk.

Equation (31) corresponds to a condition on the stability parameter Q if ρ_0 is converted into a disk gas column density σ using the scale height H given above. Then we get

$$Q = \frac{\kappa c_0}{\pi G \sigma} < \frac{2.5\mathscr{M}}{1.5(2\pi[1 + \alpha + \beta]\rho_0/\rho_{0T})^{1/2}}$$
$$\approx 0.66\mathscr{M} \approx 0.37\left(\frac{R}{r}\right)^{1/2} \tag{32}$$

for $(1 + \alpha + \beta)\rho_0/\rho_{0T} = 1$. This value of Q is comparable to that found by Kennicutt (1989) for the threshold of star formation in galactic disks, namely, $Q \approx 1.4$, if $\mathscr{M} \approx 2$ (or $R/r \sim 14$), which is approximately what we expect for energetic or magnetic rings with internal total pressure-to-density ratios comparable to the square of the shock speed, or for old rings that have slowed down to near the ambient rms velocity dispersion. Note that in most galaxy disks with $n_0 \sim 1$–$4\ \text{cm}^{-3}$, $(G\rho_0)^{-1/2}$ is longer than the duration of the high pressure phase of a supershell, so $\mathscr{M}$ will have decreased to near unity at the time of collapse. Rings or shells in the ambient medium that dissipate a significant fraction of the energy in the swept-up material ($\mathscr{M}$ large) can collapse with higher ambient Q, as can rings or shells that propagate into regions with above average densities (i.e., into molecular clouds). Nevertheless, Q serves as a useful discriminant for both stimulated and spontaneous cloud formation processes on a large scale, because they both involve a balance between gravitational and Coriolis forces.

These results imply that spontaneous star formation driven by gravitational instabilities and stimulated star formation driven by centralized pressure sources are likely to be important in the same regions of a normal galaxy disk. Thus observations of a surface density threshold for star formation (Kennicutt 1989) do not necessarily imply that all star formation results from spontaneous collapse; a significant amount of star formation throughout the disk can be from stimulated processes too. Conversely, galactic disk models with only propagating star formation are likely to be missing an equally important contribution from large-scale spontaneous processes, which should operate simultaneously in normal galaxies.

This balance between spontaneous and stimulated star formation could be very different when the ambient gas density is so large that the basic gravitational timescale of $(G\rho_0)^{-1/2}$ is less than the duration of the shell-driving pressure, which is determined by the timescale for massive stellar evolution. Then $\mathscr{M}$ could be much larger than order unity at the time of collapse because the shock-driving pressures are still strong then. If the ambient density exceeds $10^3 m_{\rm H}\ \text{cm}^{-3}$, for example, then $(G\rho_0)^{-1/2} < 3$ Myr, which is the lifetime of an early type O star. In this limit, which may apply to starburst galaxy nuclei, the analysis suggests that stimulated star formation could be much more important than spontaneous star formation from gravitational instabilities primarily because the stimulated time scales of $(G\rho_0\mathscr{M})^{-1/2}$ or $(G\rho_0\mathscr{M}^2)^{-1/2}$, depending on geometry, are much less than the spontaneous timescale $(G\rho_0)^{-1/2}$ when $\mathscr{M} \gg 1$ at the time of collapse. Nuclear disks with high densities could even have $Q > 1$ from the rapid rotation but if $\mathscr{M} \gg 1$ at the time of collapse, then Q can still be less than $0.66\mathscr{M}$ when clouds form. In this case only the stimulated star formation mechanism would operate (cf. Morris 1993).

Such a change in the ISM properties for very dense regions, leading to a change in the balance between spontaneous and stimulated star formation, would lead to a qualitative difference in the morphology of star formation. Instead of observing most star formation in the dense cores of clouds that contain the local Jean mass, star formation in very dense regions could be more scattered in the swept-up debris from other star formation.

REFERENCES

Castor, J., McCray, R., & Weaver, R. 1975, ApJ, 200, L107
Comeron, F., & Torra, J. 1994, ApJ, 423, in press
Elmegreen, B. G. 1982, in Submillimeter Wave Astronomy, ed. J. E. Beckman & J. P. Phillips (Cambridge Univ. Press), 3
———. 1985a, in Birth and Evolution of Massive Stars and Stellar Groups, ed. W. Boland & H. Van Woerden (Dordrecht: Reidel), 227
———. 1985b, in Birth and Infancy of Stars, ed. R. Lucas, A. Omont, & R. Stora (Amsterdam: Elsevier Science Pub. B.V.), 215
———. 1987, in IAU Symp. 115, Star-forming Regions, ed. M. Peimbert & J. Jugaku (Dordrecht: Reidel), 457
———. 1988, ApJ, 326, 616
———. 1992, in Star Formation in Stellar Systems, ed. G. Tenorio-Tagle, M. Prieto, & F. Sanchez (Cambridge Univ. Press), 381
Franco, J., Tenorio-Tagle, G., Bodenheimer, P., Rozyczka, M., & Mirabel, I. F. 1988, ApJ, 333, 826
Hunter, J. H., Jr., Sandford, M. T., II., Whitaker, R. W., & Klein, R. I. 1986, ApJ, 305, 309
Kennicutt, R. C. 1989, ApJ, 344, 685
MacLow, M.-M., & Norman, R. L. 1993, ApJ, 407, 207
McCray, R., & Kafatos, M. 1987, ApJ, 317, 190
Morris, M. 1993, ApJ, 408, 496
Nishi, R. 1992, Prog. Theor. Phys., 87, 347
Ögelman, H. B., & Maran, S. P. 1976, ApJ, 209, 124
Olano, C. A. 1982, A&A, 112, 195
Ostriker, J. P., & Cowie, L. L. 1981, ApJ, 243, L127
Palouš, J., Franco, J., & Tenorio-Tagle, G. 1990, A&A, 227, 175
Tenorio-Tagle, G. 1981, A&A, 94, 338
Tenorio-Tagle, G., & Bodenheimer, P. 1988, ARA&A, 26, 146
Tenorio-Tagle, G., & Palouš, J. 1987, A&A, 186, 287
Vishniac, E. T. 1983, ApJ, 274, 152
———. 1994, ApJ, submitted
Vishniac, E. T., & Ryu, D. 1989, ApJ, 337, 917
Westerlund, B. E., & Mathewson, D. S. 1966, MNRAS, 131, 371
Yoshida, T., & Habe, A. 1992, Prog. Theor. Phys., 88, 251