THE BIRTH OF MOLECULAR CLOUDS: FORMATION OF ATOMIC PRECURSORS IN COLLIDING FLOWS

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ABSTRACT

Molecular cloud complexes (MCCs) are highly structured and “turbulent.” Observational evidence suggests that MCCs are dynamically dominated systems, rather than quasi-equilibrium entities. The observed structure is more likely a consequence of the formation process than something that is imprinted after the formation of the MCC. Converging flows provide a natural mechanism to generate MCC structure. We present a detailed numerical analysis of this scenario. Our study addresses the evolution of an MCC from its birth in colliding atomic hydrogen flows up until the point when H$_2$ may begin to form. A combination of dynamical and thermal instabilities breaks up coherent flows efficiently, seeding the small-scale nonlinear density perturbations necessary for local gravitational collapse and thus allowing (close to) instantaneous star formation. Many observed properties of MCCs come as a natural consequence of this formation scenario. Since converging flows are omnipresent in the ISM, we discuss the general applicability of this mechanism, from local star formation regions to galaxy mergers.

Subject headings: ISM: clouds — ISM: molecules — methods: numerical — turbulence

1. MOTIVATION

Molecular cloud complexes (MCCs) in the Galaxy show an abundance of internal structure, in both density and velocity (e.g., Stutzki & Güsten 1990; Falgarone 1990; Mizuno et al. 1995). The density contrasts are nonlinear, independent of the importance of gravity (Williams et al. 1995), and the MCCs exhibit nontermal line widths (Falgarone & Philips 1990; Williams et al. 2000), generally interpreted as supersonic turbulence. Filaments seem to dominate the morphologies in (column) density and velocity (e.g., Bally et al. 1987; Mizuno et al. 1995; Hartmann 2002; Churchwell et al. 2004). The spatial distribution of densities and velocities is consistent with a turbulent spectrum, the details of which are a matter of debate, however (Elmegreen & Scalo 2004).

Extremely puzzling is the source of this wealth of structure. While stellar feedback certainly is a powerful driver (see, e.g., Mac Low 2004), it only acts locally and definitely only after the first stars have formed within the cloud. Moreover, stellar driving might be difficult to reconcile with the observed nearly self-similar spatial energy distribution.

External drivers suffer from the fact that the cold dense gas essentially acts like a wall to any incoming wave (e.g., Vasquez 1990; Balsara 1996; Elmegreen 1999), preventing an efficient energy transfer from the warm diffuse component to the cold dense phase (see, however, Misch & Zweibel [1994] for one-dimensional models). Hennebelle & Inutsuka (2006) argue that openings (holes or channels) in an MCC serve as inroads for Alfven waves; however, the question of what forms these channels remains.

With growing evidence that the lifetimes of molecular clouds, at least in the solar neighborhood, range around a few (2–3) Myr (Elmegreen 2000; Hartmann et al. 2001; Hartmann 2003), a picture in which MCCs are envisaged as transient objects in large-scale atomic colliding flows, rather than as well-defined entities in a quasi-equilibrium state (giant molecular cloud [GMC]), is emerging. The colliding flows would accumulate atomic gas that might eventually reach column densities high enough for H$_2$ formation, at which point the “lifetime” of the molecular cloud would start; i.e., the accumulation time could be much longer. Some of the molecular regions might become self-gravitating and form stars instantaneously (Elmegreen 1993, 2000; Ballesteros-Paredes et al. 1999a; Hartmann et al. 2001; Pringle et al. 2001; Hartmann 2003). The large-scale flows might be driven by supernova explosions or Galactic shear or might occur in galaxy interactions (for recent evidence of cloud collisions, see Looney et al. 2006).

We propose and extend the idea (see § 2) that the structures observed in MCCs arise from the very processes that form them; i.e., they arise without recourse to nonlinear perturbations introduced by hand, an approach that defers the problem to an even earlier stage. We numerically investigate the formation of structure in MCC precursors via colliding flows. To avoid the somewhat unwieldy term “precursor of molecular cloud complexes,” we abbreviate this to PoMClOc. By this we mean cold clouds (usually identified as cold neutral medium [CNM]), which may be fully atomic or may already contain traces of H$_2$. In any case, the PoMClOc’s serve as the initial stage of MCCs.

Our numerical models emphasize the ease with which colliding flows generate structure via a combination of dynamical and thermal instabilities. Due to the thermal instability, structure grows predominantly on small scales, leading to early nonlinear density perturbations as possible seeds of gravitational collapse and star formation. The resulting line widths in the cold gas when seen in projection agree with observed values; however, the cold gas does not seem to exhibit internal supersonic turbulence. Only a few percent of the energy input remains available for driving turbulent motions in the cold gas; the bulk of the energy is lost due to radiative cooling. Within the scope of the models, we find that the timescale for H$_2$ formation is limited not by a minimum temperature, but by the highly unsteady environment.

This work extends our previous study (Heitsch et al. 2005), in dimensions (here we present corresponding three-dimensional models), in resolution, and in scope of physical applications. Of course, we rely heavily on earlier work (§ 2). The physics of the problem is described in § 3, and the numerical realization and
numerical artifacts are discussed in §§ 4 and 6.1, respectively, in order to provide a background for the interpretation of the results (§ 5). These are summarized in § 6, while § 7 suggests possible routes to follow in the future.

2. PREVIOUS AND CURRENT WORK

The concept behind this study is that molecular cloud formation should be seen as a nonequilibrium process. When the perception of the interstellar medium (ISM) as a dynamical medium gained acceptance, many aspects of molecular cloud and star formation theory were revisited. Hunter (1979) and Hunter & Fleck (1982) reexamined the Jeans criteria and found that converging velocity fields could provoke the gravitational collapse of sub-Jeans mass gas clouds. Tohline et al. (1987) followed this line of thought, further emphasizing that diffuse gas cooling under compression by even mild velocity and pressure disturbances could easily be transformed to high densities. They applied this idea to externally perturbed spherical clouds. Motivated by these investigations into the role of compressive velocity fields in promoting fast cooling and gravitational collapse, Hunter et al. (1986) used two-dimensional numerical simulations to study the cooling and fragmentation of gas compressed at the interface of two identical, oppositely directed supersonic colliding gas streams.

Since these early works, which explored some consequences of a dynamical ISM, molecular cloud formation and evolution has been studied within the framework of compressible turbulence. A strongly fluctuating medium with motions on many scales results in high-density, albeit transient, regions that might be identified as molecular clouds. Large-scale (kiloparsec) models of molecular cloud formation were presented in two-dimensional simulations by Vázquez-Semadeni et al. (1995) and Passot et al. (1995), throwing light on the role of turbulence in the cycle of cloud and star formation ( Larson 1981) and specifying conditions for cloud formation in the presence of magnetic fields. But the numerical resolution needed for a detailed explanation of the structure and dynamics of molecular clouds makes it necessary to tackle the problem with smaller scale simulations (tens of pc).

For instance, the role played by thermal instability continues to be actively explored. This instability alone has been shown to be an efficient mechanism for generating substructure and even driving turbulence (Burkert & Lin 2000; Kritsuk & Norman 2002a, 2002b, 2004). As indicated by earlier work (e.g., Priest & Heyvaerts [1974] in an application to the solar atmosphere), thermal instability can be triggered by compression. Indeed, one-dimensional plane-parallel simulations of transonic converging flows (Hennebelle & Pérault 1999, 2000) show how this mechanism transforms inflowing diffuse atomic gas to dense cold gas that is long lived. Shock waves (driven, e.g., by supernova explosions; Mac Low et al. 2005; de Avillez & Breitschwerdt 2005) are another means to create cold, dense gas from warm, diffuse gas. Propagating into the warm ISM, shock waves have been shown to fragment in the presence of thermal instability (Field 1965) and linear perturbations (Koyama & Inutsuka 2000, 2002; Inutsuka & Koyama 2004), leading to $H_2$ formation within a few Myr (Bergin et al. 2004). However, the passage of a single shock might not leave in its wake enough dense, cold cloudlets to constitute an MCC. Encounters between streams of transsonically transported gas, on the other hand, might be a way to collect, compress, and cool large quantities of gas while at the same time possibly endowing the cold, dense gas with the dynamical and structural characteristics of molecular cloud complexes. This scenario has recently come into the spotlight.

In high-resolution two-dimensional simulations of supersonic colliding gas flows, Audit & Hennebelle (2005) confirm that even in a very dynamical environment a bistable medium develops as compressions initiate thermal instability. However, their study concentrates on how turbulence generates and influences thermally unstable gas, not on the reverse problem of how thermal instability could feed the turbulence characteristic of molecular clouds. In three-dimensional simulations of transonic colliding flows, Vázquez-Semadeni et al. (2006) show that a thin cold sheet, reminiscent of those observed by Heiles & Troland (2003) and Heiles (2004), forms at the junction of the two flows. It develops turbulence, which they attribute to the nonlinear thin-shell instability (NTSI; Vishniac 1994). Furthermore, they find that even in simulations without gravity, the highest density gas is overpressured with respect to the mean warm neutral medium pressure, suggesting that the ram pressure of the colliding flows is responsible.

To be fair, we should mention that favorable conditions for molecular cloud formation exist not only in converging flows, but also in “focal planes” of (sustained) MHD waves, as Elmegreen (1999) shows. However, he notes that while highly structured MCCs form in such a system, the cold gas does not acquire significant turbulence.

The notion of generating the turbulent substructure of molecular clouds by their formation process has been discussed in various contexts (e.g., Ballesteros-Paredes et al. 1999a; Hartmann et al. 2001; Audit & Hennebelle 2005; Vázquez-Semadeni et al. 2006). Audit & Hennebelle (2005) emphasized the evolution of the thermal states in the colliding flows and provided a semi-analytical model for the evolution of a gas parcel. One of the main differences with the present study is that they impose velocity perturbations on their incoming gas flow, as do Vázquez-Semadeni et al. (2006), although as “random noise” at the 1% level, whereas we defer the structure generation to the actual flow collision site. Vázquez-Semadeni et al. (2006) compared in detail analytical solutions with one-dimensional and three-dimensional models. Koyama & Inutsuka (2002) and their subsequent work discuss molecular cloud formation behind a shock-compressed layer, more closely resembling the situation of an expanding shell. In their case, the initial perturbations reside in the density field. The present study extends the “proof of concept” of “structure formation from scratch” in two dimensions presented in Heitsch et al. (2005) by focusing on the dynamical state and the turbulent properties of the PoMClOc during its formation and on the conditions necessary for the onset of $H_2$ formation.

3. PHYSICS

We restrict the models to hydrodynamics with radiative cooling, leaving out the effects of gravity and magnetic fields. Gravity would lead to further fragmentation, and magnetic fields could have a stabilizing effect (see also discussion in § 7). For this regime, then, we identify three relevant instabilities, namely, the NTSI, the Kelvin-Helmholtz instability (KHI), and the thermal instability (TI). This enumeration, of course, does not mean that the instabilities work independently of each other. Rather, the resulting dynamical system is a consequence of a combination of all three instabilities, however, in degrees depending on the local flow properties.

3.1. NTSI

The NTSI (Vishniac 1994; Fig. 1) arises in a shock-bounded slab, when ripples in a two-dimensional slab focus incoming shocked material and produce density fluctuations. Gas streaming along the x-direction will be deflected at perturbation peaks and collect in the troughs. Thus, x-momentum is transported laterally
For a weak equation of state, the slab can act as a wall to deflect the incoming gas streams, while for an adiabatic one, the bounding shocks will travel faster at the abatic one, the bounding shocks will travel faster at the act as a wall to deflect the incoming gas streams, while for an adiabatic one, the bounding shocks will travel faster at the abatic one, the bounding shocks will travel faster at the...
We chose a method based on the second-order Bhatnagar-Gross-Krook (BGK) formalism (Prendergast & Xu 1993; Slyz & Prendergast 1999; Heitsch et al. 2004; Slyz et al. 2005b), allowing control of viscosity and heat conduction. The code evolves the Navier-Stokes equations in their conservative form to second order in time and space. The hydrodynamical quantities are updated in time-unsplit form. As shown below, statistical properties of the models are resolved with respect to grid resolution, viscosity, and heat conduction, although the flow patterns change in detail, as expected in a turbulent environment.

4.1. Heating and Cooling

The heating and cooling rates are restricted to optically thin atomic lines, following Wolfire et al. (1995), so that we are able to study the precursors of MCCs up to the point when they could form H₂. Dust extinction becomes important above column densities of $N(H\text{I}) \approx 1.2 \times 10^{21} \text{ cm}^{-2}$, which are reached only in the densest regions modeled. Thus, we use the unattenuated UV radiation field for grain heating (Wolfire et al. 1995), expecting substantial uncertainties in cooling rates only for the densest regions. The ionization degree is derived from a balance between ionization by cosmic rays and recombination, assuming that Lyα photons are directly reabsorbed. Numerically, heating and cooling is implemented iteratively as a source term for the internal energy $e$ of the form

$$\partial_t e = n \Gamma(T) - n^2 \Lambda(T) \tag{2}$$

in units of energy per volume per second, where $n \Gamma$ is the heating contribution (mainly photoelectric heating from grains) and $n^2 \Lambda$ is the cooling contribution (mainly due to the $C\text{II}$ hyperfine structure line at 158 μm). Since the cooling and heating prescription has to be added outside the flux computations, it lowers the time order of the scheme. To speed up the calculations, equation (2) is tabulated on a 2048×2 grid in density and temperature. For each cell and iteration, the actual energy change is then bilinearly interpolated from this grid.

4.2. Initial and Boundary Conditions

Two opposing, uniform, identical flows in the $x$-$y$ computational plane initially collide head-on at a sinusoidal interface with wavenumber $k_x = 1$ (and $k_y = 1$ for three-dimensional models) and amplitude $\Delta$ (Fig. 1). The incoming flows are in thermal equilibrium. The system is thermally unstable for densities $1 \text{ cm}^{-3} \leq n \leq 10 \text{ cm}^{-3}$. The cooling curve covers a density range $10^{-3} \text{ cm}^{-3} \leq n \leq 10^3 \text{ cm}^{-3}$ and a temperature range $30 \text{ K} \leq T \leq 1.8 \times 10^4 \text{ K}$. The box side length is 44 pc. Thus, Coriolis forces from Galactic rotation are negligible; however, the simulation domain is large enough to accommodate a moderately sized molecular cloud complex. For an interface with $\Delta = 0$, a cold high-density slab devoid of inner structure forms. The onset of the dynamical instabilities thus can be controlled by varying the amplitude of the interface perturbation. This allows us to test turbulence generation under minimally favorable conditions. This setup might seem artificial because (1) the incoming flows are not expected to be perfectly uniform; however, we chose to defer the moment of structure generation to the last possible moment in the simulation, instead of imposing perturbations on the incoming flows (see § 4.3), and, (2) as we see below, the model run times extend considerably beyond 10 Myr. At an inflow speed of $10 \text{ km s}^{-1}$ this would correspond to a total extent of the system of 200 pc. This suggests that the initial densities of the flows are more likely to be a few cm$^{-3}$ to form molecular clouds with flows of order 50–100 pc length. Spiral density waves can also produce coherent flows of the length required, at least in principle.

Fig. 2.—Logarithmic temperature maps of the two-dimensional models 2C10c, 2C20c, and 2C30c, i.e., with inflow speeds at Mach 1, 2, and 3. The temperature range is given in log $T$, where $T$ is in kelvin. The stills were taken approximately 11.5 Myr after flow contact.

Fig. 3.—Left: Asymmetry around midplane $y = 0$ plotted against time for model sequence 2C at the six available Mach numbers for resolution $N = 512$. Right: Asymmetry around midplane $y = 0$ plotted against time for model sequence 2C at resolutions $N = 512$ and 1024.
The boundary conditions in the transverse (i.e., y and z) directions are periodic, while in the x-direction, the boundary values are defined as uniform inflow at constant density \( n_0 \) and inflow speed \( v_0 \). Consequently, the boundaries cannot treat outgoing waves or flows. Thus, the models have to be halted once the bounding shocks reach the boundaries. The corresponding timescales have been discussed by Vázquez-Semadeni et al. (2006).

To facilitate the analysis of the dynamics and time history of the cold gas, the code is equipped with Lagrangian tracer particles, which are advected after each flux update. Tracer particles are deployed at grid cell centers where the average is over all cells with \( \Delta n > 0 \) (not over the whole domain, to exclude the inflow initial conditions, which are symmetric) and \( \Delta n \) is the absolute value of the density difference between cells, which should be symmetric across the upper and lower half of the simulation midplane. The initial conditions are perfectly symmetric, but slight differences at the machine accuracy level in the reconstruction of hydrodynamical variables eventually lead to a difference in the cooling strength and thus cause perceptible asymmetries. The code preserves perfect symmetry for a purely adiabatic equation of state (i.e., without the additional heating and cooling terms). After onset, the asymmetries grow linearly in time until they reach a saturation level between \( A \approx 50 \) and 100. This corresponds approximately to the temperature (and density) contrast between the warm and cold gas and thus to the maximum asymmetry reachable for the system. Although the asymmetries increase with Mach number and, to a lesser extent, with resolution, they only appear well after the system has evolved. Koyama & Inutsuka (2002), Audit & Hennebelle (2005), and Vázquez-Semadeni et al. (2006) chose an alternative route. They added perturbations (in density or velocity) to the incoming flow, thus breaking the symmetry of the initial conditions. While the physical reason for adding perturbations to the inflow is perfectly obvious, in the present study we want to emphasize the point that even with the least possible perturbation to the flow, (turbulent) substructures are generated with ease. In a sense, this is an attempt to carry the argument to extremes.

Conversely, one could argue that imposing perturbations on the incoming flow helps to hide the symmetry breaking due to truncation errors. To get a better understanding of how our somewhat extreme initial conditions are affecting the results, we repeated model 2C20 twice, once with an interface perturbation mode of \( k_y = 32 \), instead of \( k_y = 1 \), and, in the second case, with the \( k_y = 32 \) mode overlaid on the \( k_y = 1 \) mode at a fraction of \((1/32)^{3/2}\) of the amplitude at \( k_y = 1 \), motivated by a turbulent cascade. In both cases, the system develops small-scale structures at \( k_y = 32 \), meaning that the instability grows at the smallest imposed (\( k_y = 32 \)) scales, as long as these are larger than the Field length (see § 4.4).
where $\kappa$ is the heat conductivity (see below) and $\Lambda$ is the cooling function from equation (2). With a heat conductivity of $\kappa = 2.5 \times 10^3 (T/[K])^{1/2}$ (Parker 1953), Koyama & Inutsuka (2004) conclude that for ISM conditions, a linear resolution of several thousand cells is needed (they get convergence at 16384 cells).

The heat conduction in the BGK scheme can be controlled explicitly (see Slyz et al. 2005a for an analysis) by varying the kinematic viscosity, since the Prandtl number in the code is $Pr \equiv 1$ by construction. The choice of the viscosity is controlled by two considerations, namely, (1) it should be large enough to prevent numerical artifacts on small scales and (2) it should be small enough to leave enough dynamical range. Obviously, both requirements are difficult to meet simultaneously.

To establish to what degree our models are resolved, we begin by measuring the Field lengths (eq. [4]) cellwise, excluding all cells that belong to the inflow, since their thermal timescale $\tau_c \rightarrow \infty$. Figure 4 shows the volume and mass fractions of the Field-unresolved cells as a function of time for models 2C20b and 2C30b. Note that when the volume and mass fractions agree, the bulk of the nonequilibrium gas is in the cold phase (see also Koyama & Inutsuka 2002; Audit & Hennebelle 2005). In other words, Figure 4 and the following ones refer mostly to the cold gas phase. Thus, between 10% and 20% of the mass of the cold gas are not resolved. This is acceptable as long as we devise a selection criterion for the subsequent analysis.

Figure 5 lets us estimate the size of the smallest structures we can resolve. Again, we distinguish between volume and mass fractions, which, however, in this case are nearly identical. The figure indicates that the Field length for the bulk of the mass and volume lies between 1 and 32 pc, scales that are well resolved by all our models. For our lowest resolution two-dimensional runs ($N = 512$), the resolution limit is $\Delta x = 0.34$ pc. Thus, for the subsequent discussions, we will only consider structures with sizes larger than 0.34 pc for two-dimensional runs and 0.68 pc for three-dimensional runs.

If the Field length is not resolved, structures tend to grow on grid scales; i.e., with increasing resolution, there should be more and more small-scale structures. This effect is clearly visible in Figure 6 (left), which shows a histogram of the size of cold regions ($T < 300$ K). The size of a cold region is defined as the geometric mean of its minimum and maximum diameter. The inset numbers give the average number of cold regions per line of sight, i.e., normalized to resolution. If the Field length were resolved, heat conduction would lead to a cutoff at small $L_{\text{cold}}$. Selecting for regions with $L_{\text{cold}} > 0.34$ pc, the histograms (and the average number of cold regions) agree sufficiently to proceed with the above selection criterion.

5. RESULTS

5.1. Turbulence

Molecular clouds consistently show nonthermal line widths of a few km s$^{-1}$ (e.g., Falgarone & Phillips 1990; Williams et al. 2000) that, together with temperatures of $T \approx 10$ K, are generally interpreted as supersonic turbulence. The line widths in our models are consistent with the observed values (Fig. 7). The broad-line wings are non-Gaussian. This may be a sign of intermittency (e.g., Falgarone & Phillips 1996; Lis et al. 1996).

The “observed” line width is derived from the histogram of the density-weighted line-of-sight velocity dispersion in the cold gas at $T < 300$ K (Fig. 8, filled symbols). This line width would correspond to line widths in the CNM as, e.g., traced by H$_1$. Since the internal line widths of coherent cold regions (open symbols) range around the sound speed of the cold gas (0.7 km s$^{-1}$), the internal velocity dispersions do not reach Mach numbers $M > 1$ (see also Koyama & Inutsuka 2002; Audit & Hennebelle 2005). Hence, the “supersonic” line widths are a consequence of cold
regions moving with respect to each other within a warmer and more diffuse medium, but not a result of internal supersonic turbulence in the cold gas that would eventually host star formation (Kwan & Sanders 1986; Hartmann et al. 2001). Note from Figure 9 that this result is independent of resolution and geometry. Because of the TI, to make Figures 8 and 9, cold coherent regions are identified via a temperature threshold at \( T < 300 \) K. The variations of the line-of-sight velocity dispersion with time is within the error bars shown.

Are the line widths actually indicating turbulent motions, or are we seeing the inflow motions of (already) cold gas (e.g., Vázquez-Semadeni et al. 2006)? If the gas motions are truly turbulent, the average Lyapunov exponents

\[
\langle \lambda \rangle \equiv \left\langle \frac{1}{t - t_0} \log \frac{d(t)}{d(t_0)} \right\rangle \tag{5}
\]

should be positive, indicating a growing separation between neighboring particles. In equation (5), \( t_0 \) is the start time of particle advection, and \( d \) is the corresponding separation of neighboring particles. The average refers to the simulation domain. After an initial compression phase with \( \langle \lambda \rangle < 0 \), the cold gas becomes turbulent (Fig. 11). With higher Mach number, the onset of internal turbulence in the slab is delayed, which is mirrored by the lower Lyapunov exponents. Higher resolution leads to larger \( \langle \lambda \rangle \) at the time of the onset of turbulence, which is defined as \( t(\langle \lambda \rangle > 0) \). For \( N = 1024 \), the initial separation of tracer particles is smaller, so that a larger spatial range can be covered. All exponents are positive, are independent of resolution, and converge for late times.

The power spectrum of the specific kinetic energy \( v^2 \) (Fig. 12 for model 3C20) can be used as another diagnostic of turbulence in the models. Shown are the three linear spectra, taken along the inflow direction (\( \alpha_x \)) and the transverse directions (\( \alpha_y, \alpha_z \)). The spectral index \( \alpha_x = -1.96 \) is consistent with the Fourier transform of a step function, as expected since the strong decelerations along the \( x \)-direction effectively lead to a discontinuity in \( v^2 \). (Note that for the initial conditions, not \( v^2 \), but \( u_i \) is discontinuous.) The spectral indices in the transverse directions, \( \alpha_x \) and \( \alpha_z \), are consistent with a Kolmogorov spectrum, indicating fully developed turbulence. The lower spectrum (denoted by triangles and a corresponding slope \( \alpha_z = -3.24 \)) shows the spectrum averaged over spheres of constant wavenumber \( k \). Error bars denote errors on the mean. Within the errors, the slope is still consistent with a Kolmogorov spectrum of \(-11/3\) in three dimensions. The last half-decade is dominated by numerical diffusion, leaving us approximately 1.5 decades for physical analysis.

The spectral indices differ depending on whether the spectrum is taken along the inflow or transversely. Another way to see this is to split the specific kinetic energy into a compressible (and translational) part with \( \nabla \times \vec{v} \equiv 0 \) and a solenoidal part with \( \nabla \cdot \vec{v} \equiv 0 \). For the full domain (models 2C20 and 3C20), the specific kinetic energy is clearly dominated by the compressible part, because of the inflows (Fig. 13, left). Note, however, that the solenoidal part (\textit{thin lines}) is steadily increasing. Even though the energy input rate is constant, the total specific kinetic energy (\textit{thick lines}) decays with time as a consequence of the radiative cooling (see below). The specific energy restricted to the cold gas (at \( T < 300 \) K; Fig. 13, right) is dominated by the solenoidal component, and all components increase with time, indicating that once the gas has cooled down to its minimum temperature, it starts to store the kinetic energy from the inflows. The solenoidal components are slightly larger for the three-dimensional models, which is not surprising, since the extra degree of freedom allows the gas to evade compression more efficiently.

One might speculate as to whether the total specific kinetic energy will decrease until it has reached a minimum at the moment when (most of) the gas has cooled down to the minimum temperature given by the cooling curve or whether it will find some kind of equilibrium between kinetic energy input and thermal energy loss. Since the cooling gets more and more efficient with increasing density, the latter seems unlikely. However, it is probably even more unlikely that the inflows are maintained for times long enough to establish equilibrium between kinetic energy input and thermal energy loss. Moreover, once gravity dominates the cold, dense regions, rather than reach a state of equilibrium, the clouds might pass through the phases of initial compression, turbulence generation, cooling, and, finally, gravitational collapse.
5.2. Thermal States

Structure forms in colliding flows as a result of an interplay between dynamical instabilities and TIs. The dynamical instabilities generate compressions and shear flows, while the thermal instabilities amplify density perturbations to the nonlinear regime, where they become potential seeds for self-gravitating structures. The effect of the TI on the thermal state of the gas is shown in Figure 14.

The solid line denotes the $P(n)$ relation in thermal equilibrium. It is a direct result of the heating and cooling processes included and displays the usual three regimes, namely, an “atomic” regime for $\log n < 0.5$ with an effective $\gamma_e \approx 5/3$, an “isothermal” regime for $\log n > 1.5$ with $\gamma_e \to 1$, and an unstable regime with $\gamma_e < 0$.

Each cell in the scatter plots in Figures 14 and 15 is color-coded with its thermal timescale (eq. [1]). Positive timescales denote heating (toward red colors), and negative timescales correspond to cooling (toward blue colors). At high densities, the thermal timescales are short (greenish colors), which is mirrored in the very small scatter in $\log P$ around the equilibrium curve. Moving to lower densities, the cooling timescale increases dramatically, until we reach the instability regime, which consists of a black triangular region in $P(n)$. This consists mostly of gas that has passed the bounding shocks (slightly increased densities due to the [nearly] adiabatic shock), but has not yet reached the cold dense phase. Since there is no gravity in the simulation, the ram pressure of the inflow (Fig. 14, dashed lines) limits the maximum pressure gas can reach in the system, apart from slight overshoots due to waves. A substantial amount of gas resides at higher pressure than that of the ambient inflow (higher by a factor of approximately 6 for model 2C20 and by 13 for model 2C30). This gas has been “overpressured” by the highly compressible inflow. Even the gas in the isothermal branch (i.e., gas that has cooled down to the minimum temperature) resides at pressures between the initial thermal equilibrium pressure and the ram pressure of the inflow (see also Vázquez-Semadeni et al. 2006; an effect that has traditionally been assigned to self-gravity; e.g., Pringle et al. 2001). The more or less linear black regions at $\log n \approx 0$ are a result of the adiabatic reaction of the shock-compressed gas; i.e., this gas has not yet had time to react to thermal effects (see also Young & Mac Low 2006). The dashed red line denotes the Hugoniot curve

$$\left(n_2 - \frac{\gamma - 1}{\gamma + 1} n_1\right) P_2 = \left(n_1 - \frac{\gamma - 1}{\gamma + 1} n_2\right) P_1;$$

that is, the pairs of values $(P, n)$ for the state on one side of the shock front (index 1) that are compatible with the Rankine-Hugoniot jump conditions across the shock front (index 2; Courant & Friedrichs 1948). This has also been noted by Audit & Hennebelle (2005).

From Figure 15 we see that the thermally unstable gas is generally still fast-moving; i.e., it has passed the bounding shocks, but has not yet met the slab of dense, cold material [of course, the initial inflow just shows up as a single dot in the $P(n)$ and $\epsilon(n)$ plots]. High-density gas is mostly moving at a few km s$^{-1}$, consistent with the line-of-sight velocity dispersion discussed in §5.1.

A small fraction of cells $(1.7 \times 10^{-3}$ for model 2C20 and $1.9 \times 10^{-4}$ for model 2C30) seem to exhibit velocities larger than the inflow velocity. This is hard to understand because there is no physical mechanism to accelerate gas to velocities higher than

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**Figure 11.** Left: Average Lyapunov exponents (eq. [5]) for a sequence of models in Mach number plotted against time. After the initial compression phase, the cold gas becomes turbulent. Only the first tracer particle generation is used. Models shown in this diagram were run at $N = 512$. Right: Resolution effects. Higher resolution leads to larger $\langle \Lambda \rangle$, although all $\langle \Lambda \rangle > 0$.

**Figure 12.** Specific kinetic energy spectrum for model 3C20 at 15 Myr. The linear spectral indices $\alpha_x, \alpha_y$, and $\alpha_z$ refer to the coordinate directions, and $\alpha_c$ gives the index of the spherically averaged spectrum. Along the inflow, the strong decelerations effectively lead to a step function in $(\epsilon^2)$, mirrored in the steeper index. Due to numerical dissipation, the last half-decade is not available for physical analysis.

**Figure 13.** Left: Specific kinetic energy for the whole simulation domain of models 2C20 and 3C20 plotted against time, split into the compressible part, the solenoidal part, and the total. Compressible motions dominate because of the inflows. Right: Specific kinetic energy for the cold gas ($T < 300$ K) of models 2C20 and 3C20 plotted against time, split into the compressible part, the solenoidal part, and the total. Solenoidal motions dominate in the cold gas.
inflow velocities. All of those cells coincide with the largest density/temperature jumps occurring in the simulation and have slightly increased temperatures as well. Thus, they are a consequence of a slight overshoot from the reconstruction step in the scheme. Since the fraction of faulty cells does not vary over time (once the cold gas has formed), we simply neglect these cells for the analysis.

For our later discussion of critical masses for molecular cloud formation, Figure 16 summarizes the mass content in the three thermal regimes, defined by the warm stable range \( T > 3000 \text{ K} \), the cold stable range \( T < 300 \text{ K} \), and the instability range in between. Obviously, once the cold regions form, most of the gas is locked there. The amount grows linearly with time, indicating that although the PoMClO is turbulent, it essentially acts like a slab for collecting cold material. The warm phase (Fig. 16, middle part of each panel) contains a nearly constant amount of mass with time; it is just a transitory phase in our models. At any given time only a small fraction (but still a few percent) of the total mass is passing through the unstable regime. This mass fraction depends slightly on the Mach number of the inflow, because the shorter dynamical timescales compete with the thermal timescales, so that gas that has been thrown out of thermal equilibrium has a smaller chance of attaining equilibrium conditions again. A similar effect can be seen in the temperature maps (Fig. 2). For lower Mach numbers, the transition from warm (white) to cold (black) gas is much more pronounced (see also Sánchez-Salcedo et al. 2002; Vázquez-Semadeni et al. 2006). Conversely, a high fraction of gas in the unstable regime (Heiles & Troland 2003) might indicate a highly dynamical environment. A similar effect exists for three-dimensional versus two-dimensional models. Since the compressible modes are stronger in the two-dimensional case, gas tends to be forced to cool in compressions, while in the three-dimensional case, it can evade the compression and cooling by shearing motions, which would even lead to additional heating. As Figure 16 shows, neither resolution nor geometry affects the mass fractions in the stable temperature regimes significantly. In particular, the fraction of mass in the cold phase (Fig. 16, left part of each panel) is indistinguishable for models at resolutions \( N = 512 \) and 1024.

5.3. Driving Efficiency

Comparing the kinetic energy of the inflow to that in the unstable and cold (see § 5.2) gas phases allows us to estimate the efficiency of turbulent “driving” in our models, i.e., the efficiency with which the highly ordered kinetic energy of the inflow is converted to the turbulent motions within the PoMClO. Figure 17 shows the ratio of the kinetic energy change to the energy input rate of the inflow, i.e., the efficiency,

\[
\mathcal{E} \equiv \frac{\partial \left( \frac{1}{2} n v^2 \right)_0}{n_0 v_0 A},
\]

in the cold (Fig. 17, filled symbols) and unstable (Fig. 17, open symbols) gas. The integral extends over all cells within the chosen temperature regime, and \( A \) is the area of the inflow. Overall, the efficiency \( \mathcal{E} \) in the cold gas ranges between 2% and 5%. This is
consistent with the ratio of the inflow speed to the line-of-sight velocity dispersion (see, e.g., Fig. 9). In other words, most of the energy is lost due to atomic line cooling. Gas in the unstable phase reaches an efficiency that is smaller by a factor of ≈5, an effect of the lower densities (the velocities are generally higher; see Fig. 15). There seems to be a weak trend to smaller efficiency $E$ with larger energy inflow $n_g v_0$. Larger inflow velocities result in higher compressions and thus in faster cooling, by which a growing fraction of the input energy is lost.

5.4. Conditions for H$_2$ Formation

A crucial point in our analysis is whether the cold gas reaches conditions favorable for H$_2$ formation and thus eventually for molecular cloud formation (so far we have only been talking about the precursors of molecular clouds). Three criteria need to be met.

1. For H$_2$ formation, the gas temperature has to drop below a threshold temperature, which we take to be $T_c = 300$ K (Cazaux & Tielens 2004). This might appear to be somewhat high, since the sticking probability decreases rapidly with higher temperatures and is near unity only for $T < 20$ K (Cazaux & Tielens 2004; Bergin et al. 2004). Nevertheless, we chose the above temperature limit, because it turns out that once the temperature of a gas parcel has dropped below $T = 300$ K, it will quickly cool down to the minimum temperature allowed by the cooling prescription. In that sense, the timescales give a lower limit, i.e., denote the point at which the first H$_2$ could form. Since the cooling curve only extends down to 30 K anyway, we cannot make more detailed statements about H$_2$ formation. From Figure 16 we already saw that for reasonable inflow speeds, a few $\times$ 100 $M_\odot$ pc$^{-1}$ at $T < 300$ K accumulate within approximately 8 Myr. In principle, this material would be available for H$_2$ formation, at least as far as the temperature is concerned.

2. The (now cold) gas has to stay cold long enough to allow for H$_2$ formation. Time estimates for H$_2$ formation vary. In their analysis of H$_2$ formation behind shock fronts, Bergin et al. (2004) quote timescales between 5 and 10 Myr, depending on the inflow momentum. Once H$_2$ exists, its further formation could well be a runaway process, since self-shielding is much more efficient than dust shielding. This means that the critical timescale is given by the onset of H$_2$ formation. A minimum requirement therefore is that the cold gas is not reheated during this timescale. To measure this, we followed the temperature history of the tracer particles and determined how long each particle stays cold. Figure 18 (top) is a cumulative histogram of the time intervals over which tracer particles have temperatures $T < 300$ K. Apart from a small fraction at short time intervals $\Delta T$ (these are particles at the rims of the cold regions), most of the particles stay cold for at least 6 Myr. In fact, for most models, the cold gas parcels stay cold for over 14 Myr; i.e., once it has cooled down, by far most of the gas stays cold. Model 2C35 ended 6 Myr after the initial tracer deployment, so in order to compare to the other models, all had to be evaluated at a maximum time interval of $\Delta t = 6$ Myr, since otherwise it would look as if particles are reheated in model 2C35 after 6 Myr. To summarize the top panel of Figure 18, the gas stays cold long enough to allow for H$_2$ formation.

3. Finally, the cold, dense gas could be reexposed to the ambient UV radiation field. H$_2$ formation requires a critical column density of $N$(H) $\gtrsim$ 10$^{21}$ cm$^{-2}$. This we can determine again with the tracer particles, resulting in the bottom panel of Figure 18, which combines the temperature criterion of the top panel of Figure 18 and the column density threshold. Dropping the temperature criterion does not affect the result. Thus, the critical quantity is the shielding column density, not the temperature. In other words, once the gas enters the “cold” phase, its thermal timescale is short compared to the dynamical timescale, so that the gas stays isothermal (see Fig. 14). However, due to the continuous restructuring of the cloud, gas is repeatedly reexposed to the UV radiation field, and the column density of the cold gas only stays above 10$^{22}$ cm$^{-2}$ for about 1 Myr. This is a direct consequence of the PoMClO/C's highly dynamical nature. Note that Figure 18 gives a pessimistic view. Once a small fraction of the particles has reached conditions beneficial for H$_2$ formation, self-shielding will set in. The analysis up to now only includes shielding by dust. The inflow velocity of model 2C10 is too low to reach sufficiently large column densities for H$_2$ formation in the elapsed time of the simulation. As Bergin et al. (2004) found in one dimension, average column densities...
must achieve values of $10^{21}$ cm$^{-2}$ for H$_2$ formation to occur because of the shielding requirement.

6. A SUMMARY AND DISCUSSION

Molecular clouds in our Galaxy are complex and highly structured, with broad, nonthermal line widths suggesting substantial turbulent motions. Thus, molecular clouds are not very likely static entities and might not necessarily be in an equilibrium state, but their properties could well be determined by their formation process. We presented numerical models of the formation of precursors of molecular cloud complexes (PoMCloC’s), in large-scale colliding H i flows. We now summarize our findings and discuss their astrophysical relevance.

6.1. Effects of Boundary Conditions

In § 4 we discussed to what extent numerical artifacts might influence our conclusions. One last effect needs to be mentioned, namely, the choice of the boundary conditions. While the boundary conditions in the horizontal direction are prescribed by the inflow, we are free to define the boundaries in the transverse direction. As mentioned in § 4.2, our standard choice is periodic boundary conditions; i.e., material leaving the simulation domain at the bottom reenters at the top and vice versa (the same holds for the tracer particles). This might raise the question of whether the level of turbulence and the amount of cold gas in the simulations is a “closed box” effect in the sense that if the boundaries were open, the compressed gas could leave the simulation domain before cooling down and contributing to the cold gas mass.

In order to assess the effects of this specific boundary choice, we ran a set of models with open boundary conditions in the direction transverse to the inflow; i.e., gas is free to leave the simulation domain. Models s2C20 and s2C30 (Fig. 19, top) have the same box size as models 2C20 and 2C30 (Fig. 2), while in models l2C20 and l2C30 the simulation domain was enlarged in the transverse direction (Fig. 19, bottom) in such a way that the inflow region has the same size as in the square models (e.g., 2C20 and 2C30), but above and below this inflow region we place “inactive” regions into which the compressed gas can stream. In other words, not only are the boundaries in the transverse direction open, but there is no inflow in the $x$-direction in the “inactive” regions (in fact, they are open boundaries; i.e., material is free to leave the box).

The morphologies of models with open and periodic boundaries are slightly different, but not disquietingly so. For the non-periodic models, note that while there is a clear outflow observable in l2C20 and l2C30, its signature is weak in s2C20 and s2C30. The open boundaries are implemented via a constant extrapolation of the last active cells, so that the pressure gradient at the boundary in the transverse direction is constant. However, because for models l2C20 and l2C30 we implement an “inactive” region above and below the inflow region, the gas in the dense compressed region, i.e., the inflow region, sees a pressure gradient along the transverse direction at the boundaries between the “active” and “inactive” regions.

For models l2C20 and l2C30 dense gas is definitely squeezed out between the colliding flows into the “inactive” regions. How does this affect the mass budget? Figure 20 shows the mass per length contained in the same three temperature regimes as in Figure 16, but for models 2C20, s2C20, and l2C20.

Clearly, the total mass (i.e., the sum over the three temperature regimes) increases linearly, dominated by the growth of the cold mass fraction. Only model l2C20 shows some deviation from a linear evolution at later times. The cooling timescale is much shorter than the flow timescale and the sound crossing time, so that the gas is compressed and cooled down before it can feel the boundaries. In other words, gas in the incoming flow cools efficiently, and structures within the slabs form independent of the boundary conditions. We conclude that our discussions of the mass budget and the level of turbulence are nearly independent of the specific choice of boundary conditions.
models, despite the open boundary conditions of models s2C20 and l2C20. For convenience, the mass fraction in the unstable regime has been multiplied by 10. The total mass in the simulation domain increases (more or less) linearly for all models, despite the open boundary conditions of models s2C20 and l2C20.

6.2. Turbulence

The observed line-of-sight velocity dispersions of a few km s\(^{-1}\) in the cold gas are reproduced by the models (Fig. 7). Yet, the internal velocity dispersions are generally subsonic; i.e., it seems the term “supersonic” is not necessarily an accurate description of the hydrodynamical state of the cold gas (Figs. 8 and 9). Hartmann (2002) argues that because of the ages and small spatial dispersions of young stars in Taurus, their velocity dispersions relative to their natal gas are very likely subsonic. The turbulent line widths amount only to a fraction of the inflow velocity.

The subsonic internal velocity dispersions come as a surprise, and the question arises of whether self-gravity could lead to higher internal dispersions. While this is definitely something to test, gravity, to drive internal motions, would have to act on local nonlinear perturbations within the cold gas, for whose growth there simply might not be enough time available (Burkert & Hartmann 2004).

The above result would affect some aspects of the model of turbulence-controlled star formation (see Mac Low & Klessen 2004). There supersonic turbulence leads to nonlinear (isothermal) shock compressions, which in turn trigger (at least partly) fragmentation and subsequent gravitational collapse. By this fragmentation mechanism star formation is “localized” in the sense that before the whole cloud can collapse, the nonlinear density perturbations caused by supersonic turbulence have formed stars (see also Burkert & Hartmann 2004). In the absence of supersonic turbulence in the cold regions, the nonlinear density perturbations must arise during the formation process of the cloud. These perturbations could be envisaged as the cold filaments, which are a consequence of the combination of dynamical and thermal instabilities. The subsonic internal motions are mirrored by the specific kinetic energy modes (Fig. 13) and spectra (Fig. 12). They clearly point to weakly compressible turbulence in the cold gas.

Turbulence is generated with ease in the colliding flow scenario, and the resulting PoMCloC reproduces observational bulk quantities. One possible consequence of turbulence as an initial condition effect is the efficient mixing between warm and cold phases in the complex. There is an ongoing debate about the degree of mixing between atomic and molecular hydrogen in MCCs. While there is no doubt about the intermittent density (and velocity) structure in molecular clouds, what the diffuse, possibly warmer, medium is made of is less clear. Turbulent mixing would point to the CNM, i.e., H\(_{\text{i}}\), while H\(_{2}\) self-shielding and H\(_{2}\) formation timescale arguments would favor a fully molecular cloud. The “openings” in the clouds (no matter whether they are, e.g., channels, gaps, or holes) might play a crucial role in the internal dynamics of the clouds. Hennebelle & Inutsuka (2006) argue that the channels of warm material in MCCs allow energy (in their case, in the form of Alfvén waves) to enter the cloud and thus drive the internal motions. However, this would only refer to motions of the cold gas regions with respect to each other, but not necessarily to internal motions in the cold gas.

6.3. Thermal States

Because of its very short thermal timescales, the dense gas closely follows the thermal equilibrium relation between pressure and density (Fig. 14). This gas phase is approximately isothermal and corresponds to the dense, isolated regions that we identify as precursors of molecular gas.

A perceptible amount of gas resides in the thermally unstable phase, consistent with observations by Heiles & Troland (2003) and consistent with the modeling results of other groups (e.g., Dib & Burkert 2005; Vázquez-Semadeni et al. 2006). The thermally unstable gas generally travels at high velocities (Fig. 15), meaning that it corresponds to gas that has passed the bounding shock, but has not yet “fallen onto” the cold, dense objects (where it would cool down to the isothermal cloud temperature). The amount of unstable gas varies with the inflow Mach number, which is a consequence of competing dynamical and thermal timescales.

6.4. Conditions for H\(_{2}\) Formation

Given a sufficiently high inflow speed, colliding flows are an efficient way to assemble enough cold gas to form molecular clouds. Despite the highly structured objects, the amount of cold mass increases approximately linearly with time as if collected in a cold slab (Fig. 16). Once the gas has cooled below \(T \approx 300\) K in our models, it stays cold (Fig. 18). The cold material has very short thermal timescales; i.e., it is nearly isothermal [in Fig. 14, \((d\log P)/(d\log n) \rightarrow 1\) for \(n > 100\, \text{cm}^{-3}\)].

The main limiting factor for H\(_{2}\) formation in our models is not the temperature of the gas, but the highly time-dependent column density. Because of the dynamical nature of the clouds, gas is continuously reexposed to the ambient radiation field. Most of the gas manages to stay above shielding column densities of \(N > 10^{21}\, \text{cm}^{-2}\) only for 2 Myr at most (Fig. 18). However, these estimates are pessimistic. (1) We only considered extinction by dust, not self-shielding by already formed H\(_{2}\), and (2) molecular clouds might not form “from scratch,” i.e., only from H\(_{\text{i}}\). Instead, the inflows might already contain a substantial amount of H\(_{2}\) (see, e.g., Pringle et al. 2001). This would not only shift the mass budget in favor of H\(_{2}\), but also provide efficient self-shielding early in the molecular cloud formation process.

6.5. Cloud Formation Timescales

Even more than 10 Myr after flow contact time, in all our runs the total mass collected in the cold gas falls short of typical molecular cloud masses by a factor of approximately 10. The average column density essentially evolves one-dimensionally, i.e., \(N(t) = n_{\text{ngf}} \approx 1 \times 10^{20}\, \text{cm}^{-2}/(t/\text{Myr})\) for the fastest inflow speed in our models. Thus, at least 10 Myr are needed to reach an average column density high enough for self-shielding (since the dynamics leads to local focusing of the gas flows, isolated regions might reach this stage earlier, but not markedly so; see Bergin et al. 2004). Two remedies come to mind. (1) The cold mass growth rate depends nearly linearly on the inflow momentum (Fig. 16). However, the observed velocities and densities in the Galactic H\(_{\text{i}}\) limit
the inflow momentum to approximately $n_0 v_0 < 30 \text{ km s}^{-1} \text{ cm}^{-3}$, or, equivalently, $n_0 v_{\text{rot}} \approx 10^{21} \text{ cm}^{-2}$. Thus, molecular clouds might preferentially form in flows, where the ambient atomic density is higher than the average density, i.e., $n_0 = 3 \text{ cm}^{-3}$ and $v_0 = 10 \text{ km s}^{-1}$. Then, the mechanism described above might well dominate the formation of low-mass molecular clouds. More massive objects need either more time or a different parameter regime, e.g., as in galaxy mergers. (2) The models presented here do not account for gravity. Since a slableike geometry is favored by construction in the colliding flow scenario, gravity would lead to collapse in the plane, causing the densities (and column densities) needed for shielding to be rapidly reached. In that sense, any timescales given for conditions favorable for H$_2$ formation are upper estimates. Note that the lateral collapse can set in much later, and if the cloud is large enough and sufficiently dense substructures exist, it might only occur after stars have formed locally.

We would like to stress the point that since the presented models are concerned with the formation of precursors of molecular clouds, the somewhat largish assembly timescales do not constitute a problem for short molecular cloud lifetimes. As Bergin et al. (2004) discussed, the ratio of the PoMClOc assembly timescale to the actual molecular cloud lifetime can be as large as the ratio of atomic to molecular gas (approximately a factor of 4 in the solar neighborhood; Savage et al. 1977; Dame 1993). For a molecular cloud lifetime of say, 3 Myr in the solar neighborhood (Hartmann et al. 2001 and references therein), the accumulation time for the PoMClOc would be expected to range around $\geq 10$ Myr. Obviously, the molecular cloud lifetime starts only once molecules are being formed.

Our models generally run longer than 10 Myr. At an inflow velocity of 10 km s$^{-1}$, this corresponds to a coherent flow of approximately 100 pc, or, strictly speaking, twice that length for a colliding flow. Although higher flow velocities and/or densities (note that at least with the densities we seem to be on the lower side of the parameter range; see, e.g., Bergin et al. 2004) would allow shorter timescales and thus shorter coherence lengths, the question nevertheless arises of whether such coherent flows are realistic. On smaller scales (and higher densities), the interaction of winds of nearby stars might provide a possible source (see, e.g., Churchwell et al. 2004), while on larger scales, the most probable source would be interacting supernova shells. The somewhat longish assembly timescales suggest that the initial densities of the flows are more likely to be a few cm$^{-3}$ to form molecular clouds with flows of order 50–100 pc length. Spiral density waves can also produce coherent flows of the required length, at least in principle.

7. FUTURE WORK

The main purpose of this work was to demonstrate that precursors of molecular cloud complexes forming in colliding flows reproduce observational measures such as line widths and column densities without any recourse to initially imposed perturbations. The combination of dynamical and thermal instabilities efficiently generates the nonlinear substructures necessary for a fast gravitational collapse and a (nearly) instantaneous onset of star formation, thus meeting the short timescales required by observations (e.g., Hartmann 2003). Once the material is cold and dense, it seems difficult to drive highly supersonic motions (unless there is an internal driver, i.e., stars; Vázquez 1990; Balsara 1996; see, however, Miesch & Zweibel 1994), or even to mix it with warm gas (Heitsch et al. 2006). However, the cold and warm gas must be mixed efficiently in order to allow the cold material to move “freely” at velocities that are supersonic with respect to the temperature of the cold gas. Thus, the cold substructures must arise during the formation process of the cloud, from compression and cooling in the warm phase. We demonstrated without any recourse to perturbed initial conditions that colliding flows provide a generic physical mechanism for creating nonlinear structure in PoMClOc’s. In their further evolution (i.e., beyond the models presented here), these nonlinear perturbations will be the seeds for gravitational collapse and star formation. Thus, our models show that from a dynamical point of view, the concept of “short” cloud lifetimes is feasible. Ultimately, one can begin to see how realistic the concept of “turbulent support” of molecular clouds actually is (see, e.g., Ballesteros-Paredes et al. 1999b).

An MCC in the solar neighborhood has an average column density of $N(\text{H}_2) \approx 1.4 \times 10^{21} \text{ cm}^{-2}$ (McKee 1999), coinciding with the dust-shielding column density necessary for H$_2$ formation. Once this value is reached, H$_2$ formation is expected to begin. Within 2 or 3 more Myr, star formation will have occurred, and the cloud will have dispersed again. In this scenario, the average column density of an ensemble of MCCs should be roughly constant, implying a mass-size relation of $M \propto L^2$. The latter would be expected for a Larson (1981) relation of the type $\sigma \propto L^{2/5}$, consistent with observed scalings (e.g., Elmegreen & Scalo 2004).

Clearly, this study leaves us with many unanswered questions. Gravity might lead to global edge effects (Burkert & Hartmann 2004), as well as to local “supersonic” velocities and to additional instabilities (see, e.g., Hunter et al. 1997; Huchtstaedt et al. 2005). Actual H$_2$ formation would further lower the temperature in the cold gas, in which case supersonic velocities might eventually be reached. The inflows might already be partially molecular, in which case the filling factor of H$_2$ should be low (Pringle et al. 2001). Finally, magnetic fields could influence and possibly control the structure formation (e.g., Vázquez-Semadeni et al. 1995; Passot et al. 1995; Elmegreen 1999; Hartmann et al. 2001).

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